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The little Randall–Sundrum model at the large hadron collider

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ABSTRACT

We present a predictive warped model of flavor that is cut off at an ultraviolet scale $\mathcal{O}(10^3)$ TeV. This “Little Randall–Sundrum (LRS)” model is a volume-truncation, by a factor $y \approx 6$, of the RS scenario and is holographically dual to dynamics with number of colors larger by y . The LRS couplings between Kaluza–Klein states and the Standard Model fields, including the proton constituents, are explicitly calculable without *ad hoc* assumptions. Assuming separate gauge and flavor dynamics, a number of unwanted contributions to precision electroweak, $Zb\bar{b}$ and flavor observables are suppressed in the LRS framework, compared with the corresponding RS case. An important consequence of the LRS truncation, independent of precise details, is a significant enhancement of the clean (golden) di-lepton LHC signals, by $\mathcal{O}(y^3)$, due to a larger “ ρ -photon” mixing and a smaller inter-composite coupling.

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Electroweak symmetry breaking (EWSB) in the Standard Model (SM) via the Higgs condensate $v \equiv \langle H \rangle \simeq 250$ GeV is economical and consistent with data. However, quantum effects render v quadratically sensitive to an ultraviolet (UV) cutoff scale Λ_{UV} . For Λ_{UV} near the gravity scale $\bar{M}_P \sim 10^{18}$ GeV, a severe “hierarchy” $\mathcal{O}(10^{-32})$ arises. One may question the urgency of this problem, as physics close to \bar{M}_P is unknown and inaccessible in the near future. Nonetheless, precision electroweak (EW) data require $\Lambda_{UV} \gtrsim 10$ TeV, near well-tested scales, posing a challenge to a natural Higgs sector. This is often called the little hierarchy; for some recent proposals to address this problem see [1]. Precision flavor data demand $\Lambda_{UV} \gtrsim 10^2\text{--}10^3$ TeV, posing a much more severe “weak-flavor” hierarchy.

The Randall–Sundrum (RS) model [2] was originally proposed to solve the hierarchy problem and yielded distinct collider signatures [3]. However, with 4D-sequestered fermions [4,5], tension with precision data generates a little hierarchy [4,6], the resolution of which led to the inclusion of SM fermions [7] and gauge fields [4,5] in the 5D bulk. This also provided an attractive explanation of the SM flavor structure [7,8], but made the RS model less accessible to experiments [9–14]. In addition, the generic theory requires more structure to be consistent with oblique and non-oblique precision tests [15,16] and constraints from flavor changing neutral currents [17]. In what follows, “RS” denotes the original hierarchy model and all of its extensions.

While the RS construction has a compelling appeal, as it allows a simultaneous resolution of the SM hierarchy and flavor puzzles, it

is premised on a strong assumption. That is, warping extends over many orders of magnitude, without any basic change in physics, from the weak scale to the Planck scale. Surely this assumption needs to be put to an experimental test and we will discuss below how this may indeed be possible, in a warped scenario with various attractive features.

In this work, we use a volume-truncated RS background only to address the hierarchy between the weak (IR) and flavor (UV) scales. SM couplings to new physics, and hence the LHC phenomenology, are explicitly set by the flavor structure without *ad hoc* assumptions. The 5D UV scale M_5 is taken to be $\mathcal{O}(10^3)$ TeV to suppress light-flavor operators in this “Little Randall–Sundrum (LRS)” model [19]. We note that all EW and flavor data are compatible with having $\Lambda_{UV} \sim M_5$, where additional physics may arise.

Keeping Yukawa dynamics unchanged by our truncation, a number of unwanted contributions to precision EW and flavor data are suppressed within the LRS scenario, compared to the RS counterpart. In fact, we will show that any specific RS model is always more constrained than its corresponding LRS counterpart. An exciting consequence of the LRS truncation is a much improved prospect for discovery at the LHC, via clean di-lepton “golden” modes, since the couplings of the KK gauge bosons to light fermions are enhanced while their couplings to the heavy fields are suppressed. Here and below we assume, for simplicity, an IR-brane Higgs and tree level matching for the gauge couplings, as discussed later. Also, we will focus on the quark sector, however, leptons can be included straightforwardly.

