LHC phenomenology of lowest massive Regge recurrences in the Randall-Sundrum orbifold

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We consider string realizations of the Randall-Sundrum effective theory for electroweak symmetry breaking and explore the search for the lowest massive Regge excitation of the gluon and of the extra (color singlet) gauge boson inherent in D-brane constructions. In these curved backgrounds, the higher-spin Regge recurrences of standard model fields localized near the IR brane are warped down to close to the TeV range and hence can be produced at collider experiments. Assuming that the theory is weakly coupled, we make use of four gauge boson amplitudes evaluated near the first Regge pole to determine the discovery potential of LHC. We study the inclusive dijet mass spectrum in the central rapidity region $|y_{\rm jet}| < 1.0$ for dijet masses $M \geq 2.5$ TeV. We find that with an integrated luminosity of 100 fb^{-1}, the 5σ discovery reach can be as high as 4.7 TeV. Observations of resonant structures in $pp \to$ directγ + jet can provide interesting corroboration for string physics up to 3.0 TeV. We also study the ratio of dijet mass spectra at small and large scattering angles. We show that with the first fb^{-1} such a ratio can probe lowest-lying Regge states for masses ~2.5 TeV.

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I. INTRODUCTION

The saga of the standard model (SM) is still exhilarating because it leaves all questions of consequence unanswered. Perhaps the most evident of unanswered questions is why the weak interactions are weak. The well-known nonzero vacuum expectation value of the scalar Higgs doublet condensate, $\langle H \rangle = v = 2.46 \times 10^2$ GeV, sets the scale of electroweak interactions. However, due to the quadratic sensitivity of the Higgs mass to quantum corrections from an arbitrarily high mass scale, one is faced with the gauge hierarchy problem: the question of why $v \ll M_{\rm Pl}$, where $M_{\rm Pl} = 1.22 \times 10^{19}$ GeV is the Planck mass. The traditional view is to adopt $M_{\rm Pl}$ as the fundamental mass setting the scale of the unified theory incorporating gravity and attempt to derive v through some dynamical mechanism (e.g. renormalization group evolution). In recent years, however, a new framework with a diametrically opposite viewpoint has been proposed, in which v is instead the fundamental scale of nature [1]. D-brane string compactifications with low string scale and large extra dimensions allow a definite representation of this innovative premise [2].

TeV-scale superstring theory provides a brane-world description of the SM, which is localized on membranes extending in $p+3$ spatial dimensions, the so-called D-branes [3]. Gauge interactions emerge as excitations of open strings with endpoints attached on 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 $(p+3)$-branes and transverse dimensions). The apparent weakness of gravity at energies below a few TeV can then be understood as a consequence of the gravitational force "leaking" into the transverse large compact dimensions of spacetime. This is possible only if the intrinsic scale of string excitations is also of order a few TeV. Should nature be so cooperative, a whole tower of infinite string excitations will open up at this low mass threshold, and new particles of spin J follow the well-known Regge trajectories of vibrating strings: $J = J_0 + \alpha' M^2$, where α' is the Regge slope parameter that determines the fundamental string mass scale

$$M_s = \frac{1}{\sqrt{\alpha'}}. \tag{1}$$

Only one assumption is necessary to build up a solid framework: the string coupling must be small for the validity of perturbation theory in the computations of scattering amplitudes. In this case, black hole production and other strong gravity effects occur at energies above the string scale, therefore at least the few lowest Regge recurrences are available for examination, free from interference with some complex quantum gravitational phenomena.

In proton collisions at the Large Hadron Collider (LHC), Regge states will be produced as soon as the energies of some partonic subprocesses cross the threshold at $\hat{s} > M_s^2$. In a series of recent publications [4–9] we have computed open string scattering amplitudes in D-brane models and have discussed the associated phenomenological aspects of low mass string Regge recurrences related to experimental searches for physics beyond the SM [10]. We developed our program in the simplest way, by working within the construct of a minimal model in which we considered scattering processes which take place on the (color) $U(3)_a$ stack of D-branes, which is intersected by the (weak doublet) $U(2)_b$ stack of D-branes, as well by a third (weak singlet) $U(1)_c$ stack of D-brane. In the bosonic sector, the open strings terminating on the $U(3)_a$ stack contain the standard gluons g and an additional $U(1)_a$ gauge boson C; on the $U(2)_b$ stacks the open strings correspond to the weak gauge bosons W, and again an additional $U(1)_b$ gauge field. So the associated gauge groups for these stacks are $SU(3)_C \times U(1)_a$, $SU(2)_L \times U(1)_b$, and $U(1)_c$, respectively; the physical hypercharge is a linear combination of $U(1)_a$, $U(1)_b$, and $U(1)_c$. The fermionic matter consists of open strings, which stretch between different stacks of D$(p+3)$-branes and are hence located at the intersection points.

In canonical D-brane constructions the large hierarchy between the weak scale and the fundamental scale of gravity is eliminated through the large volume of the transverse dimensions. An alternative explanation to solve the gauge hierarchy problem was suggested by Randall and Sundrum (herein RS) [11]. The RS setup has the shape of a gravitational condenser: two branes, which rigidly reside at $S^1/\mathbb{Z}_2$ orbifold fixed point boundaries $y = 0$ and $y = \pi r_c$ (the UV and IR branes, respectively), gravitationally repel each other and are stabilized by a slab of anti–de Sitter (AdS) space. The metric satisfying this ansatz (in horospherical coordinates) is given by

$$ds^2 = e^{-2k|y|}\eta_{\mu\nu}dx^\mu dx^\nu + dy^2, \tag{2}$$

where k is the AdS curvature scale, which is somewhat smaller than the fundamental five-dimensional Planck mass $M^\star_{\rm Pl} \sim M_{\rm Pl}$.[1] In this setup the distance scales get exponentially redshifted as one moves from the UV brane toward the IR brane. Such exponential suppression can then naturally explain why the observed physical scales are so much smaller than the Planck scale. For example, if the five-dimensional Higgs condensate $v_5 \sim k$ is IR-localized, the observed four-dimensional value will be obtained from $e^{-k\pi r_c}\langle H_5\rangle$, and the observed hierarchy between the gravitational and electroweak mass scales is reproduced if $kr_c \approx 12$. The most distinct signal of this setup is the appearance of a tower of spin-2 resonances, corresponding to the Kaluza-Klein (KK) states of the five-dimensional graviton, which have masses and couplings driven by the TeV scale. These KK gravitons couple to all SM fields universally, yielding striking predictions for collider experiments [12].

As originally noted in [13], to address the hierarchy problem it is sufficient to keep the Higgs near the IR brane. Interestingly, if the remaining gauge bosons and fermions are allowed to propagate into the warped dimension, one can also formulate an attractive mechanism to explain the flavor mass hierarchy [14]. The idea here is that the light fermions are localized near the UV brane. This raises the effective cutoff scale for operators composed of these fields far above the TeV regime, providing an efficient mechanism to suppress unwanted operators, such as those mediating flavor changing neutral currents (FCNC) processes, related to tightly constrained light flavors. Moreover, this results in small four-dimensional Yukawa couplings to the Higgs, even if there are no small five-dimensional Yukawa couplings. The top quark is IR-localized to obtain a large four-dimensional top Yukawa coupling. Because the fermion profiles depend exponentially on the bulk masses, this provides an understanding of the hierarchy of fermion masses (and mixing) without hierarchies in the fundamental five-dimensional parameters, solving the SM flavor puzzle.

