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Warped Phenomenology *

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Abstract

We explore the phenomenology associated with the recently proposed localized gravity model of Randall and Sundrum where gravity propagates in a 5-dimensional non-factorizable geometry and generates the 4-dimensional weak-Planck scale hierarchy by an exponential function of the compactification radius, called a warp factor. The Kaluza-Klein tower of gravitons which emerge in this scenario have strikingly different properties than in the factorizable case with large extra dimensions. We derive the form of the graviton tower interactions with the Standard Model fields and examine their direct production in Drell-Yan and dijet events at the Tevatron and LHC as well as the KK spectrum line-shape at high-energy linear e^+e^- colliders. In the case where the first KK excitation is observed, we outline the procedure to uniquely determine the parameters of this scenario. We also investigate the effect of KK tower exchanges in contact interaction searches. We find that present experiments can place meaningful constraints on the parameters of this model.

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The large disparity between the electroweak and apparent fundamental scale of gravity, known as the hierarchy problem, is a primary mystery of particle physics. Traditionally, new symmetries, particles, or interactions have been introduced at the electroweak scale to stabilize this hierarchy. However, it is possible that our 4-dimensional vision of gravity does not represent the full theory and that the observed value of the Planck scale, M_{Pl} is not truly fundamental. A scenario of this type due to Arkani-Hamed, Dimopoulos, and Dvali[1] (ADD) proposes the existence of n additional compact dimensions and relates the fundamental $4+n$ dimensional Planck scale, M, to our effective 4-dimensional value through the volume of the compactified dimensions, $M_{Pl}^2 = V_n M^{2+n}$. Setting $M \sim$ TeV to remove the above hierarchy necessitates a large size for the extra dimensions with a compactification scale of $\mu_c = 1/r_c \sim$ eV$-$MeV for $n = 2-7$. This, unfortunately, introduces another hierarchy between μ_c and M, which must somehow be stabilized. Nonetheless, this scenario has received much attention as it affords concrete phenomenological tests. Since it is experimentally determined that the Standard Model (SM) fields do not feel the effects of additional dimensions of this size, they are confined to a wall, or 3-brane, while gravity is allowed to propagate freely in the full higher-dimensional space, or bulk. Kaluza-Klein (KK) towers of gravitons, which can interact with the wall fields, result from compactification of the bulk. The coupling of each KK excitation is M_{Pl} suppressed, however the mode spacing is determined by μ_c and is thus very small compared to typical collider energies. This allows the summation over an enormous number of KK states which can be exchanged or emitted in a physical process, thereby reducing the summed suppression from $1/M_{Pl}$ to $1/M$, or $\sim$TeV^{-1}. This has resulted in a vast array of phenomenological[2] and astrophysical[3] studies with present collider data bounding $M \gtrsim 1$ TeV for all n and Supernova 1987A cooling and γ ray flux constraints setting $M \gtrsim 50-110$ TeV for $n = 2$ only.

An alternative higher dimensional scenario has recently been proposed by Randall

and Sundrum[4] (RS), where the hierarchy is generated by an exponential function of the compactification radius, called a warp factor. They assume a 5-dimensional non-factorizable geometry, based on a slice of AdS_5 spacetime. Two 3-branes, one being visible with the other being hidden, with opposite tensions reside at S_1/Z_2 orbifold fixed points, taken to be $\phi = 0, \pi$, where ϕ is the angular coordinate parameterizing the extra dimension. The solution to Einstein's equations for this configuration, maintaining 4-dimensional Poincare invariance, is given by the 5-dimensional metric

$$ds^2 = e^{-2\sigma(\phi)}\eta_{\mu\nu}dx^\mu dx^\nu + r_c^2 d\phi^2 \,, \tag{1}$$