The RS background is a slice of AdS_5 , bounded by two Minkowski 3-branes, with the metric [2] $ds^2 = e^{-2\sigma} \eta_{\mu\nu} dx^\mu dx^\nu - r_c^2 d\phi^2$, where $\sigma = kr_c|\phi|$, k is the 5D curvature scale, r_c is the radius of compactification, and $\phi \in [0, \pi]$. The UV (Planck) brane is at $\phi = 0$

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and the IR (TeV) brane is at $\phi = \pi$. Going from the UV brane to the IR brane, the 4D scale redshifts from $k \lesssim M_5$ to the weak scale $\kappa \equiv ke^{-kr_c\pi} \sim \text{TeV}$. Solution to the hierarchy problem requires $kr_c \simeq 11$, with the Higgs at or near the IR brane. A natural flavor structure is obtained, using bulk fermions with non-zero vector-like masses m_i , $i = u, d, \dots$ [7,8]. The resulting zero-mode fermions are exponentially localized in 5D, parameterized by $c_i \equiv m_i/k$. One may choose $c_i \sim 1$ so that light fermions are UV-localized and have small overlaps with the IR-localized Higgs. Due to the warping, the light-flavor cutoff scale is then much larger than the IR/weak scale. This suppresses dangerous light-flavor operators and yields the correct fermion mass hierarchy with $\mathcal{O}(1)$ parameters. However, not all precision data are accommodated with bulk fermions.

Oblique corrections: Here we would like to go over the important constraints on warped phenomenology from precision electroweak data. There are various contributions that can be parameterized in terms of the oblique Peskin–Takeuchi (S, T) parameters [20] and we will discuss them in turn.

We begin by considering the case where the only gauged symmetries in the bulk are those of the 5D SM. First of all, there is a contribution that comes from the *tree-level* mixing of the gauge zero modes with the higher KK modes. In the RS model, these are given by [15],

$$S_{\text{tree}} \approx 2\pi (v/\kappa)^2 \left[1 - \frac{1}{kr_c\pi} + \xi(c) \right], \quad (1)$$

$$T_{\text{tree}} \approx \frac{\pi}{2 \cos^2 \theta_W} (v/\kappa)^2 \left[kr_c\pi - \frac{1}{kr_c\pi} + \xi(c) \right], \quad (2)$$

where

$$\xi(c) \equiv \frac{(2c-1)/(3-2c)}{1 - e^{kr_c\pi(2c-1)}} \left(2kr_c\pi - \frac{5-2c}{3-2c} \right) \quad (3)$$

encodes fermion localization; $\cos^2 \theta_W \simeq 0.77$. For all realistic warped fermion profiles of interest in this work, $\xi(c) \ll 1$.

Without a “bulk” custodial symmetry, there is also a UV-sensitive loop contribution δT to the T parameter. This UV-sensitivity can be absorbed into a higher dimension operator. Assuming that this operator is generated by strong dynamics at the 4D cutoff scale Λ , it will have the form

$$16\pi^2 \frac{(D^\mu H)^\dagger H (H^\dagger D_\mu H)}{\Lambda^2}. \quad (4)$$

Current data favor $|S| \sim |T| \sim 0.1\text{--}0.3$ [21] summed over all contributions. We see that T_{tree} from Eq. (2) is the dominant tree-level constraint, given its volume enhancement $kr_c\pi$, which is roughly a factor of 35 in the RS model. To reduce the size of T_{tree} in this setup then requires increasing κ (m_{KK}) to values that lead to a severe little hierarchy and null LHC signals. Alternatively, we see that reducing $kr_c\pi$ yields a significant suppression, keeping KK masses fixed. In our LRS construct, we will truncate the volume to $kr_c\pi = 6$. Note that $m_{\text{KK}} = \kappa_{\text{KK}} \kappa$, where $\kappa_{\text{KK}}^{\text{RS}} = 2.45, 5.56, \dots$ [4,5] and $\kappa_{\text{KK}}^{\text{LRS}} = 2.70, 5.87, \dots$. Then, for $m_{\text{KK}} \approx 5$ TeV, the RS model yields $(S, T)_{\text{tree}} \approx (0.1, 1.1)$, whereas for the LRS model $(S, T)_{\text{tree}} \approx (0.1, 0.2)$, from Eqs. (1) and (2).