The RS setup has also been used to construct warped Higgsless models, where the electroweak symmetry is broken by boundary conditions on the five-dimensional gauge fields [15]. Gauge fields are allowed to propagate within all five dimensions. The electroweak gauge structure of the minimal viable model is $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, where $U(1)_{B-L}$ corresponds to gauging baryon minus lepton number. Boundary conditions on the bulk gauge fields are chosen so that the $SU(2)_L \times SU(2)_R$ symmetry is broken on the IR brane to the diagonal subgroup $SU(2)_D$, and the $SU(2)_R \times U(1)_{B-L}$ symmetry is broken down to the usual $U(1)_Y$ hypercharge in the UV brane to ensure that the low-energy gauge group without electroweak symmetry breaking is $SU(2)_L \times U(1)_Y$. The $SU(3)_C$ QCD group is unbroken everywhere, i.e., in the warped dimension and on the branes. The spectrum of electroweak vector bosons consists of a single massless photon along with KK towers of charged W_n and neutral Z_n states. The SM massive W and Z vectors, which get masses from the $SU(2)_L \times U(1)_Y$-violating boundary condition on the IR brane, are identified with the lowest KK modes of the W_n and Z_n towers. SM fermions extend into all dimensions, and they have explicit mass terms that are allowed by the nonchiral structure of the theory in the bulk and on the IR brane. The most serious challenge to construct viable models of Higgsless electroweak symmetry breaking is satisfying the constraints from precision electroweak measurements [16]. Mixing of the W and Z with higher KK modes changes their couplings to fermions relative to the SM. Heavier KK modes ($m_{\rm KK} \gtrsim 1$ TeV) are preferred to reduce these deviations to an acceptable level, but the KK modes cannot be too heavy ($m_{\rm KK} \lesssim 1$ TeV) if they are to unitarize vector boson scattering. Both requirements can be satisfied simultaneously if there are localized kinetic terms on each of the branes [17], and if the SM fermions (with the exception of the right-handed top quark) have approximately flat profiles in the extra dimension [18]. In this case, the first vector boson KK modes above the Z and W typically have masses ≈ 0.5–1.5 TeV [18]. Additional model structure is needed to generate a sufficiently large

[1]Greek subscripts extend over ordinary four-dimensional spacetime and are raised and lowered with the flat Minkowskian metric $\eta_{\mu\nu}$, whereas Latin subscripts span the full five-dimensional space and are raised and lowered with the full metric g_{MN}.

top quark mass while not overly disrupting the measured $Zb_L\bar{b}_L$ coupling. Some examples include new toplike custodial bulk fermions [19], or a second warped bulk space on the other side of the UV brane with its own IR brane [20].

In string realizations of extended RS models of hierarchy and flavor, we do expect the higher-spin Regge recurrences of the SM fields localized near the IR brane to be redshifted close to the TeV scale and therefore be directly produced at the LHC [21,22]. In this paper, we explore the search for the lowest massive Regge excitation in inclusive γ + jet and dijet mass spectra. In the spirit of [21–23], we *assume* that the RS orbifold arises as part of the compactification manifold in a weakly coupled string theory. We further *assume* that the compactification radii of the other five dimensions are $\mathcal{O}(M_s^{\star -1})$ and therefore can be safely integrated out.[2] With this in mind, the basic relation between the curvature of the warped internal space, the string scale (the mass of the Regge states), and the five-dimensional Planck mass is

$$k \ll M_s^\star = \frac{1}{\sqrt{\alpha'^\star}} \ll M_{\rm Pl}^\star, \quad (3)$$

where $\alpha'^\star$ is the slope of the associated Regge trajectory. The first inequality permits the warping to leave intact the basic string properties (such as the dual resonant structure) of the perturbative scattering amplitudes. The infinite tower of open string Regge excitations have the same quantum numbers under the SM gauge group as the gluons and the quarks, but in general higher spins, and their masses are just square-root-of-integer multiples of the string mass $M_s^\star$. In what follows, the first Regge excitations of the gluon (g), the extra $U(1)$ boson tied to the color stack (C), and quarks (q) will be indicated with g^*, C^*, q^*, respectively. In this paper we complement model independent searches of top-production via q^* excitation [21,25] by analyzing tree-level four-point amplitudes relevant to γ + jet and dijet final states. We make use of four gauge boson amplitudes evaluated near the first resonant pole to determine the discovery potential of LHC for g^* and C^* excitations. We study the inclusive dijet mass spectrum in the central rapidity region $|y_{\rm jet}| < 1.0$ for dijet masses $M \geq 2.5$ TeV. We find that with an integrated luminosity of 100 fb^{-1}, the 5σ discovery reach can be as high as 4.7 TeV. Observations of resonant structures in $pp \to$ directγ + jet can provide interesting corroboration for string physics up to 3.0 TeV. We also study the ratio of dijet mass spectra at small and large (center-of-mass) scattering angles. We show that with the first fb^{-1} such a ratio can probe lowest-lying Regge states for masses $\sim$2.5 TeV. The outline of the paper is as follows: in Sec. II we collect all the relevant formulas leading to four gauge boson string amplitudes, the analysis of the LHC is carried out in Sec. III, and we summarize in Sec. IV.

[2]The dearth of string constructions for a transition to the RS compactification [24] makes a full comparison between the string scale and internal dimension radii difficult. A recent study [23] of a range of models seems to indicate that $M_s^\star r_c \sim 1$ is viable.

II. FOUR-POINT AMPLITUDES OF GAUGE BOSONS

The most direct way to compute the amplitude for the scattering of four gauge bosons is to consider the case of polarized particles because all nonvanishing contributions can be then generated from a single, maximally helicity violating (MHV) amplitude—the so-called *partial* MHV amplitude [26]. In canonical D-brane constructions, all string effects are encapsulated in this amplitude in one ‘‘form factor’’ function of Mandelstam variables s, t, u (constrained by $s + t + u = 0$)[3]:

$$V(s,t,u) = \frac{su}{tM_s^2} B(-s/M_s^2, -u/M_s^2) = \frac{\Gamma(1-s/M_s^2)\Gamma(1-u/M_s^2)}{\Gamma(1+t/M_s^2)}. \quad (4)$$

The physical content of the form factor becomes clear after using the well-known expansion in terms of s-channel resonances [27]

$$B(-s/M_s^2, -u/M_s^2) = -\sum_{n=0}^{\infty} \frac{M_s^{2-2n}}{n!} \frac{1}{s - nM_s^2} \bigg[\prod_{J=1}^{n}(u + M_s^2 J)\bigg], \quad (5)$$

which exhibits s-channel poles associated to the propagation of virtual Regge excitations with masses $\sqrt{n}M_s$. Thus near the nth level pole ($s \to nM_s^2$)

$$V(s,t,u) \approx \frac{1}{s - nM_s^2} \times \frac{M_s^{2-2n}}{(n-1)!} \prod_{J=0}^{n-1}(u + M_s^2 J). \quad (6)$$

In specific amplitudes, the residues combine with the remaining kinematic factors, reflecting the spin content of particles exchanged in the s-channel, ranging from $J = 0$ to $J = n + 1$. Unfortunately, Veneziano amplitudes only apply to strings propagating on flat Minkowski backgrounds, and their generalization to warped spaces is presently unknown. In the absence of concrete string theory constructions, we describe the lowest-lying Regge excitations of SM gauge bosons following the bottom-up approach advocated in [22]. In the limit where k is taken to zero, this innovative approach reproduces the string effects encapsulated in (4).