where the Greek indices run over ordinary 4-dimensional spacetime, $\sigma(\phi) = kr_c|\phi|$ with r_c being the compactification radius of the extra dimension, and $0 \leq |\phi| \leq \pi$. Here k is a scale of order the Planck scale and relates the 5-dimensional Planck scale M to the cosmological constant. Similar configurations have also been found to arise in M/string-theory[5]. An extension of this scenario where the higher dimensional space is non-compact, *i.e.*, $r_c \to \infty$, is discussed in Ref. [6] and several aspects of this and related ideas have been investigated in Ref. [7]. Examination of the action in the 4-dimensional effective theory in the RS scenario yields[4]

$$\overline{M}_{Pl}^2 = \frac{M^3}{k}(1 - e^{-2kr_c\pi}) \tag{2}$$

for the reduced effective 4-D Planck scale. Assuming that we live on the 3-brane located at $|\phi| = \pi$, it is found that a field on this brane with the fundamental mass parameter m_0 will appear to have the physical mass $m = e^{-kr_c\pi}m_0$. TeV scales are thus generated from fundamental scales of order M_{Pl} via a geometrical exponential factor and the observed scale hierarchy is reproduced if $kr_c \simeq 12$. Hence, due to the exponential nature of the warp factor, no additional large hierarchies are generated. In fact, it has been demonstrated[8] that the size of μ_c in this scenario can be stabilized without fine tuning of parameters, making this

theory very attractive.

The graviton KK spectrum is quite different in this scenario than in the case with factorizable geometry, resulting in a distinctive phenomenology. As we will see below, the masses and couplings of each individual KK excitation are determined by the scale $\Lambda_\pi = \overline{M}_{Pl} e^{-k r_c \pi} \sim$ TeV. This implies that these KK states can be separately produced on resonance with observable rates at colliders up to the kinematic limit. We will examine the cases of KK graviton production in Drell-Yan and dijet events at hadron colliders as well as the KK spectrum line-shape at high-energy linear e^+e^- colliders. In the circumstance where a resonance is observed, we outline the procedure to be employed in order to uniquely determine the parameters of this model. In the case where no direct production is observed, we compute the bounds on the parameter space in the contact interaction limit. We find that data from present accelerators already place meaningful constraints on the parameter space of this scenario. The phenomenology of these KK gravitons is similar in spirit to the production of traditional KK excitations of the SM gauge fields[9], but differs in detail because of the form of the KK wavefunction due to the non-factorizable metric and their spin.

We now calculate the mass spectrum and couplings of the graviton KK modes in the effective 4-dimensional theory on the 3-brane at $\phi = \pi$. The starting point is the 5-dimensional Einstein's equation for the RS configuration, which is given in Ref. [4]. We parameterize the tensor fluctuations $h_{\alpha\beta}$ by taking a linear expansion of the flat metric about its Minkowski value, $\hat{G}_{\alpha\beta} = e^{-2\sigma} (\eta_{\alpha\beta} + \kappa^* h_{\alpha\beta})$, where κ^* is an expansion parameter. In order to obtain the mass spectrum of the tensor fluctuations, we consider the 4-dimensional $\alpha\beta$ components of Einstein's equation with the replacement $G_{\alpha\beta} \to \hat{G}_{\alpha\beta}$, keeping terms up to $\mathcal{O}(\kappa^*)$. We work in the gauge with $\partial^\alpha h_{\alpha\beta} = h^\alpha_\alpha = 0$. Upon compactification the graviton

field $h_{\alpha\beta}$ is expanded into a KK tower

$$h_{\alpha\beta}(x,\phi) = \sum_{n=0}^{\infty} h_{\alpha\beta}^{(n)}(x)\,\frac{\chi^{(n)}(\phi)}{\sqrt{r_c}}\,, \tag{3}$$

where the $h_{\alpha\beta}^{(n)}(x)$ correspond to the KK modes of the graviton on the background of Minkowski space on the 3-brane. In a gauge where $\eta^{\alpha\beta}\partial_\alpha h_{\beta\gamma}^{(n)} = \eta^{\alpha\beta} h_{\alpha\beta}^{(n)} = 0$, the equation of motion of $h_{\alpha\beta}^{(n)}$ is given by

$$\left(\eta^{\alpha\beta}\partial_\alpha\partial_\beta - m_n^2\right) h_{\mu\nu}^{(n)}(x) = 0\,, \tag{4}$$

corresponding to the states with masses $m_n \geq 0$. Using the KK expansion (3) for $h_{\alpha\beta}$ in $\hat{G}_{\alpha\beta}$, Einstein's equation in conjunction with the above equation of motion yields the following differential equation for $\chi^{(n)}(\phi)$

$$\frac{-1}{r_c^2}\frac{d}{d\phi}\left(e^{-4\sigma}\frac{d\chi^{(n)}}{d\phi}\right) = m_n^2\, e^{-2\sigma}\chi^{(n)}\,. \tag{5}$$

