Probing warped extra dimension via $gg \to h$ and $h \to \gamma\gamma$ at LHC

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A B S T R A C T

The processes $gg \to h$ and $h \to \gamma\gamma$ are of paramount importance in the context of Higgs search at the LHC. These processes are loop driven and hence could be sensitive to the presence of any new colored fermion states having a large coupling with the Higgs. Such a scenario arises in a warped extra-dimensional theory, where the Higgs is confined to the TeV brane and the hierarchy of fermion masses is addressed by localizing them at different positions in the bulk. We show that the Yukawa coupling of the Higgs with the fermion Kaluza–Klein (KK) states can be order one irrespective of their zero mode masses. We observe that the $gg \to h$ and $h \to \gamma\gamma$ rates are substantially altered if the KK states lie within the reach of LHC. We provide both intuitive and numerical comparison between the RS and UED scenarios as regards their quantitative impact in such processes.

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1. Introduction

For an intermediate mass (< 150 GeV) Higgs boson, the relevance of its production at the CERN Large Hadron Collider (LHC) via gluon fusion ($gg \to h$) and its subsequent decay into two photons ($h \to \gamma\gamma$) cannot be over-emphasized. Since these are loop induced processes, a natural question arises as how sensitive these processes are to the existence of new physics. In this Letter, we explore such a possibility by embedding the Standard Model (SM) in a Randall–Sundrum (RS) warped geometry [1], where the bulk is a slice of Anti-de Sitter space (AdS$_5$) accessible to some or all SM particles [2,3]. The virtues of such a scenario include a resolution of the gauge hierarchy problem caused by the warp factor [1], and an explanation of the hierarchy of fermion masses by their respective localizations in the bulk keeping the Higgs confined at the TeV brane [4]. Besides, the smallness of the neutrino masses could be explained [5], and light KK states would lead to interesting signals at LHC [6]. We demonstrate in this Letter that the loop contribution of the KK towers of quarks and gauge bosons emerging from the compactification would have a sizable numerical impact on the $gg \to h$ and $h \to \gamma\gamma$ rates. This happens because the Higgs coupling to a pair of KK fermion–antifermion is not suppressed by the zero mode fermion mass and can easily be order one [7]. The underlying reason is simple. Although the zero mode wave-functions of different flavors have varying overlap at the TeV brane depending on the zero mode masses, the KK profiles of all fermions have

a significant presence at the TeV brane where the Higgs resides. As a result, the KK Yukawa couplings of different flavors are not only all large, they are also roughly universal. This large universal Yukawa coupling in the RS scenario constitutes the corner-stone of our study. On the contrary, in flat Universal Extra Dimension (UED) only the KK top Yukawa coupling is large, others being suppressed by the respective zero mode fermion masses. We provide comparative plots to demonstrate how the warping in RS fares against the flatness of UED for the processes under consideration.

2. Warped extra dimension

The extra coordinate y is compactified on an $S^1/\mathbb{Z}_2$ orbifold of radius R, with $-\pi R \leqslant y \leqslant \pi R$. Two 3-branes reside at the orbifold fixed points at $y = (0, \pi R)$. The space–time between the two branes is a slice of AdS$_5$ geometry, and the 5d metric is given by [1],

$$ds^2 = e^{-2\sigma} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2, \quad \text{where } \sigma = k|y|. \tag{1}$$

Above, $1/k$ is the AdS$_5$ curvature radius, and $\eta_{\mu\nu} = \text{diag}(-1, 1, 1, 1)$. The natural mass scale associated with the $y = 0$ brane is the Planck scale (M_P), while the effective mass scale associated with the $y = \pi R$ brane is $M_P e^{-\pi kR}$, which is of the order of a TeV for $kR \simeq 12$. To address the fermion mass hierarchy, the Higgs boson has to be confined to the TeV brane, thus solving the gauge hierarchy problem in the same stroke. The bulk contains the fermions and gauge bosons. After integrating out the y-dependence, the 4d Lagrangian can be written in terms of the zero modes and their KK towers. A generic 5d field can be decomposed as (only fermions and gauge bosons are relevant for us) [3]