Even though the LRS truncation has suppressed the tree-level KK-tower mixing contribution, we must still address the loop and cutoff scale effects encoded in Eq. (4). These are the same in both the RS and the LRS scenarios, since the flavor sector is assumed to be independent of the gauge dynamics and hence unchanged by the LRS truncation. Barring unnatural cancellations, demanding that the cutoff contribution is less than $\mathcal{O}(0.1)$ pushes m_{KK} to values of $\mathcal{O}(10)$ TeV.

Ref. [15] attributed $T \gtrsim 1$, in the RS model, to the absence of bulk “custodial” protection and postulated a $SU(2)_L \times SU(2)_R \times U(1)_X$ 5D symmetry to eliminate tree-level contributions to T . It

turns out that the loop contribution δT , governed by fermion KK modes, still remains, but is no longer UV-sensitive. Also, given the gauged 5D custodial symmetry, there is no cutoff contribution of the form in Eq. (4), at the IR-boundary. Ref. [15] concluded that including the SM effects, for $m_{\text{KK}}^{\text{RS}} \sim 3\text{--}4$ TeV and a light Higgs, S and T can be accommodated at an acceptable level. Note that with a bulk custodial symmetry, the LRS construct will enjoy the same level of agreement with the oblique data as the models in Ref. [15].

The suppression of T_{tree} in the LRS scenario, without bulk custodial symmetry, can be understood as follows. After EWSB, the entire KK tower of states mixes and the zero mode (SM) wavefunctions get deformed away from a constant. This generates a large tree-level oblique correction to T , controlled by the KK-Higgs (i.e., KK-IR brane) coupling. The deformation of the SM W/Z wavefunction is then proportional to $kr_c\pi$. Hence, the LRS contributions get suppressed by a factor

$$y \equiv \frac{kr_c\pi|_{\text{RS}}}{kr_c\pi|_{\text{LRS}}} \approx 6. \quad (5)$$

Note that S from Eq. (1) is basically the same in the RS and LRS models. This is because the dominant contributions to S come from a universal shift in the light fermion-gauge field couplings [15]. This shift depends on the product of the mixing between the zero modes and KK gauge states (after EWSB) and the universal couplings of KK gauge fields to the light fermions. While the former decreases with shrinking volume the latter is increased so that the product is unchanged.

In brief, the LRS truncation does suppress the contribution from KK-tower mixing in the gauge sector, compared to the RS case, quite efficiently. However, a bulk custodial symmetry is still required in both the RS and LRS setups to control loop and cutoff scale contributions to T and to bring the scale of KK masses down to ~ 3 TeV [15,22]. As we will show next, a more dramatic improvement can be achieved regarding the non-oblique and precision flavor observables.

Non-oblique correction and flavor physics: Refs. [23,25] have shown that RS models flow to next-to-minimal-flavor violation, where flavor changing effects are primarily from mixing with the third generation. The extra flavor breaking sources are quasi-aligned with their SM counterparts and the misalignment is at most of order the CKM matrix, but new sources of CP violation are present. Non-oblique $Zb\bar{b}$ and FCNC constraints are more involved since they depend on the amount of flavor non-universality, determined by the fermion zero-mode IR-brane profile values, $f_{Q,u,d}$ [23].

As shown in Refs. [23,24], once the overall scale of 5D Yukawa coupling λ_5 and t_R -localization c -parameter are fixed, the localization of all other fermions is set by the measured masses and CKM mixing angles, assuming anarchical 5D Yukawa matrices via the relation $m_{u,d} \propto F_Q \lambda_5^{u,d} F_{u,d}$, where m_x denotes 4D masses and $x = u, d$ correspond to the up and down quarks respectively; f_x are eigenvalues of F_x . Here, we will keep λ_5 unscaled by LRS-truncation and fixed at its “RS value” [17,23] (see discussion below). Then, the amount of non-universality is unchanged, but the strength of the KK-mediated effects get decreased like the truncated LRS volume, parameterized by $kr_c\pi$. This is the reason the non-universal precision observables are significantly suppressed in our model. For concreteness, in Table 1, we give a set of c_x and f_x values that reproduce the SM quark masses and mixing angles; $\lambda_5/k = 2$ in accordance with Ref. [23].