The orthonormality condition for $\chi^{(n)}$ is found to be $\int_{-\pi}^{\pi} d\phi\; e^{-2\sigma}\chi^{(m)}\chi^{(n)} = \delta_{mn}$. In deriving Eq. (5), we have used $(d\sigma/d\phi)^2 = (kr_c)^2$ and $d^2\sigma/d\phi^2 = 2kr_c\,[\delta(\phi) - \delta(\phi-\pi)]$, as required by the orbifold symmetry for $\phi \in [-\pi,\pi]$ [4]. The solutions for $\chi^{(n)}$ are then given by[10]

$$\chi^{(n)}(\phi) = \frac{e^{2\sigma}}{N_n}\left[J_2(z_n) + \alpha_n\, Y_2(z_n)\right], \tag{6}$$

where J_2 and Y_2 are Bessel functions of order 2, $z_n(\phi) = m_n e^{\sigma(\phi)}/k$, N_n represents the wavefunction normalization, and α_n are constant coefficients.

Defining $x_n \equiv z_n(\pi)$, and working in the limit that $m_n/k \ll 1$ and $e^{kr_c\pi} \gg 1$, the requirement that the first derivative of $\chi^{(n)}$ be continuous at the orbifold fixed points yields

$$\alpha_n \sim x_n^2 e^{-2kr_c\pi}\,, \qquad \text{and} \qquad J_1(x_n) = 0\,, \tag{7}$$

so that the x_n are simply roots of the Bessel function of order 1. Note that the masses of the graviton KK excitations, given by $m_n = kx_n e^{-kr_c\pi}$, are dependent on the roots of J_1 and are not equally spaced, contrasted to most KK models with one extra dimension. For $x_n \ll e^{kr_c\pi}$, we see that $\alpha_n \ll 1$, and hence $Y_2(z_n)$ can be neglected compared to $J_2(z_n)$ in Eq. (6). We thus obtain for the normalization

$$N_n \simeq \frac{e^{kr_c\pi}}{\sqrt{kr_c}}\, J_2(x_n)\ ;\ n > 0\,, \tag{8}$$

and the normalization of the zero mode is simply $N_0 = 1/\sqrt{kr_c}$.

Having found the solutions for $\chi^{(n)}$, we can now derive the interactions of $h^{(n)}_{\alpha\beta}$ with the matter fields on the 3-brane. Starting with the 5-dimensional action and imposing the constraint that we live on the brane at $\phi = \pi$, we find the usual form of the interaction Lagrangian in the 4-dimensional effective theory,

$$\mathcal{L} = -\frac{1}{M^{3/2}} T^{\alpha\beta}(x) h_{\alpha\beta}(x, \phi = \pi)\,, \tag{9}$$

where $T_{\alpha\beta}(x)$ is the symmetric conserved Minkowski space energy-momentum tensor of the matter fields and we have used the definition $\kappa^* = 2/M^{3/2}$. Expanding the graviton field into the KK states of Eq. (3) and using the above normalization in Eq. (8) for $\chi^{(n)}(\phi)$ we find via Eq. (2)

$$\mathcal{L} = -\frac{1}{\overline{M}_{Pl}} T^{\alpha\beta}(x) h^{(0)}_{\alpha\beta}(x) - \frac{1}{\Lambda_\pi} T^{\alpha\beta}(x) \sum_{n=1}^{\infty} h^{(n)}_{\alpha\beta}(x)\,. \tag{10}$$

Here we see that the zero mode separates from the sum and couples with the usual 4-dimensional strength, $\overline{M}_{Pl}^{-1}$, however, all the massive KK states are only suppressed by Λ_π^{-1}, where we find that $\Lambda_\pi = e^{-kr_c\pi}\overline{M}_{Pl}$, which is of order the weak scale.