[3]For simplicity, we drop carets for the parton subprocesses.

Consider a free (noninteracting) massive spin-2 field B_{MN} in curved five-dimensional spacetime,

$$\mathcal{L} = \frac{1}{4}H^{LMN}H_{LMN} - \frac{1}{2}H^{LM}{}_{M}H_{LN}{}^{N} + \frac{1}{2}m^2[(B_M{}^M)^2 - B^{MN}B_{MN}], \tag{7}$$

where $H_{LMN} = \nabla_L B_{MN} - \nabla_M B_{LN}$ is the field strength tensor and $m \equiv M_s^\star$ is the mass of the lightest Regge excitation. This field can be further decomposed according to its spins ($J = 0$, $J = 1$, and $J = 2$) in four dimensions. The tensor, vector, and scalar components are $B_{\mu\nu}$, $B_{\mu 5}$, and B_{55}, respectively. The Lagrangian (7) contains terms which mix these components. Such mixed terms need to be canceled for a consistent KK decomposition. As shown in [22], the action can be factorized as

$$S = S_{J=2} \oplus S_{J=1,J=0}, \tag{8}$$

where the five-dimensional Lagrangian for $J = 2$ is given by

$$S_{J=2} = \int d^5x \bigg\{ e^{2k|y|}\bigg[\frac{1}{4}H^{\lambda\mu\nu}H_{\lambda\mu\nu} - \frac{1}{2}\Big(1 - \frac{2}{\xi}\Big)H^{\lambda\mu}{}_{\mu}H_{\lambda\nu}{}^{\nu}\bigg] + \frac{1}{2}B_\mu{}^\mu(-\partial_y^2 + 4k^2 + m^2)B_\nu{}^\nu - \frac{1}{2}B^{\mu\nu}(-\partial_y^2 + 4k^2 + m^2)B_{\mu\nu} + 2k[\delta(y) - \delta(y - \pi r_c)][B^{\mu\nu}B_{\mu\nu} - (B_\mu{}^\mu)^2]\bigg\}, \tag{9}$$

and ξ is a parameter in the gauge fixing term. The field $B_{\mu\nu}$ can be decomposed according to its wave function in the warped dimension

$$B_{\mu\nu} = \frac{1}{\sqrt{\pi r_c}} \sum_{n=1}^{\infty} B_{\mu\nu}^{(n)} f^{(n)}(y). \tag{10}$$

The equation of motion is

$$e^{2k|y|}D_{\mu\nu}{}^{\alpha\beta}B_{\alpha\beta} + \{-\partial_y^2 + 4k^2 + m^2 - 4k[\delta(y) - \delta(y - \pi r_c)]\}B_{\mu\nu} = 0, \tag{11}$$

where $D_{\mu\nu}{}^{\alpha\beta}$ is an operator from the first two terms in (9). A massless spin-2 field has the equation of motion of $D_{\mu\nu}{}^{\alpha\beta}B_{\alpha\beta} = 0$. So the masses are given by the eigenvalues of the operator

$$e^{-2k|y|}\{-\partial_y^2 + 4k^2 + m^2 - 4k[\delta(y) - \delta(y - \pi r_c)]\}, \tag{12}$$

with mode functions $f^{(n)}$ satisfying the following equation:

$$-f^{(n)\prime\prime} + (4k^2 + m^2)f^{(n)} - 4k[\delta(y) - \delta(y - \pi r_c)]f^{(n)} = (\mu^{(n)})^2 e^{2k|y|} f^{(n)}, \tag{13}$$

and associated inner product

$$\frac{1}{\pi r_c}\int_0^{\pi r_c} dy e^{2k|y|} f^{(n)} f^{(m)} = \delta^{nm}, \tag{14}$$

from the orthonormal condition. For this choice of $f^{(n)}$, we have [from the last three terms in (9)]

$$\begin{aligned}\int d^5x \ldots &= -\frac{1}{2}\int d^4x dy B^{\mu\nu}\{-\partial_y^2 + 4k^2 + m^2 - 4k[\delta(y) - \delta(y - \pi r_c)]\}B_{\mu\nu} \\ &= -\frac{1}{2}\int d^4x dy B^{(m)\mu\nu}(x)B_{\mu\nu}^{(n)}(x)\frac{1}{\pi r_c} \times \sum_{n=1}^{\infty}\sum_{m=1}^{\infty}(\mu^{(n)})^2 e^{2k|y|}f^{(m)}(y)f^{(n)}(y) \\ &= -\frac{1}{2}\int d^4x \sum_{n=1}^{\infty}(\mu^{(n)})^2 B^{(n)\mu\nu}(x)B_{\mu\nu}^{(n)}(x),\end{aligned} \tag{15}$$

where in the last line, we use (14). The integration of $H_{\lambda\mu\nu}H^{\lambda\mu\nu}$ is trivial because there is no y-derivative. Hence, after the extra dimension is integrated out, Eq. (9) can be reduced to a four-dimensional Lagrangian of free spin-2 fields (with different masses $\mu^{(n)}$)

$$S_{J=2} = \int d^4x \sum_{n=1}^{\infty}\bigg\{\frac{1}{4}H^{(n)\lambda\mu\nu}H_{\lambda\mu\nu}^{(n)} - \frac{1}{2}\Big(1 - \frac{2}{\xi}\Big)H^{(n)\lambda\mu}{}_{\mu}H_{\lambda\nu}^{(n)\nu} + \frac{1}{2}(\mu^{(n)})^2[B_\mu^{(n)\mu}B_\nu^{(n)\nu} - B^{(n)\mu\nu}B_{\mu\nu}^{(n)}]\bigg\}, \tag{16}$$

where $\xi \to \infty$ when computing the scattering amplitude. The general solution of (13) is a Bessel function [22]

$$f^{(n)}(y) = \frac{1}{N}\bigg[J_\nu\Big(\frac{\mu^{(n)}}{\Lambda_{\rm IR}}w\Big) + cJ_{-\nu}\Big(\frac{\mu^{(n)}}{\Lambda_{\rm IR}}w\Big)\bigg], \tag{17}$$

where N is the normalization constant, c is an integration constant (each of these constants implicitly depends upon the level n), $\Lambda_{\rm IR} = ke^{-\pi k r_c}$, and $w = e^{k(|y| - \pi r_c)}$, $\in [e^{-k\pi r_c}, 1]$. The order of the Bessel function is $\nu \equiv \sqrt{4 + \mathfrak{m}^2}$, where $\mathfrak{m} = m/k$ is the string scale in units of the RS curvature. With appropriate boundary conditions, the masses $\mu^{(n)}$ and the explicit form of $f^{(n)}$ can be obtained.