Bulk-RS constraints from the $Zb\bar{b}$ coupling require non standard fermion representations under the custodial symmetry, as well as a Z_2 symmetry [16], in order to have $m_{\text{KK}}^{\text{RS}} \sim 3$ TeV; otherwise, $m_{\text{KK}}^{\text{RS}} \gtrsim 5$ TeV [15]. Without the custodial symmetry, there are various contributions to $Zb\bar{b}$. The first originate from the gauge

Table 1

The eigenvalues of c_x and f_x which roughly yield the right masses and CKM elements at the m_Z scale [18]

Flavor	c_Q, f_Q	c_u, f_u	c_d, f_d
I	1.45, 0.003	$1.7, 8 \times 10^{-4}$	1.52, 0.002
II	1.17, 0.015	0.86, 0.071	1.26, 0.009
III	0.52, 0.28	-0.19, 0.83	1.14, 0.018

zero-mode-KK-tower mixing due to EWSB and the enhanced coupling of the gauge KK modes to IR-localized b_L . These corrections are proportional to $(kr_c\pi/m_{KK}^2)f_{Q3}^2$ [15], where f_{Q3} is assumed to have the RS value. Thus, the LRS bound from these contributions is

$$m_{KK}^{\text{LRS}} \gtrsim m_{KK}^{\text{RS}}/\sqrt{y} \quad (Zb\bar{b}). \quad (6)$$

With $\sqrt{y} \approx 2.4$ we get $m_{KK}^{\text{LRS}} \gtrsim 2$ TeV.

The second type of correction to $Zb\bar{b}$ is due to $\mathcal{O}(1)$ mixing between b_L and the exotic $SU(2)_R$ partner of t_R .¹ This latter contribution will be absent for a choice of representation in which t_R is a L - R isosinglet [16]. Note that without a bulk custodial symmetry, there is no exotic t_R partner. However, there is a third type of correction to $Zb\bar{b}$ from the mixing of the KK modes of b_R and the b_L zero mode. This contribution is not truncated in the LRS model and is of order $4[(v/\sqrt{2})\lambda_5 k f_{Q3}/m_{KK}(b_R)]^2$. To keep deviations in the $Zb\bar{b}$ coupling below 0.3%, we then need $m_{KK}(b_R) \gtrsim 4$ TeV. Interestingly, $m_{KK} \gtrsim 3$ TeV for gauge fields already implies the former bound for the KK modes of b_R , in the LRS framework presented here. We hence conclude that all of the above constraints from $Zb\bar{b}$ can be satisfied for gauge sector $m_{KK} \gtrsim 3$ TeV, without any protective symmetries (however, as discussed above, this would be inconsistent with the bound from the T parameter).

We finally review the strongest constraints on generic bulk RS models, from $\Delta F = 2$ processes due to tree level exchange of KK gluons. Ref. [23] showed that, with an IR localized Higgs, the ratio of RS and SM $(V-A) \times (V-A)$ contributions $h^{\text{RS}} \propto (F_Q^2)_{ij}^2$ (in the down quark mass basis). One can write

$$h^{\text{RS}} = \frac{M_{12}^{\text{RS}}}{M_{12}^{\text{SM}}} \sim 0.5 \times \frac{kr_c\pi}{35} \left(\frac{3 \text{ TeV}}{m_{KK}} \right)^2 \left(\frac{f_{Q3}}{0.3} \right)^4. \quad (7)$$

At present, $h^{\text{RS}} \lesssim 0.3$ [25–27]. However, the dominant contribution $\delta(\epsilon_K)$ to ϵ_K from $(V-A) \times (V+A)$ operators [26] is given by $\delta(\epsilon_K) \propto kr_c\pi (F_Q^2)_{12} (F_d^2)_{12}$ [17] which is $\mathcal{O}(20)$ times smaller. This is not enough due to a matrix element chiral enhancement of $\mathcal{O}(11)$ and a $\mathcal{O}(7)$ factor from the running between the KK and weak scales, requiring $m_{KK}^{\text{RS}} \gtrsim 8$ TeV. In the LRS case, both contributions are suppressed by y and thus $m_{KK}^{\text{LRS}} \gtrsim 3$ TeV, no stronger than the oblique constraints. We note that the RS CP electric dipole moment problem [23] persists in our setup, as it is governed by 5D Yukawa interactions which are unchanged (for possible solutions to this problem see [17,28]).