Our calculations have been performed with the assumption $k < M$ with $M \sim \overline{M}_{Pl}$, so that the 5-dimensional curvature is small compared to M and the solution for the bulk metric can be trusted[4]. This implies that the ratio $k/\overline{M}_{Pl}$ cannot be too large and we take $k/\overline{M}_{Pl} \leq 1$ in our analysis below. As we will see, the value of this ratio is central to the phenomenological investigation of this model. In order to get a feel for the natural size of this parameter, we perform a simple estimate using string theoretic arguments. The string scale M_s can be related[11] to $\overline{M}_{Pl}$ in 4-dimensional heterotic string theories by $M_s \sim g_{_{YM}} \overline{M}_{Pl}$, where $g_{_{YM}}$ is the 4-dimensional Yang-Mills gauge coupling constant, and the tension τ_3 of a D 3-brane is given by

$$\tau_3 = \frac{M_s^4}{g\,(2\pi)^3}\,, \tag{11}$$

where g is the string coupling constant. For $g_{_{YM}} \sim 0.7$ and $g \sim 1$, we find $\tau_3 \sim 10^{-3}\,\overline{M}_{Pl}^4$. In the RS scenario, the magnitude of the 3-brane tension is given by $V = 24\,\overline{M}_{Pl}^2 k^2$. Requiring that $V = \tau_3$, suggests

$$\frac{k}{\overline{M}_{Pl}} \sim 10^{-2}. \tag{12}$$

We take the range $0.01 \leq k/\overline{M}_{Pl} \leq 1$ in our phenomenological analysis, however, the above discussion suggests that string theoretic and curvature considerations favor the lower end of this range. We note that recent work[12] on gauge unification in a modified RS scenario also favors smaller values for this ratio.

Constraints on the parameters of this model can be obtained by direct collider searches for the first graviton excitation at the Tevatron or LHC. The cleanest signal for graviton resonance production will be either an excess in Drell-Yan events, $q\bar{q}, gg \to G^{(1)} \to \ell^+\ell^-$ (in analogy to searches for extra neutral gauge bosons), or in the dijet channel, $q\bar{q}, gg \to G^{(1)} \to q\bar{q}, gg$. Note that gluon-gluon initiated processes now contribute to Drell-Yan production.

This differs from the ADD scheme where individual resonances associated with graviton exchange are not observable due to the tiny mode spacing. Using the above Lagrangian (10), the production cross section, decay widths, and branching fractions relevant for graviton production can be obtained in a straightforward manner. We assume that the first excitation only decays into SM states, so that for a fixed value of the first graviton excitation mass, m_1, the value of $k/\overline{M}_{Pl}$ completely determines all of the above quantities. In fact, the total width is found to be proportional to $(k/\overline{M}_{Pl})^2$. Keeping in mind that theoretic arguments favor a smaller value for this parameter, and to get a handle on the possible constraints that arise from these channels, we employ the narrow width approximation. This is strictly valid only for values of $k/\overline{M}_{Pl} \lesssim 0.3$ but well approximates the true search reach obtained via a more complete analysis[13]. We then compare our results with the existing Tevatron bounds[14]. The lack of any signal for a new resonance in either the Drell-Yan or dijet channel in the data then provides a constraint on $k/\overline{M}_{Pl}$ for any given value of m_1 as shown in Fig. 1(a). We also perform a similar analysis to estimate the future 95% C.L. parameter exclusion regions at both Run II at the Tevatron and at the LHC under the assumption that no signal is found; these results are displayed in Figs. 1(a) and (b). The dijet constraints for Run II were estimated by a simple luminosity (and $\sqrt{s}$) rescaling of the published Run I results. Note that the Drell-Yan and dijet channels play complementary roles at the Tevatron in obtaining these limits. We expect a dijet search at the LHC to yield poor results due to the large QCD background at this higher center-of-mass energy.