We now turn to the discussion of $J = 0$. In the effective four-dimensional theory there is one real scalar $\text{Re}(\phi)$, which comes from the five-dimensional scalar and couples to the gluon strength F^2. In addition, there is one pseudoscalar axion $A_\star^5$, which in four dimensions couples as $A_\star^5 F^* F$, with ${}^*F = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$. This pseudoscalar axion comes from the fifth component of a massive vector $A_\star$, with coupling $\epsilon_{\mu\nu\rho\sigma 5}F^{\mu\nu}F^{\rho\sigma}A_\star^5$. Then, $\text{Re}(\phi)$ and $\text{Im}(\phi) \equiv A_\star^5$ combine to one complex scalar ϕ which couples as $\phi(F + i^*F)(F + i^*F) + \text{c.c.}$; this ensures that ϕ and its complex conjugate ϕ^* couple only to the

$++$ $(--)$ helicity combinations, respectively.[4] Both the scalar and the pseudoscalar will be affected in the same way by warping, because they sit in one supersymmetry multiplet. Thus, to determine the $J = 0$ contribution, we study the effect of warping on a dilatonlike scalar a with the coupling $\mathrm{Re}(\phi)F^2$.

The Klein-Gordon equation for a scalar ϕ in the RS spacetime is

$$\frac{1}{\sqrt{g}}\partial_M\sqrt{g}\partial^M\phi + m^2\phi = 0; \tag{18}$$

more explicitly, it is

$$e^{2k|y|}\partial^\mu\partial_\mu\phi + [-\partial_y^2 + 4k\,\mathrm{sgn}(y)\partial_y + m^2]\phi = 0. \tag{19}$$

The field ϕ can be decomposed according to its wave function in the warped dimension

$$\phi(x, y) = \frac{1}{\sqrt{\pi r_c}}\sum_{n=1}^{\infty}\phi^{(n)}(x)h^{(n)}(y). \tag{20}$$

One can choose the mode functions $h^{(n)}$ satisfying the following equation:

$$-h^{(n)\prime\prime} + 4k\,\mathrm{sgn}(y)h^{(n)\prime} + m^2h^{(n)} = (\mu^{(n)})^2e^{2k|y|}h^{(n)}. \tag{21}$$

With a change of variable $x = \frac{1}{k}e^{k|y|}$, we have

$$\frac{d}{dy} = kx\frac{d}{dx}, \qquad \frac{d^2}{dy^2} = k^2x^2\frac{d^2}{dy^2} + k^2x\frac{d}{dx}, \tag{22}$$

so (21) can be written as

$$x^2h^{(n)\prime\prime} + 3xh^{(n)\prime} + [(\mu^{(n)})^2x^2 - \mathfrak{m}^2](\mu^{(n)}) = 0. \tag{23}$$

The solution to this equation is

$$h^{(n)}(x) = \frac{1}{N}(\mu^{(n)}x)^2\{J_\nu(\mu^{(n)}x) + CJ_{-\nu}(\mu^{(n)}x)\} \equiv x^2\tilde{f}^{(n)}(x), \tag{24}$$

where N is a normalization constant and C an integration constant. For later convenience, we also define a new function $\tilde{f}^{(n)}$. The boundary conditions are

$$h^{(n)\prime}(0+) - h^{(n)\prime}(0-) = 0 \tag{25}$$

and

$$h^{(n)\prime}(-\pi r_c+) - h^{(n)\prime}(\pi r_c-) = 0, \tag{26}$$

where the prime is the derivative with respect to y. As in the case of $B_{\mu\nu}$, the mass $\mu^{(n)}$ is determined from the second boundary condition

$$x^2\tilde{f}^{(n)\prime}(-\pi r_c+) - 2xkx\tilde{f}^{(n)}(-\pi r_c+) - x^2\tilde{f}^{(n)\prime}(\pi r_c-) - 2xkx\tilde{f}^{(n)}(\pi r_c-) = 0, \tag{27}$$

or

$$\tilde{f}^{(n)\prime}(-\pi r_c+) - \tilde{f}^{(n)\prime}(\pi r_c-) = 4k\tilde{f}^{(n)}(\pi r_c), \tag{28}$$

which is essentially the boundary condition for $B_{\mu\nu}$ [22]. As a result, the mass of ϕ is exactly the same as that of $B_{\mu\nu}$. Note that $h^{(n)}(x)$ can be expressed as

$$h^{(n)} = e^{2k|y|}f^{(n)}, \tag{29}$$

where $f^{(n)}$ are the mode functions for $B_{\mu\nu}$. So $h^{(n)}(x)$ are normalized as

$$\frac{1}{\pi r_c}\int_0^{\pi r_c} dye^{-2k|y|}h^{(n)}h^{(m)} = \delta^{nm}. \tag{30}$$

This gives a canonical kinetic term for $\phi^{(n)}$ (because of the different powers of $e^{2k|y|}$).

In this paper we will restrict our calculations to incoming QCD gluons. We then obtain the decomposition of the QCD gauge field. Gauge freedom can be used to set $A_5 = 0$ [28]. This is consistent with the gauge invariant equation $\oint dx^5A_5 = 0$, which results from the assumption that A_5 is a $\mathbb{Z}_2$-odd function of the extra dimension. In this gauge, the four-dimensional vector zero mode has a constant profile in the bulk

$$A_\mu(x, y) = \frac{1}{\sqrt{\pi r_c}}A_\mu^{(0)}(x) + \cdots, \tag{31}$$

and the gluon field strength takes the familiar form $F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + g_af^{abc}A_\mu^bA_\nu^c$, with $a = 1, \ldots, 8$.

The coupling of the five-dimensional field B_{MN} to the gluon is given by

$$S_{ggg^*(C^*)} = \int d^5x\sqrt{-g}\frac{g_5}{\sqrt{2}M_s^\star}C^{abc}\Big(F^{aAC}F_C^{bB} - \frac{1}{4}F^{aCD}F_{CD}^bg^{AB}\Big)B_{AB}^c, \tag{32}$$

where $C^{abc} = 2[\mathrm{Tr}(T^aT^bT^c) + \mathrm{Tr}(T^aT^bT^c)]$ is the color factor, T^a are the generators of the fundamental representation of $U(3)$ (normalized here according to $\mathrm{Tr}(T^aT^b) = \frac{1}{2}\delta^{ab}$), and $F_{AB}^a = \partial_AA_B^a - \partial_BA_A^a + g_5f^{abc}A_A^bA_B^c$. Note that the color indices on the field strength F run from 1 to 8; on the tensor field B, $U(3)$ indices ($c = 0, \ldots, 8$) are permitted (with $c = 0$ corresponding to the tensor excitation C^*).[5] Hence, g_5 is related to the Yang-Mills QCD coupling g_a according to $g_5 = g_a\sqrt{\pi r_c}$. The factor $g_5/\sqrt{2}M_s^\star$ is determined by matching the $gg \to g^*(C^*)$ amplitude to the

[4]We may trace the origin of the $J = 0$ contribution to components of B_{MN} and other fields of the 10-dimensional theory. Instead, we proceed by simply using the correspondence with the tree-level string theory and identify the vertex function through comparison with the tree-level $J = 0$ pole. As described in the text this has the correct helicity structure.