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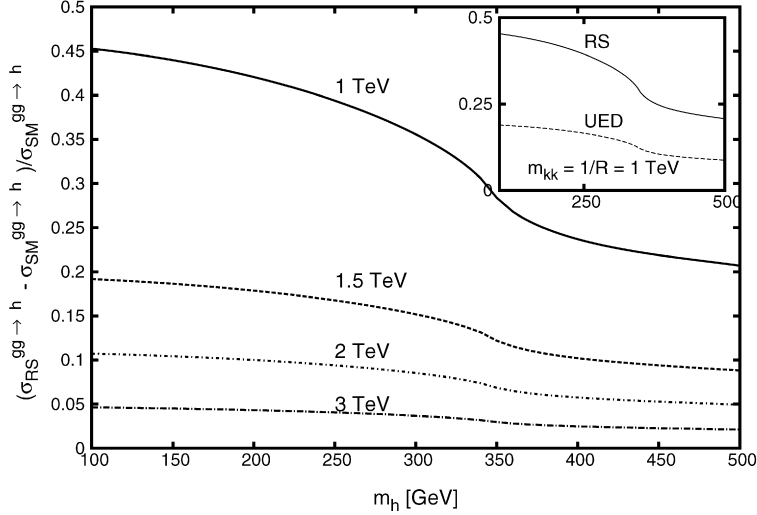
Fig. 1. The fractional deviation (from the SM) of the $gg \to h$ production cross section in RS is plotted against the Higgs mass. The four curves correspond to four different choices of m_{KK}. In the inset, we have compared the UED contribution for $1/R = 1$ TeV with the RS contribution for $m_{KK} = 1$ TeV.

$$\Phi(x^\mu, y) = \frac{1}{\sqrt{2\pi R}} \sum_{n=0}^{\infty} \Phi^{(n)}(x^\mu) f_n(y),$$

$$\text{where } f_n(y) = \frac{e^{s\sigma/2}}{N_n} \left[J_\alpha\left(\frac{m_n}{k} e^\sigma\right) + b_\alpha(m_n) Y_\alpha\left(\frac{m_n}{k} e^\sigma\right) \right], \quad (2)$$

with $s = (1, 2)$ for $\Phi = \{e^{-2\sigma} \Psi_{L,R}, A_\mu\}$. Above,

$$b_\alpha = -\frac{(-r + \frac{s}{2}) J_\alpha(\frac{m_n}{k}) + \frac{m_n}{k} J'_\alpha(\frac{m_n}{k})}{(-r + \frac{s}{2}) Y_\alpha(\frac{m_n}{k}) + \frac{m_n}{k} Y'_\alpha(\frac{m_n}{k})},$$

$$\text{and} \quad N_n \simeq \frac{1}{\sqrt{\pi^2 R m_n e^{-\pi kR}}}, \quad (3)$$

where $r = (\mp c, 0)$ and $\alpha = (c \pm \frac{1}{2}, 1)$ for $\mathbb{Z}_2$ even/odd modes. By imposing boundary conditions on $f_n(y)$ in Eq. (2) and in the limit $m_n \ll k$ and $kR \gg 1$, one obtains the KK masses as (for $n = 1, 2, \ldots$),

$$m_n \simeq \left(n + \frac{1}{2}\left|c - \frac{1}{2}\right| - \frac{1}{4}\right) \pi k e^{-\pi kR} \quad \text{(fermions)},$$

$$m_n \simeq \left(n - \frac{1}{4}\right) \pi k e^{-\pi kR} \quad \text{(gauge bosons)}. \quad (4)$$

Now, the Yukawa part of the action with two 5d Dirac fermions $\Psi_{iL}(x, y)$ and $\Psi_{iR}(x, y)$ for each flavor i is given by [3]

$$S_y = \int d^4 x \int dy \sqrt{-g} \, \lambda_{ij(5d)} H(x)$$
$$\times \left(\bar{\Psi}_{iL}(x, y) \Psi_{jR}(x, y) + \text{h.c.}\right) \delta(y - \pi R). \quad (5)$$

For simplicity we ignore flavor mixing, and further assume $c_{iL} = c_{iR} = c_i$. The Yukawa coupling of the zero mode fermions turns out to be [3]

$$\lambda_i = \lambda_{i(5d)} k (1/2 - c_i) \left(1 - e^{(2c_i - 1)\pi kR}\right)^{-1}. \quad (6)$$

Assuming the 5d coupling $\lambda_{i(5d)} k \sim 1$, one can trade the zero mode fermion masses in favor of the corresponding c_i ($c_q = 0.62, 0.61, 0.51, 0.56, -0.49, 0.48$ for $q = u, d, c, s, t, b$). This is how the fermion mass hierarchy problem is addressed. Next, we derive the Yukawa coupling of the nth KK fermion for $m_n \ll k \sim M_P$, $kR \gg 1$ and $\lambda_{i(5d)} k \sim 1$ as (with a tacit assumption of KK number conservation to avoid any divergence in KK sum)

$$\lambda_i^{(n)} \sim \cos^2\left(\left[n - \frac{|c - \frac{1}{2}| - |c \mp \frac{1}{2}|}{2} - \frac{1}{2}\right]\pi\right), \quad (7)$$

where $\mp$ correspond to $\mathbb{Z}_2$ odd/even KK modes. Thus the KK Yukawa couplings for $\mathbb{Z}_2$ even KK modes, regardless of their flavors and KK numbers, are roughly equal to unity for the values of c_q quoted above.