The discovery of the first graviton excitation as a resonance at a collider will immediately allow the determination of all of the fundamental model parameters through measurements of its mass and width, m_1 and Γ_1, respectively. To demonstrate this, we make use of the two relations $\Lambda_\pi = m_1\overline{M}_{Pl}/kx_1$ and $\Gamma_1 = \rho m_1 x_1^2 (k/\overline{M}_{Pl})^2$, where x_1 is the first non-zero

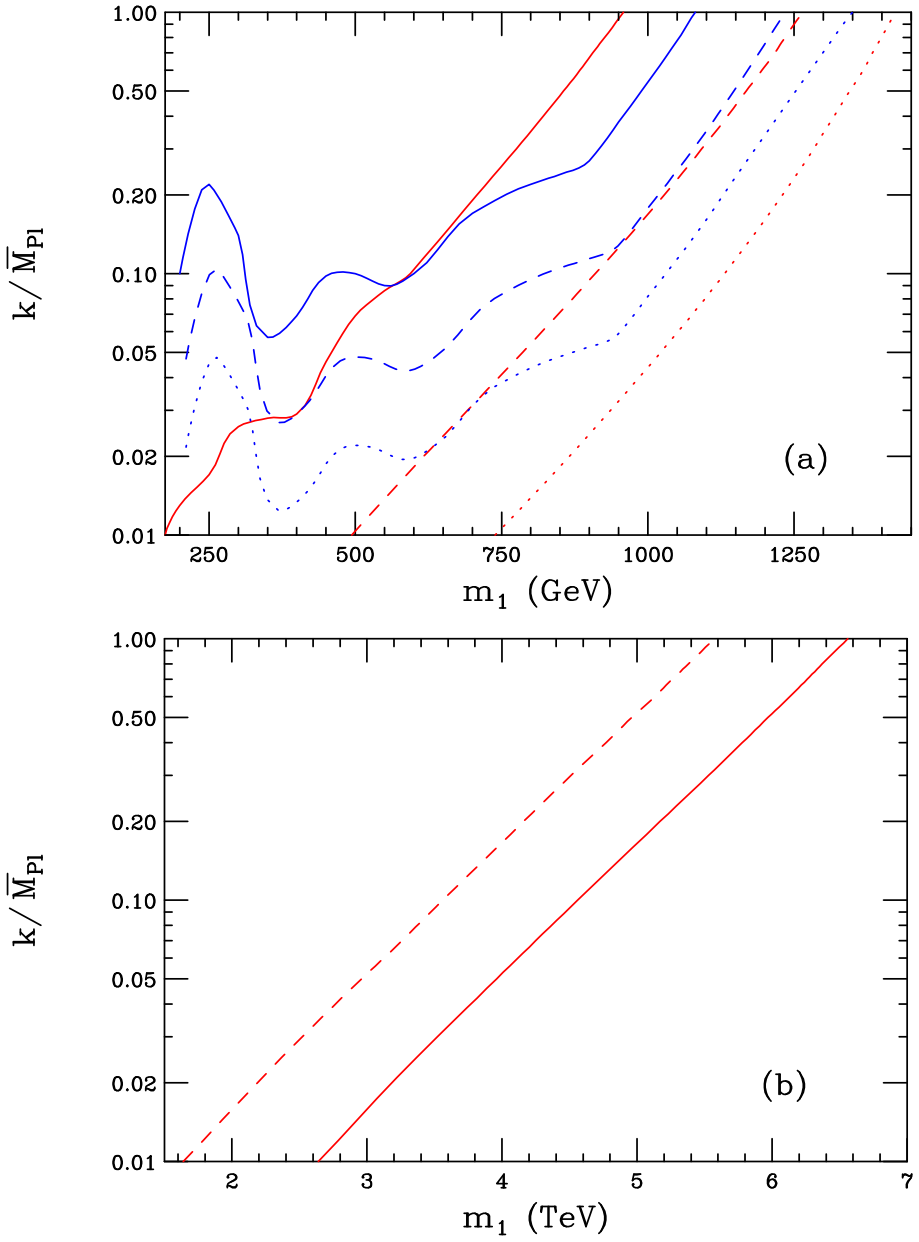

Figure 1: Exclusion regions for resonance production of the first KK graviton excitation in (a) the Drell-Yan (corresponding to the diagonal lines) and dijet (represented by the bumpy curves) channels at the Tevatron and (b) Drell-Yan production at the LHC. (a) The solid curves represent the results for Run I, while the dashed, dotted curves correspond to Run II with 2, 30 fb^{-1} of integrated luminosity, respectively. (b) The dashed, solid curves correspond to 10, 100 fb^{-1} . The excluded region lies above and to the left of the curves.

root of the J_1 Bessel function and ρ is a constant which depends on the number of open decay channels; it is fixed provided we assume that the graviton decays only to SM fields. Using these relations we immediately find that $r_c = -\log[m_1/kx_1]/k\pi$ with $k = \overline{M}_{Pl}[\Gamma_1/m_1\rho x_1^2]^{1/2}$. In addition, the spin-2 nature of the graviton can be determined via angular distributions of its decay products.