[5]As can be verified from the four-point function [4] there is no coupling $gg \to C$, however the composite nature of C^* and g^* permits, respectively, $gg \to C^*$ and $gC \to g^*$ couplings, with color globally preserve.

s-channel pole term in the string (tree-level) amplitude [29]. Thus, the four-dimensional coupling term is found to be

$$\mathcal{L}_{gg\to g^*(C^*)} = \frac{g^{(0)}}{\sqrt{2}\tilde{M}_s} C^{abc}\Big[\Big(F^{a\gamma}F_{\gamma}{}^{\beta} - \frac{1}{4}F^{a\gamma\delta}F^{b}_{\gamma\delta}\eta^{\alpha\beta}\Big)B^{c(0)}_{\alpha\beta} + \frac{1}{2}\Big(\phi^{c(0)}F^{a\mu\nu}F^{b}_{\mu\nu} + \frac{1}{2}\bar{\phi}^{c(0)}F^{a\mu\nu}F^{b\rho\sigma}\epsilon_{\mu\nu\rho\sigma}\Big)\Big], \quad (33)$$

where $\tilde{M}_s = e^{-k\pi r_c}M_s^\star \sim 1$ TeV is the redshifted string scale, $g^{(0)}$ follows from the integration of the zero mode $f^{(0)}(y)$ of $B^{(0)}_{\mu\nu}$, and $\bar{\phi}^{c(0)}$ is the zero mode for the imaginary part of the complex scalar. Since each field in (32) contributes to the integration with a factor $(\pi r_c)^{-1/2}$ we obtain

$$g^{(0)} = \frac{g_a e^{-\pi k r_c}}{\pi r_c}\int_0^{\pi r_c} dy e^{2ky} f^{(0)}(y). \quad (34)$$

The coupling (33) gives three vertices,

$$i\frac{\sqrt{2}g^{(0)}}{\tilde{M}_s}C^{abc}\Big(\Sigma^{\alpha\beta} - \frac{1}{4}\eta^{\alpha\beta}\Sigma_{\gamma}{}^{\gamma}\Big)b_{\alpha\beta}, \qquad i\frac{g^{(0)}}{\sqrt{2}\tilde{M}_s}C^{abc}\Sigma_{\mu}{}^{\mu},$$
$$i\frac{g^{(0)}}{\sqrt{2}\tilde{M}_s}C^{abc}4\epsilon_{\mu\nu\rho\sigma}k_1^\mu\epsilon_1^\nu k_2^\rho\epsilon_2^\sigma, \quad (35)$$

where $b_{\alpha\beta}$ is a polarization of $B^{c(0)}_{\alpha\beta}$, k_i^μ and ϵ_i^ν (with $i = 1, 2$) are, respectively, the momentum and polarization of the incoming gluons, $\Sigma^{\alpha\beta} = (k_1^\alpha\epsilon_1^\gamma - k_1^\gamma\epsilon_1^\alpha)(k_{2\gamma}\epsilon_2^\beta - k_2^\beta\epsilon_{2\gamma}) + (\alpha \leftrightarrow \beta)$, and its trace $\Sigma_\gamma{}^\gamma = 4(k_1\cdot\epsilon_2)(k_2\cdot\epsilon_1) - 4(\epsilon_1\cdot\epsilon_2)(k_1\cdot k_2)$ [30]. As in the $J = 2$ case, the coupling is determined by matching to the $J = 0$ pole term in the tree-level string amplitude.

Finally, we note that the $J = 1$ resonant level exists, but is not accessible in purely gluonic scattering [5].

The s-channel pole terms of the average square amplitudes contributing to γ + jet and dijet production at the LHC can be obtained from the general formulae given in Ref. [6]. The 4-gluon average square amplitude is given by

$$|\mathcal{M}(gg\to gg)|^2 = 2\Big(\frac{g^{(0)}}{\tilde{M}_s}\Big)^4\Big(\frac{N^2 - 4 + (12/N^2)}{N^2 - 1}\Big)\frac{s^4 + t^4 + u^4}{(s-\mu^2)^2}, \quad (36)$$

where to simplify notation we have dropped the superscript indicating the lowest massive Regge excitation, i.e., $\mu \equiv \mu^{(0)}$. 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 [5]. After this is done, the contribution of $gg \to gg$ is as follows:

$$|\mathcal{M}(gg\to gg)|^2 = \frac{19}{12}\Big(\frac{g^{(0)}}{\tilde{M}_s}\Big)^4\Big\{W^{gg\to gg}_{g^*}\Big[\frac{s^4}{(s-\mu^2)^2 + (\Gamma^{J=0}_{g^*}\mu)^2} + \frac{t^4+u^4}{(s-\mu^2)^2 + (\Gamma^{J=2}_{g^*}\mu)^2}\Big] + W^{gg\to gg}_{C^*}\Big[\frac{s^4}{(s-\mu^2)^2 + (\Gamma^{J=0}_{C^*}\mu)^2} + \frac{t^4+u^4}{(s-\mu^2)^2 + (\Gamma^{J=2}_{C^*}\mu)^2}\Big]\Big\}, \quad (37)$$

where

$$\Gamma^{J=0}_{g^*} = 75\Big(\frac{g^{(0)}\mu}{g_a\tilde{M}_s}\Big)\Big(\frac{\mu}{\text{TeV}}\Big)\ \text{GeV},$$
$$\Gamma^{J=0}_{C^*} = 150\Big(\frac{g^{(0)}\mu}{g_a\tilde{M}_s}\Big)\Big(\frac{\mu}{\text{TeV}}\Big)\ \text{GeV},$$
$$\Gamma^{J=2}_{g^*} = 45\Big(\frac{g^{(0)}\mu}{g_a\tilde{M}_s}\Big)\Big(\frac{\mu}{\text{TeV}}\Big)\ \text{GeV},$$
$$\Gamma^{J=2}_{C^*} = 75\Big(\frac{g^{(0)}\mu}{g_a\tilde{M}_s}\Big)\Big(\frac{\mu}{\text{TeV}}\Big)\ \text{GeV}$$

are the total decay widths for intermediate states g^*, C^* (with angular momentum J) [5,22]. The associated weights of these intermediate states are given in terms of the probabilities for the various entrance and exit channels

$$\frac{N^2 - 4 + 12/N^2}{N^2-1} = \frac{16}{(N^2-1)^2}\Big[(N^2-1)\Big(\frac{N^2-4}{4N}\Big)^2 + \Big(\frac{N^2-1}{2N}\Big)^2\Big] \propto \frac{16}{(N^2-1)^2}[(N^2-1)(\Gamma_{g^*\to gg})^2 + (\Gamma_{C^*\to gg})^2], \quad (38)$$

yielding

$$W^{gg\to gg}_{g^*} = \frac{8(\Gamma_{g^*\to gg})^2}{8(\Gamma_{g^*\to gg})^2 + (\Gamma_{C^*\to gg})^2} = 0.44,$$
$$W^{gg\to gg}_{C^*} = \frac{(\Gamma_{C^*\to gg})^2}{8(\Gamma_{g^*\to gg})^2 + (\Gamma_{C^*\to gg})^2} = 0.56,$$

where superscripts $J = 2$ are understood to be inserted on all the Γ's.