Phenomenology: Gauge KK modes couple to UV-localized light fermions (important initial states at colliders), with strength $g_{KK} \sim g_4/\sqrt{kr_c\pi}$; g_4 is a typical 4D SM gauge coupling. We get $g_{KK}^{\text{RS}} \sim g_4/6$ and $g_{KK}^{\text{LRS}} \sim g_4/2.5$. In particular, the UV-brane values of the normalized first gauge KK wavefunctions $\chi^{(1)}$ are

$$\chi_{\text{RS}}^{(1)}(\phi=0) = -0.08; \quad \chi_{\text{LRS}}^{(1)}(\phi=0) = -0.20. \quad (8)$$

This leads to improvements in the LHC sensitivity, at fixed m_{KK} , for the following reasons: (i) Typically broad states [10] become narrower by a factor $y \sim (0.2/0.08)^2$, (ii) branching ratio (BR) into light states (e.g., e^+e^-) increases by a factor y^2 , (iii) from (i) and (ii) one can show that the signal S goes up by $y^3 \sim 250$, while

Table 2

Summarized comparison of constraints and predictions in the RS and the LRS scenarios. For simplicity and definiteness, the Higgs is assumed to be on the IR-brane. The constraints correspond to lower bounds on gauge KK masses, in TeV. Here, we assume a custodial symmetry for the T parameter; a left–right Z_2 symmetry is imposed to protect the $Zb\bar{b}$ coupling, unless denoted by *. The predictions in the last row correspond to a Z' of mass {2, 5} TeV, respectively

Constraint/prediction	RS	LRS
T parameter	3	3
S parameter	3	3
$Z \rightarrow b\bar{b}$	3	3*
ϵ_K	8	3
S/B for $Z' \rightarrow t^+t^-$	{0.3, -}	{ $\mathcal{O}(100)$, $\mathcal{O}(100)$ }

the background B drops as $1/y$, over the resonance width. Hence, S/B in the LRS model is expected to go up by a factor $y^4 \sim 1500$, a remarkable enhancement! These features lead to a larger LRS discovery reach and a way to test the validity of this setup and the underlying assumptions. As the LHC reach for KK gluons in bulk-flavor RS models is 3–4 TeV [10], the corresponding LRS reach can be as big as ~ 5 TeV. The enhanced g_{KK} in the LRS model could also allow access to the elusive EW gauge KK modes [14]. For example, the $Z' \rightarrow \ell^+\ell^-$, $\ell = e, \mu$, golden decay modes which were close to hopeless within the RS case [14] could lead to discovery in the LRS setup. Using the same cuts as in Ref. [14], we find roughly 2000 (3) events with S/B and $S/\sqrt{B} \gg 1$ [in fact $\mathcal{O}(100)$] for $M_{Z'} = 2$ (5) TeV and 100 fb^{-1} .² Note that given the significance of the signal, a discovery would be unambiguous over this range. Enhancement in production rate is expected for the SM KK fermions whose LHC discovery, in the RS model, would be quite challenging [29]. Furthermore, any bulk couplings mediated via λ_5 are relatively stronger due to LRS scaling. This would, in principle, be a direct test of our setup. Also, BR's of the neutral modes into composite states such as $W_L W_L, Z_L h$ and $t\bar{t}$ compared to those into light fermions will provide a robust test of the LRS construct.

TeV-scale spin-2 “graviton” resonances are distinct RS signatures [3,9]. Since $M_5 \sim 10^3$ TeV, we eliminate the zero mode graviton, using UV-brane Dirichlet boundary conditions; we find $\chi_G^{(\text{LRS})} \simeq \chi_G^{(\text{RS})} = 3.83, 7.02, \dots$. Gluon-fusion production of KK gravitons is dominant at hadron colliders, with a cross section proportional to $[(k/M_5)/kr_c\pi]^2 (k/M_5)(x_G/m_G)^2$ [9]. In a generic LRS model, both k/M_5 and $kr_c\pi$ shrink by a factor of y , decreasing this cross section by a factor of $\mathcal{O}(y)$. Thus, observation of the LRS spin-2 resonances is unlikely at the LHC (probably more unlikely than in the RS case [9]).