To exhibit how the tower of graviton excitations may appear at a collider, Fig. 2 displays the cross section for $e^+e^- \to \mu^+\mu^-$ as a function of $\sqrt{s}$, assuming $m_1 = 600$ GeV and taking various values of $k/\overline{M}_{Pl}$ for purposes of demonstration. We see that for small values of $k/\overline{M}_{Pl}$ the gravitons appear as ever widening peaks and are almost regularly spaced, with the widths and the spacing both being dependent on successive roots of J_1. However, as $k/\overline{M}_{Pl}$ grows, the peaks become too wide to be identified as true resonances and the classic KK signature of successive peaks becomes lost. Instead, it would appear experimentally that there is an overall large enhancement of the cross section, similar to what might be expected from a contact interaction. One may worry that at some point the cross section may grow so large as to violate the partial wave unitarity bound of $\sigma_U = 20\pi/s$, which is appropriate[15] to the case of initial and final fermion states with helicity of 1. However, even for values of $k/\overline{M}_{Pl}$ as large as unity we find that unitarity will not be violated until $\sqrt{s}$ is at least several TeV.

In the circumstance that gravitons are too massive to be directly produced at colliders, their contributions to fermion pair production may still be felt via virtual exchange. For smaller values of $k/\overline{M}_{Pl}$, this would be similar to observing the effects of the SM Z boson before the resonance turns on, or for larger values, to searching for contact interactions. The 4-fermion matrix element is easily computed from the Lagrangian (10) and is seen to reproduce that derived for the scenario of ADD with large extra factorizable dimensions[16]