3. Contribution of KK states to $\sigma(gg \to h)$

The process $gg \to h$ proceeds through fermion triangle loops. The SM expression of the cross section is given by ($\tau_q \equiv 4m_q^2/m_H^2$)

$$\sigma_{gg \to h}^{SM} = \frac{\alpha_s^2}{576\pi v^2} \left| \sum_q A_q(\tau_q) \right|^2,$$

$$\text{where } A_q(\tau_q)\big|_{SM} = 2\tau_q \big[1 + (1 - \tau_q) f(\tau_q)\big], \quad (8)$$

$$\text{with } f(\tau) = \arcsin^2\left(\frac{1}{\sqrt{\tau}}\right) \quad \text{for } \tau \geqslant 1,$$

$$\text{and} \quad f(\tau) = -\frac{1}{4}\left[\ln\left(\frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}}\right) - i\pi\right]^2 \quad \text{for } \tau < 1. \quad (9)$$

Above, α_s is the QCD coupling at the Higgs mass scale, v is the Higgs vacuum expectation value and A_q is the loop amplitude from the qth quark. In the SM, the dominant contribution comes from the top quark loop. Now, there will be additional contributions from the KK quarks. Importantly, due to the large universal KK Yukawa couplings, not only the KK top but also the KK modes of other quarks would have sizable contribution. Indeed, the lightest modes ($n = 1$) would have dominant contributions. Setting the KK Yukawa couplings to unity, as suggested by Eq. (7), we derive the amplitude of the nth KK mediation of the qth flavor, with the same normalization of Eq. (8), as

$$A_q(\tau_{q_n})\big|_{KK} = \frac{4v^2}{m_h^2} \big[1 + (1 - \tau_{q_n}) f(\tau_{q_n})\big]. \quad (10)$$

In 5d the sum over n yields a finite result. Eq. (10) is different from the UED result [8] in two ways: (i) we have set the KK Yukawa coupling to unity irrespective of quark flavors, while in UED the KK Yukawa coupling is controlled by zero mode masses; (ii) in UED there is an additional factor of 2 because both $\mathbb{Z}_2$ even and odd KK modes contribute, while in RS the odd modes do not couple to the brane-localized Higgs. In Fig. 1, we have plotted the variation with m_h of the deviation of the production cross section $\sigma_{RS}(gg \to h)$ from its SM expectation $\sigma_{SM}(gg \to h)$ normalized by the SM value. The dominant QCD correction cancels
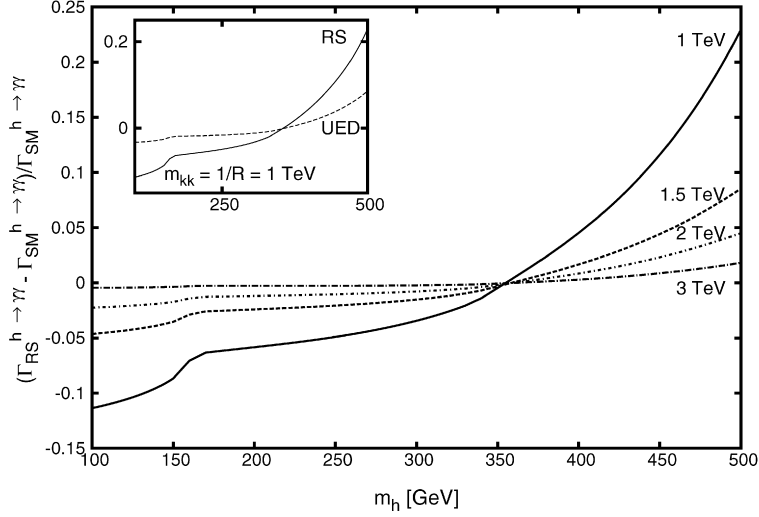
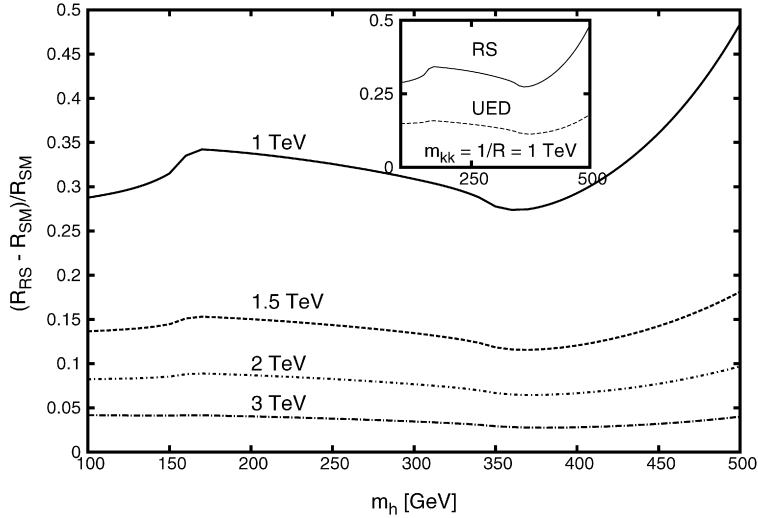**Fig. 2.** Same as in Fig. 1, except that the fractional deviation in $h \to \gamma\gamma$ decay width is plotted.