We also note that there can be LRS collider signatures in terms of deviations from the SM top-quark couplings. Within the RS scenario, there are two types of contributions to $t \rightarrow cZ$ of similar size [30]. One is through mixing between Z zero and KK modes which will be suppressed by y and probably unobservable in the LRS framework. However, the second one proceeds through mixing between t_R and the KK modes of t_L . This mixing is controlled by the 5D Yukawa which is left unchanged in our LRS construct and, therefore, $t \rightarrow cZ$ should be within the reach of the LHC. A similar effect yields an $\mathcal{O}(20\%)$ shift in the $Zt_R\bar{t}_R$ coupling which is probably beyond the LHC sensitivity but may be observed at a future linear collider [31].

Holography: We now present a qualitative discussion of the LRS, using the AdS/CFT correspondence [32], following previous interpretations of geometric RS results [33,34] in the dual context of a strongly coupled large N 4D gauge theory [35]. We begin by studying the effects of LRS-truncation on the classical geometric relation

¹ We thank K. Agashe for emphasizing this point.

² Similar LHC reaches were found, in a different context, by M. Piai in Ref. [19].

between g_4 and the 5D gauge coupling g_5 [36]:

$$1/g_4^2 = \tau_{UV} + \tau_{IR} + \log(k/\kappa)/(kg_5^2). \quad (9)$$

Here, τ_{UV} and τ_{IR} will be treated as small UV and IR quantum threshold corrections, respectively. For a generic comparison of couplings, we neglect $\tau_{UV,IR}$ and keep g_4 fixed to its measured value. Thus, reducing the volume suppression $kr_c\pi$ (the log) requires lowering the value of kg_5^2 . In the dual CFT, this is interpreted as the contribution of CFT “quarks” to the running of external gauge couplings from the fundamental scale, M_5 , down to the TeV scale (just like the contribution of SM quarks to α_{QED} running) [33,34,37]. Thus the relation $\sqrt{kg_5^2} \sim 4\pi/\sqrt{N}$ should hold between the dual theories. Explicit calculations [4,5] confirm that couplings among gauge KK modes, i.e., IR localized fields, are enhanced, compared to the corresponding zero mode gauge coupling, by $\sqrt{kr_c\pi} \sim \sqrt{kg_5^2}$. This, in the dual CFT picture, corresponds to the coupling of three composites given by $4\pi/\sqrt{N}$ at large N [35,38]. Consequently, the truncated LRS volume is dual to $N^{\text{LRS}} \sim yN^{\text{RS}} > N^{\text{RS}}$, making the inter-composite interactions weaker. The weakened CFT interactions with the Higgs, a composite state, account for the decrease in T from Eq. (2).

In our LRS construct, we held the 5D Yukawa coupling λ_5 unscaled, lowering the 4D IR-brane cutoff to about 10 TeV, as in the RS case. In the dual language, this corresponds to separate dynamics for this sector, characterized by a “flavor” CFT with $N_F \sim N^{\text{RS}} < N^{\text{LRS}}$ ($N_F \sim 3-4$ [17,23]). This independent CFT is linked to dynamical breaking of the 5D flavor group which is not completely broken by the bulk masses (unlike $\langle H \rangle$ which breaks the “chiral” symmetry) [23]. If the Higgs is dynamically realized as a pseudo-Goldstone boson (PGB) a similar different scaling for its potential should be applied, i.e., the dynamics which generates the PGB potential is characterized by $N = N_F$. Otherwise, increasing N would induce a more severe fine tuning for the PGB potential [39]. However, in all the known models including the most realistic ones [39] the dominant contributions to this potential come from the top sector, corresponding to the flavor CFT, and not from the weakly gauged one, consistent with the above assumptions [40].³ This also explains why keeping $S \sim (v/f_\pi)^2 N^{\text{LRS}}$ at the RS value does not lead to extra fine-tuning of v , since the “decay constant” f_π here is from the weakly gauged dynamics which does not govern the Higgs potential. The constancy of the S parameter under truncation can be understood as follows. The main contribution to S is from the universal vertex corrections [15]. This is controlled by gauge zero-KK mode mixing, which scales as $1/\sqrt{N}$, and the universal KK couplings to light fermions, which scales as \sqrt{N} (see below). Therefore, S remains unchanged.