As we pointed out in the introduction, the hypercharge is a color composite state containing the photon. The s-channel pole term of the average square amplitude contributing to $gg \to \gamma$ + jet is given by [4]

$$|\mathcal{M}(gg \to g\gamma)|^2 = \frac{5}{3} Q^2 \left(\frac{g^{(0)}}{\tilde{M}_s}\right)^4 \left[\frac{s^4}{(s-\mu^2)^2 + (\Gamma_{g^*}^{J=0}\mu)^2} + \frac{t^4 + u^4}{(s-\mu^2)^2 + (\Gamma_{g^*}^{J=2}\mu)^2}\right], \tag{39}$$

where $Q = \sqrt{1/6}\kappa \cos\theta_W$ is the product of the $U(1)$ charge of the fundamental representation ($\sqrt{1/6}$) followed by successive projections onto the hypercharge (κ) and then onto the photon ($\cos\theta_W$). The $C - Y$ mixing coefficient is model dependent: in the minimal $U(3) \times Sp(1) \times U(1)$ model it is quite small, around $\kappa \simeq 0.12$ for couplings evaluated at the Z mass [31], which is modestly enhanced to $\kappa \simeq 0.14$ as a result of renormalization-group running of the couplings up to 2.5 TeV [4]. It should be noted that in models possessing an additional $U(1)$ which partners $SU(2)_L$ on a $U(2)$ brane [32], the various assignment of the charges can result in values of κ which can differ considerably from 0.12. For the phenomenological analysis that follows we set $\kappa^2 = 0.02$.

III. LHC DISCOVERY REACH

The most important parameter to determine the LHC discovery reach for string recurrences is the mass of the lowest-lying Regge excitation, which depends on $\Lambda_{\rm IR}$ and $\mathfrak{m}$. For fixed $\mathfrak{m}$ the mass of g^* and C^* excitations is to a very good approximation a linear function of $\Lambda_{\rm IR}$ [22]. As we already remarked in the introduction, in Higgsless models $\Lambda_{\rm IR}$ is subject to significant constraints from electroweak data. The KK excitations of the vector gauge bosons must be near 1 TeV to simultaneously satisfy unitarity and electroweak constraints. This leads to $\Lambda_{\rm IR} \approx 0.5$ TeV. Similarly, to avoid precision electroweak constraints in scenarios where the Higgs is IR-localized the lightest KK excitation mass is $\gtrsim 3$ TeV [33], yielding $\Lambda_{\rm IR} \gtrsim 1$ TeV. From (3) we obtain the condition $\mathfrak{m} \gg 1$ for string propagation on a smooth geometric background. Nevertheless, as in many examples in various arenas of physics, $\mathfrak{m} \sim$ a few may in fact be sufficient, depending on the behavior of the leading corrections to the geometric limit. In our phenomenological study we will follow [22] and set $\mathfrak{m} \gtrsim 3$, which leads to $\mu^{(0)} \approx 5\Lambda_{\rm IR}$, $g^{(0)}/g_a \simeq 0.1$, and $\mu^{(0)} = 5\mathfrak{m}^{-1}\tilde{M}_s \simeq 1.7\tilde{M}_s$.

Given the particular nature of the process we are considering, the production of a TeV Regge state and its subsequent 2-body decay, one would hope that the resonance would be visible in data binned according to the invariant mass M of the dijet, after setting cuts on the different jet rapidities, $|y_1|$, $|y_2| \le 1$ [34] and transverse momenta $p_{\rm T}^{1,2} > 50$ GeV. With the definitions $Y \equiv \frac{1}{2}(y_1 + y_2)$ and $y \equiv \frac{1}{2}(y_1 - y_2)$, the cross section per interval of M for $pp \to$ dijet is given by

$$\frac{d\sigma}{dM} = M\tau \sum_{ijkl}\left[\int_{-Y_{\rm max}}^{0} dY f_i(x_a, M) f_j(x_b, M) \times \int_{-(y_{\rm max}+Y)}^{y_{\rm max}+Y} dy \frac{d\sigma}{d\hat{t}}\bigg|_{ij\to kl} \frac{1}{\cosh^2 y} + \int_0^{Y_{\rm max}} dY f_i(x_a, M) f_j(x_b, M) \times \int_{-(y_{\rm max}-Y)}^{y_{\rm max}-Y} dy \frac{d\sigma}{d\hat{t}}\bigg|_{ij\to kl} \frac{1}{\cosh^2 y}\right], \tag{40}$$

where $\tau = M^2/s$, $x_a = \sqrt{\tau}e^Y$, $x_b = \sqrt{\tau}e^{-Y}$, and

$$|\mathcal{M}(ij \to kl)|^2 = 16\pi\hat{s}^2 \frac{d\sigma}{d\hat{t}}\bigg|_{ij\to kl}. \tag{41}$$

In this section we reinstate the caret notation ($\hat{s}$, $\hat{t}$, $\hat{u}$) to specify partonic subprocesses. The Y integration range in Eq. (40), $Y_{\rm max} = \min\{\ln(1/\sqrt{\tau}), y_{\rm max}\}$, comes from requiring $x_a, x_b < 1$ together with the rapidity cuts $y_{\rm min} < |y_1|$, $|y_2| < y_{\rm max}$. The kinematics of the scattering also provides the relation $M = 2p_T\cosh y$, which when combined with $p_T = M/2\sin\theta^* = M/2\sqrt{1-\cos^2\theta^*}$, yields $\cosh y = (1-\cos^2\theta^*)^{-1/2}$, where θ^* is the center-of-mass scattering angle. Finally, the Mandelstam invariants occurring in the cross section are given by $\hat{s} = M^2$, $\hat{t} = -\frac{1}{2}M^2e^{-y}/\cosh y$, and $\hat{u} = -\frac{1}{2}M^2e^{+y}/\cosh y$.

Standard bump-hunting methods, such as calculating cumulative cross sections

$$\sigma(M_0) = \int_{M_0}^{\infty} \frac{d\sigma}{dM} dM \tag{42}$$

and searching for regions with significant deviations from the QCD background, may allow one to find an interval of M suspected of containing a bump. With the establishment of such a region, one may calculate a signal-to-noise ratio, with the signal rate estimated in the invariant mass window $[\mu^{(0)} - 2\Gamma, \mu^{(0)} + 2\Gamma]$. As usual, the noise is defined as the square root of the number of background events in the same dijet mass interval for the same integrated luminosity. The QCD background has been calculated at the partonic level considering all SM contributions to dijet final states [7]. Our calculation, making use of the CTEQ6 parton distribution functions [35] agrees with that presented in [34].

The top curve in Fig. 1 shows the behavior of the signal-to-noise (S/N) ratio as a function of the lowest massive Regge excitation, for 100 fb^{-1} of integrated luminosity and $\sqrt{s} = 14$ TeV. *Regge excitations with masses* $\mu^{(0)} \lesssim 4.7$ TeV *are open to discovery at the* $\geq 5\sigma$ *level*. This implies that in the Higgsless model discovery would be possible in a wide range of the presently unconstrained parameter space, whereas in the model with a Higgs localized on the IR-brane the LHC discovery potential would be only marginal. The bottom curve in Fig. 1 shows the S/N ratio in the $pp \to$ direct γ + jet channel. To accommodate the minimal acceptance cuts on final state photons from the CMS and ATLAS proposals [36], we set $|y_{\rm max}| < 2.4$.

The approximate equality of the background due to misidentified π^0's and the QCD background [37], across a range of large p_T^γ as implemented in Ref. [4], is maintained as an approximate equality over a range of γ-jet invariant masses with the rapidity cuts imposed. Observations of resonant structures in $pp \to \text{direct}\,\gamma + \text{jet}$ can provide interesting corroboration for string physics up to 3.0 TeV. Before proceeding, we stress that the results shown in Fig. 1 are conservative, in the sense that we have not included in the signal the stringy contributions of processes containing fermions. These will be somewhat more model dependent since they require details of the SM pattern of masses and mixings, but we expect that these contributions can potentially increase the reach of LHC for discovery of Regge recurrences.