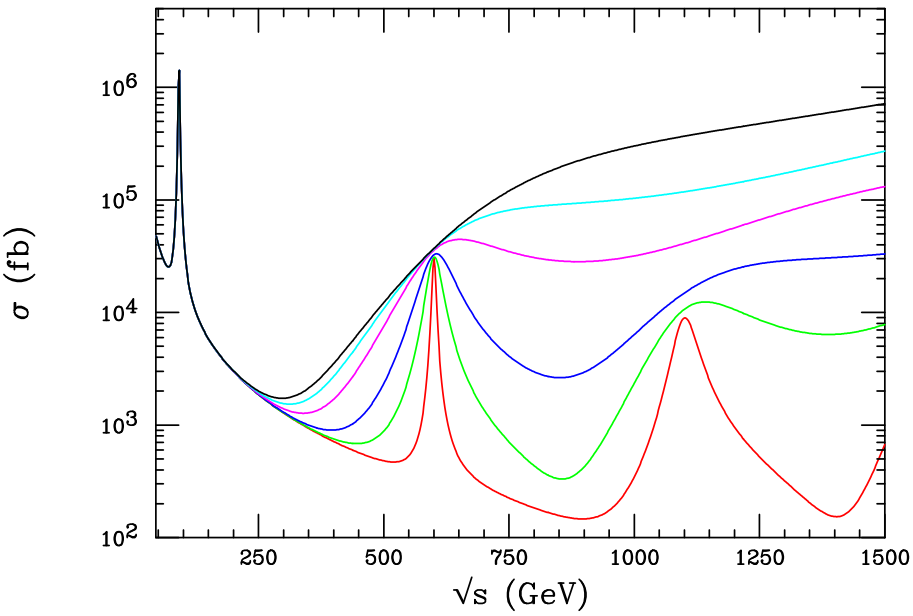

Figure 2: The cross section for $e^+e^- \to \mu^+\mu^-$ including the exchange of a tower of KK gravitons, taking the mass of the first mode to be 600 GeV, as a function of $\sqrt{s}$. From top to bottom the curves correspond to $k/\overline{M}_{Pl} = 1.0,\ 0.7,\ 0.5,\ 0.3,\ 0.2,\ 0.1$.

with the replacement

$$\frac{\lambda}{M_s^4} \to \frac{i^2}{8\Lambda_\pi^2} \sum_{n=1}^{\infty} \frac{1}{s - m_n^2}\,. \tag{13}$$

The advantage in this scenario over the factorizable case is that there are no divergences associated with performing the sum since there is only one new dimension, and hence uncertainties associated with the introduction of a cut-off do not appear. In the limit of $m_n^2 \gg s$, the sum over the KK graviton propagators becomes $[k\Lambda_\pi/\overline{M}_{Pl}]^{-2} \sum_n 1/x_n^2$ which rapidly converges. The 95% C.L. search reach in the $\Lambda_\pi - k/\overline{M}_{Pl}$ plane are given in Fig. 3 for various (a) e^+e^- and (b) hadron colliders. In e^+e^- annihilation we have examined the unpolarized (and polarized for the case of high energy linear colliders) angular and τ polarization distributions, summing over e, μ, τ, c, b (and t, if kinematically accessible) final states, and included initial state radiation, heavy quark tagging efficiencies, an angular cut around the beam pipe, and 90% beam polarization where applicable. For hadron colliders we examined the lepton

pair invariant mass spectrum and forward-backward asymmetry in Drell-Yan production, for both e and μ final states. We also investigated the case where the first two excitations are too close to the collider center-of-mass energy to use the approximation $m_n^2 \gg s$. The bounds in e^+e^- annihilation for this case are given by the solid curves in Fig. 3(a). We see that there is very little difference in the resulting constraints.

As a last point, we note that whereas graviton tower emission was an important probe of the ADD scenario, this is no longer true in the RS model since the graviton states are so massive and can be individually examined on resonance.

In this paper we have explored the phenomenological implications of the Randall-Sundrum localized gravity model of non-factorizable 5-dimensional spacetime, and contrasted it with the ADD scenario. We (*i*) derived the interaction of the KK tower of gravitons with the SM fields, (*ii*) obtained limits on the model parameters using existing data from colliders, both through direct production searches and via virtual exchange contributions, and estimated what future colliders can do to extend these bounds. (*iii*) We described the appearance of KK tower production at high energy linear colliders, the possible loss of the conventional KK signature of successive peaks due to the ever growing widths of these excitations, and (*iv*) demonstrated how measurements of the properties of the first KK state would completely determine the model parameters.

We find the scenario of gravity localization to be theoretically very attractive, and even more importantly, to have distinctive experimental tests. We hope that future experiment will eventually reveal the existence of higher dimensional spacetime.