The non-oblique and FCNC contributions depend on the amount of non-universality in the couplings of the KK states to different generations. On the CFT side, this corresponds to the amount of partial compositeness for a given N_F . The amount of compositeness follows from the observed masses and mixing angles [23], once λ_5 is set and the location of t_R is decided. By fixing these to the RS value, the LRS amount of partial compositeness is then unchanged, and hence the non-universal observables are suppressed by truncation. This, generically, yields a better agreement with the data. It implies that interactions proportional to λ_5 (such as between the Higgs and two KK fermions) are stronger than the corresponding KK gauge interactions.

An important consequence of volume truncation is enhanced ρ -photon mixing, proportional to \sqrt{N} [33,37], leading to larger couplings of light SM fermions to gauge composite/KK modes.

³ However, for the particular minimal realization of a PGB Higgs as an A_5 , flavor and gauge dynamics are of the same origin which implies only a single value of N .

The composite (KK) partial widths into elementary fermions scales as N , while the total width drops as $1/N$. Hence, $S \sim N^3$ and $\mathcal{B} \sim 1/N$, over the resonance width. Both effects yield stronger LRS signals at the LHC than for the RS case, since $y = N^{\text{LRS}}/N^{\text{RS}} \gg 1$, as discussed before. This is analogous to how $e^+e^- \rightarrow \rho \rightarrow \mu^+\mu^-$ is modified when N_c is increased.

Finally, we emphasize that unless mentioned explicitly (as for λ_5) we have rescaled all the couplings in the theory according to the LRS value of N . This is why we did not get an enhancement in the KK graviton production. Also, we have neglected brane-kinetic terms to allow a transparent comparison of our model with generic RS models where such terms have sub-dominant effects on the above observables. Lastly, we note that $M_5 \sim 10^3$ TeV does not suppress baryon and lepton number violation sufficiently. Such issues lie beyond the scope of this work, but can be addressed with discrete symmetries or in a UV-completion of the LRS model. However, dimension-9 operators suppressed by the LRS M_5 lead to $n-\bar{n}$ oscillations at acceptable levels [41] and may be accessible in near future experiments.

In summary, we presented the “Little Randall–Sundrum (LRS)” model of hierarchy between the flavor and weak scales which is much less constrained than \tilde{M}_P -weak warped scenarios. Here, the 5D cutoff scale $M_5 \sim 10^3$ TeV is chosen to suppress unwanted light-quark operators sufficiently and the weak scale is obtained from $\mathcal{O}(M_5)$ scales by warping; the flavor puzzle is addressed by fermion localization, as in the RS model. Even without a bulk custodial symmetry, the “tree-level” lower bound on LRS gauge KK masses is at ~ 5 TeV; the RS bound is $m_{\text{KK}}^{\text{RS}} \gtrsim 12$ TeV. Loop and higher dimension contributions to T raise $m_{\text{KK}} \gtrsim \mathcal{O}(10)$ TeV, without a custodial symmetry, in both RS and LRS models. Custodial symmetry can be imposed if desired, leading to an oblique lower bound at ~ 3 TeV. As we kept the overall Yukawa scale unchanged, the most severe RS-type constraints are much better behaved here: non-oblique constraints from $Zb\bar{b}$, without a protective Z_2 symmetry, are absent for $m_{\text{KK}} \gtrsim 3$ TeV and the FCNC bounds are largely relaxed. We have summarized the comparison between the RS and LRS frameworks regarding various precision constraints in Table 2. As can be seen from this table, the LRS framework is never more constrained than its RS counterpart and in many cases it is much more compatible with data. Typical LRS graviton KK modes are more elusive than those of \tilde{M}_P -weak warped models. However, the light-fermion LHC-production rate and BR’s for LRS gauge KK modes are much bigger than the corresponding RS values and yield a signal $\sim y^3$ times larger; $y \sim 6$ is the LRS truncation factor. We hence conclude that the LRS model is a good candidate for new physics and may soon be uncovered at the LHC, or perhaps probed at a future linear collider.

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