QCD parton-parton cross sections are dominated by t-channel exchanges that produce dijet angular distributions which peak at small center-of-mass scattering angles. In contrast, nonstandard contact interactions or excitations of resonances result in a more isotropic distribution. In terms of rapidity variable for standard transverse momentum cuts, dijets resulting from QCD processes will preferentially populate the large rapidity region, while the new processes generate events more uniformly distributed in the entire rapidity region. To analyze the details of the rapidity space the D0 Collaboration introduced a new parameter [38],

$$R = \frac{d\sigma/dM|_{(|y_1|,|y_2|<0.5)}}{d\sigma/dM|_{(0.5<|y_1|,|y_2|<1.0)}}, \tag{43}$$

the ratio of the number of events, in a given dijet mass bin, for both rapidities $|y_1|$, $|y_2| < 0.5$ and both rapidities

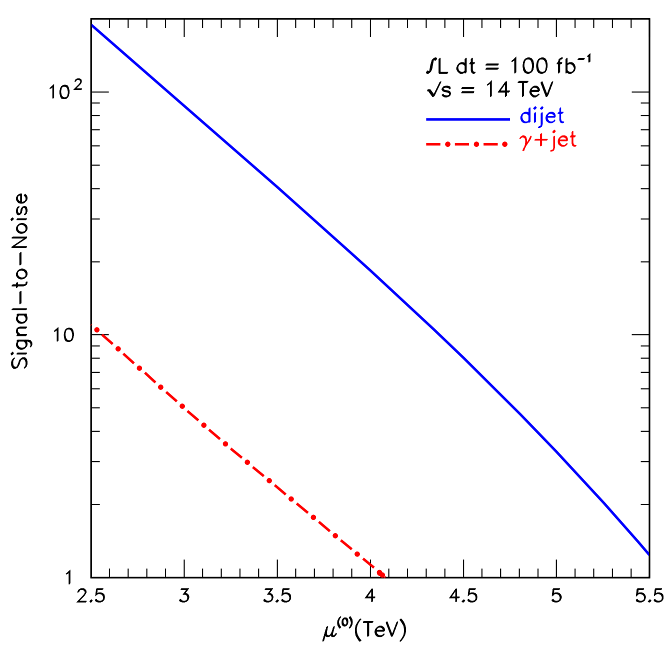

FIG. 1 (color online). $pp \to$ dijet and $pp \to \gamma +$ jet signal-to-noise ratio for 100 fb^{-1} integrated luminosity.

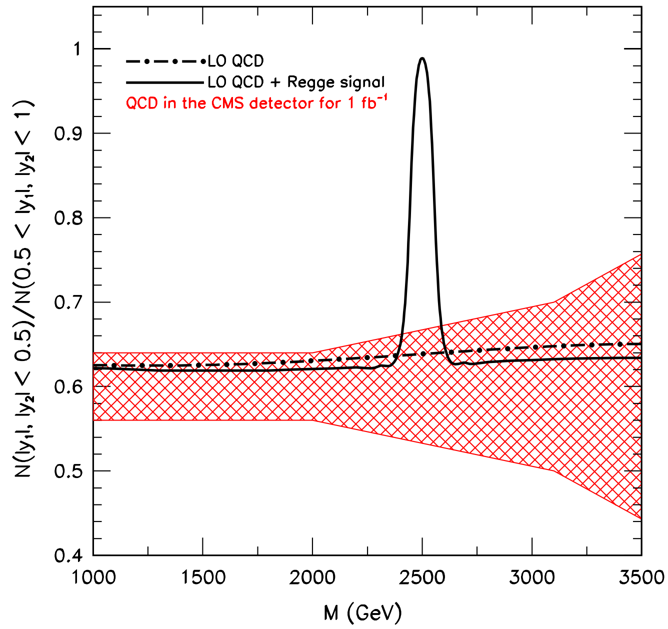

FIG. 2 (color online). For a luminosity of 100 fb^{-1}, the expected statistical error (shaded region) of the dijet ratio of QCD in the CMS detector [40] is compared with LO QCD (dot-dashed line) and LO QCD plus lowest massive Regge excitation (solid line), for $\mu^{(0)} = 2.5$ TeV.

$0.5 < |y_1|$, $|y_2| < 1.0$. The ratio R is a genuine measure of the most sensitive part of the angular distribution, providing a single number that can be measure as a function of the dijet invariant mass [39].

In Fig. 2 we compare the results from a full CMS detector simulation of the ratio R, with predictions from leading-order (LO) QCD and contributions to the g^* and C^* excitations. The synthetic population was generated with PYTHIA, passed through the full CMS detector simulation and reconstructed with the ORCA reconstruction package [40]. It is clear that with the first fb^{-1} of data collected at the LHC, the R parameter will be able to probe lowest-lying Regge excitations for $\mu^{(0)} \sim 2.5$ TeV.[6]

IV. CONCLUSIONS

In this paper, we have extended the work in Refs. [21,22] on an approximate calculation of string amplitudes in the RS geometry to include the $J = 0$ contribution to bosonic four-point functions. We have carried out a phenomenological analysis of the resonant contributions to dijet

[6]It should be noted that the R parameter serves only as the crudest discriminator between QCD and stringy behavior of the cross section. More detailed analyses of the rapidity dependence of the final state jets are in order. In a recent paper [41] the behavior of the stringy amplitudes (for flat geometries) with respect to the rapidity difference y has been discussed. Results were presented for the separate contributions of the 1/2 and 3/2 resonances for the dominant $qg \to qg$ process, as well as for the combined cross sections. It remains to compare these to QCD.

production at the LHC, and found that for an integrated luminosity of 100 fb^{-1}, discovery of the resonant signal at signal-to-noise of 5σ is possible for resonant masses of up to nearly 5 TeV. However, it should be noted that this is possible only for the Higgsless model. For the model with the Higgs on or near the IR brane, the requirement $\Lambda_{\mathrm{IR}} \geq 1$ TeV combined with the relation $\mu \simeq 5\Lambda_{\mathrm{IR}}$ implies $\mu > 5$ TeV, greatly narrowing the possible region of discovery.

In addition to the Regge recurrences there are of course KK modes of SM particles and gravitons propagating in the s-channel, which at this point we have omitted from consideration. Their importance can be gauged by their masses relative to $\mu^{(0)}$. The ratio of string to KK masses is model dependent, but in general there could be several cases where the $\mu^{(0)}/m_{\mathrm{KK}}$ ratio is around a few [23]. This relation can be illustrated by comparing with the masses of the KK states of the graviton: $m_G^{(n)} = x_n \Lambda_{\mathrm{IR}}$, with x_n being the nth roots of the Bessel function J_1 [12]. We find that $\mu^{(0)}/m_G^{(1)} \sim 1.25$. This implies that the KK contribution is not significantly enhanced over the Regge contribution [21], and so here we have limited our discussion to the Regge case.