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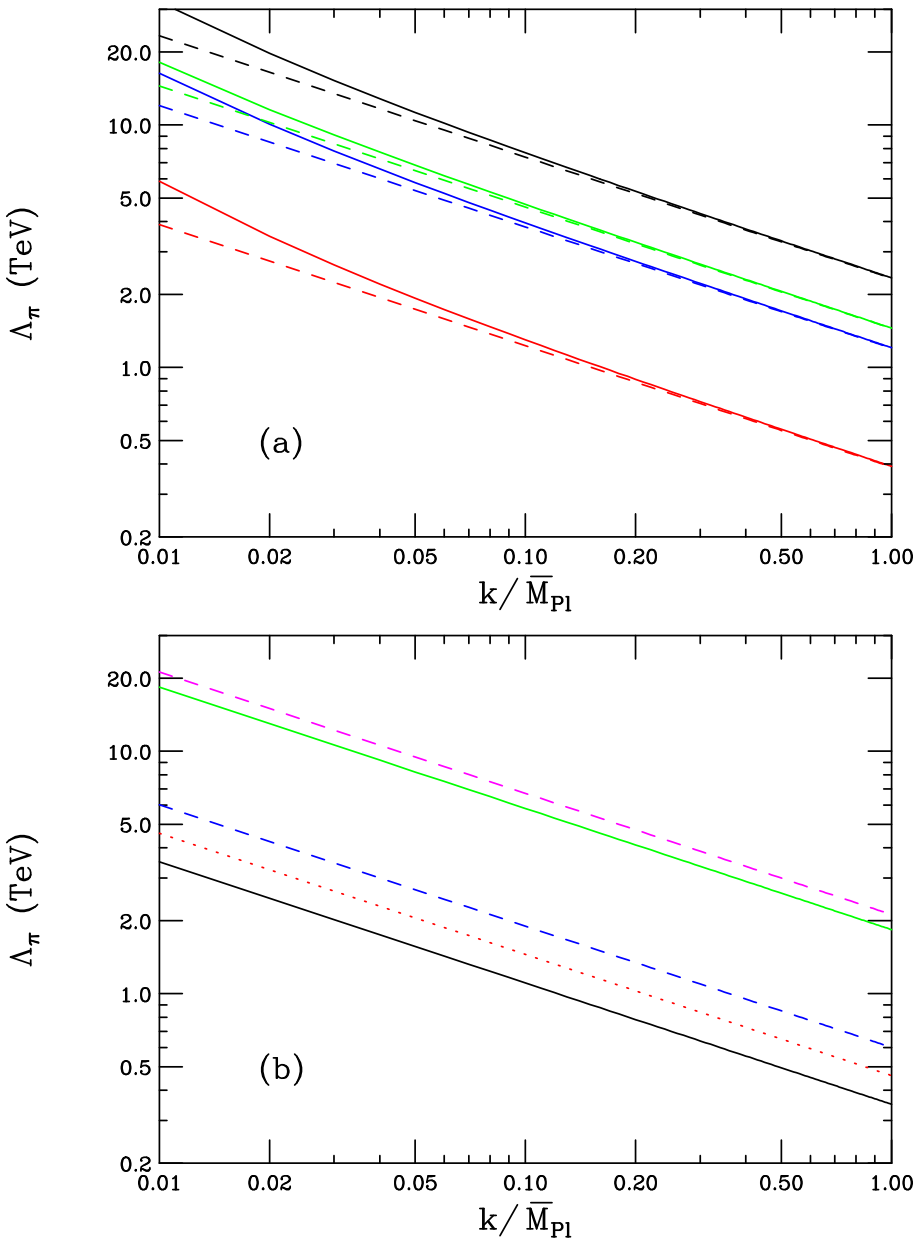

Figure 3: Constraints in the $\Lambda_\pi - k/\overline{M}_{Pl}$ plane from virtual exchange of the tower of KK gravitons. The excluded region lies below the curves. (a) The dashed curves assume the entire tower lies far above $\sqrt{s}$, while the solid curves correspond to the case where the first two excitations are close to $\sqrt{s}$. From bottom to top the pairs of curves correspond to LEP II at 195 GeV with 2.5 fb^{-1} of integrated luminosity; a linear e^+e^- collider at 500 GeV with 75 fb^{-1}; 500 GeV with 500 fb^{-1}; and 1 TeV with 200 fb^{-1}. (b) From bottom to top the curves correspond to the Tevatron Run I with 110 pb^{-1}, Run II with 2 fb^{-1}, Run II with 30 fb^{-1}, the LHC with 10 fb^{-1}, and 100 fb^{-1}.

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