Finally, the large amount of data required for discovery may be traced to a strong difference at the phenomenological level between the RS scenario and the flat space result: the effective four-dimensional coupling constant $g^{(0)} \simeq 0.1 g_a$. For a given resonance mass, we also have $\tilde{M}_s \simeq 0.6\mu$. The net result, following from Eq. (37), is that for a given resonance mass, the RS cross section is a factor of $(0.1/0.6)^4 \approx 10^{-3}$ times that of the flat case scenario. (There is also some effect from the narrowing of the total widths.) The drastic reduction of the effective coupling is a direct result of permitting the gluon field to propagate in the warped bulk.

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[1] N. Arkani-Hamed, S. Dimopoulos, and G. R. Dvali, Phys. Lett. B **429**, 263 (1998).

[2] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. R. Dvali, Phys. Lett. B **436**, 257 (1998); G. Shiu and S. H. H. Tye, Phys. Rev. D **58**, 106007 (1998).

[3] R. Blumenhagen, B. Körs, D. Lüst, and S. Stieberger, Phys. Rep. **445**, 1 (2007).

[4] L. A. Anchordoqui, H. Goldberg, S. Nawata, and T. R. Taylor, Phys. Rev. Lett. **100**, 171603 (2008); Phys. Rev. D **78**, 016005 (2008);

[5] L. A. Anchordoqui, H. Goldberg, and T. R. Taylor, Phys. Lett. B **668**, 373 (2008).

[6] D. Lüst, S. Stieberger, and T. R. Taylor, Nucl. Phys. **B808**, 1 (2009).

[7] L. A. Anchordoqui, H. Goldberg, D. Lüst, S. Nawata, S. Stieberger, and T. R. Taylor, Phys. Rev. Lett. **101**, 241803 (2008); Nucl. Phys. **B821**, 181 (2009).

[8] D. Lust, O. Schlotterer, S. Stieberger, and T. R. Taylor, Nucl. Phys. **B828**, 139 (2010).

[9] L. A. Anchordoqui, H. Goldberg, D. Lust, S. Stieberger, and T. R. Taylor, Mod. Phys. Lett. A **24**, 2481 (2009).

[10] We have also shown that supersymmetric D-brane constructions embrace an interesting dark matter phenomenology: L. A. Anchordoqui, H. Goldberg, D. Hooper, D. Marfatia, and T. R. Taylor, Phys. Lett. B **683**, 321 (2010).

[11] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999).

[12] H. Davoudiasl, J. L. Hewett, and T. G. Rizzo, Phys. Rev. Lett. **84**, 2080 (2000).

[13] W. D. Goldberger and M. B. Wise, Phys. Rev. D **60**, 107505 (1999).

[14] Y. Grossman and M. Neubert, Phys. Lett. B **474**, 361 (2000); T. Gherghetta and A. Pomarol, Nucl. Phys. B **586**, 141 (2000); S. J. Huber and Q. Shafi, Phys. Lett. B **498**, 256 (2001).

[15] C. Csaki, C. Grojean, L. Pilo, and J. Terning, Phys. Rev. Lett. **92**, 101802 (2004); C. Csaki, C. Grojean, J. Hubisz, Y. Shirman, and J. Terning, Phys. Rev. D **70**, 015012 (2004).

[16] Y. Nomura, J. High Energy Phys. 11 (2003) 050; R. Barbieri, A. Pomarol, and R. Rattazzi, Phys. Lett. B **591**, 141 (2004); H. Davoudiasl, J. L. Hewett, B. Lillie, and T. G. Rizzo, Phys. Rev. D **70**, 015006 (2004); J. High Energy Phys. 05 (2004) 015; J. L. Hewett, B. Lillie, and T. G. Rizzo, J. High Energy Phys. 10 (2004) 014; G. Burdman and Y. Nomura, Phys. Rev. D **69**, 115013 (2004).

[17] G. Cacciapaglia, C. Csaki, C. Grojean, and J. Terning, Phys. Rev. D **70**, 075014 (2004).

[18] G. Cacciapaglia, C. Csaki, C. Grojean, and J. Terning, Phys. Rev. D **71**, 035015 (2005).

[19] K. Agashe, R. Contino, L. Da Rold, and A. Pomarol, Phys. Lett. B **641**, 62 (2006); G. Cacciapaglia, C. Csaki, G. Marandella, and J. Terning, Phys. Rev. D **75**, 015003 (2007).

[20] G. Cacciapaglia, C. Csaki, C. Grojean, M. Reece, and J. Terning, Phys. Rev. D **72**, 095018 (2005).
[21] B. Hassanain, J. March-Russell, and J. G. Rosa, J. High Energy Phys. 07 (2009) 077.
[22] M. Perelstein and A. Spray, J. High Energy Phys. 10 (2009) 096.
[23] M. Reece and L. T. Wang, J. High Energy Phys. 07 (2010) 040.
[24] H. L. Verlinde, Nucl. Phys. **B580**, 264 (2000); S. B. Giddings, S. Kachru, and J. Polchinski, Phys. Rev. D **66**, 106006 (2002); S. Kachru, D. Simic, and S. P. Trivedi, J. High Energy Phys. 05 (2010) 067.
[25] The discovery potential of string resonances via top quark pair production in the context of canonical D-brane constructions has been recently established: Z. Dong, T. Han, M. x. Huang, and G. Shiu, J. High Energy Phys. 09 (2010) 048.
[26] S. J. Parke and T. R. Taylor, Phys. Rev. Lett. **56**, 2459 (1986).
[27] G. Veneziano, Nuovo Cimento A **57**, 190 (1968).
[28] H. Davoudiasl, J. L. Hewett, and T. G. Rizzo, Phys. Lett. B **473**, 43 (2000); A. Pomarol, Phys. Lett. B **486**, 153 (2000).
[29] See Ref. [22] for some caveats pertaining to this approach.
[30] S. Cullen, M. Perelstein, and M. E. Peskin, Phys. Rev. D **62**, 055012 (2000).
[31] D. Berenstein and S. Pinansky, Phys. Rev. D **75**, 095009 (2007).
[32] See, e.g., I. Antoniadis, E. Kiritsis, and T. N. Tomaras, Phys. Lett. B **486**, 186 (2000).
[33] K. Agashe, A. Delgado, M. J. May, and R. Sundrum, J. High Energy Phys. 08 (2003) 050.
[34] A. Bhatti *et al.*, J. Phys. G **36**, 015004 (2009).
[35] J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. Nadolsky, and W. K. Tung, J. High Energy Phys. 07 (2002) 012.
[36] G. L. Bayatian *et al.* (CMS Collaboration), J. Phys. G **34**, 995 (2007); W. W. Armstrong *et al.* (ATLAS Collaboration), Report No. CERN/LHCC 94-43 (unpublished).
[37] P. Gupta, B. C. Choudhary, S. Chatterji, and S. Bhattacharya, Eur. Phys. J. C **53**, 49 (2007).
[38] B. Abbott *et al.* (D0 Collaboration), Phys. Rev. Lett. **82**, 2457 (1999).
[39] An illustration of the use of this parameter in a heuristic model where standard model amplitudes are modified by a Veneziano form factor has been presented by P. Meade and L. Randall, J. High Energy Phys. 05 (2008) 003.
[40] S. Esen and R. Harris, CMS Note 2006/071, 2006 (unpublished).
[41] N. Kitazawa, J. High Energy Phys. 10 (2010) 051.