Fig. 3. Same as in Fig. 1, except that the fractional deviation in $R = \sigma_{gg \to h} \Gamma_{h \to \gamma\gamma}$ has been plotted.

in this normalization. We have chosen four reference values of m_{KK} ($= 1.0, 1.5, 2.0$ and 3.0 TeV), where m_{KK} is the KK mass of the $n = 1$ gauge bosons, which also happens to be the lightest KK mass in the bulk (corresponding to the conformal limit, $c = 1/2$ for fermions). For m_h below 150 GeV, the deviation is quite substantial (close to 45%) for $m_{KK} = 1$ TeV. For larger $m_{KK} = 1.5$ (3.0) TeV, the effect is still recognizable, around 18% (5%). In the inset, we exhibit a comparison between RS and UED contributions to the same observable, where the KK mass scales of the two scenarios, namely m_{KK} for RS and $1/R$ for UED, have been assumed to be identical ($= 1$ TeV). For $m_h < 150$ GeV, the RS contribution is about 2.5 times larger than the UED contribution, while the margin slightly goes down with increasing m_h. This factor 2.5 can be understood in the following way: In RS, five $n = 1$ KK flavors (except the KK top) have mass around m_{KK} with order one Yukawa coupling. So naively we would expect a factor of 5 enhancement relative to UED. But in UED both $\mathbb{Z}_2$ even and odd modes contribute. This reduces the overall enhancement factor in RS over UED to about 2.5.

4. Contribution of KK states to $\Gamma(h \to \gamma\gamma)$

The $h \to \gamma\gamma$ process proceeds through fermion triangles as well as via gauge loops along with the associated ghosts. The de-

cay width in the SM can be written as

$$\Gamma_{h \to \gamma\gamma} = \frac{\alpha m_h^3}{256\pi^3 v^2} \left| \sum_f N_c^f Q_f^2 A_f(\tau_f) + A_W(\tau_W) \right|^2, \tag{11}$$

where α is the electromagnetic coupling at the Higgs mass scale. The expression for A_f is given in Eq. (8), and the dominant SM contribution to A_f comes from the top quark loop. The W-loop amplitude in the SM is given by

$$A_W(\tau_W)\big|_{SM} = -\left[2 + 3\tau_W + 3\tau_W(2 - \tau_W)f(\tau_W)\right]. \tag{12}$$

We derive the KK contribution of the gauge sector as

$$A_W(\tau_{W_n})\big|_{KK} = -\left[2 + 3\tau_W + 3\tau_W(2 - \tau_{W_n})f(\tau_{W_n}) - 2(\tau_{W_n} - \tau_W)f(\tau_{W_n})\right]. \tag{13}$$

Again, the sum over n yields finite result and in the limit of large KK mass the KK contribution decouples. Our Eq. (13) is very different from the corresponding UED expression [8], primarily because the Higgs is confined at the brane in the present scenario while it resides in the bulk in UED. In Fig. 2, we have plotted the decay width $\Gamma(h \to \gamma\gamma)$ in RS relative (and normalized as well) to the SM. Again, the four choices of m_{KK} are 1.0, 1.5, 2.0 and

3.0 TeV. There is a partial cancellation between quark and gauge boson loops, both in real and imaginary parts, not only for the zero mode but also for each KK mode. The meeting of the four curves just above the $m_h = 2m_t$ threshold is a consequence of the above cancellation and at the meeting point the SM contribution overwhelms the KK contribution. Unlike in Fig. 1, we witness both suppression and enhancement with respect to the SM contribution. The inset carries an illustration how RS fares against UED for identical KK masses.

Next we construct a variable $R = \sigma_{gg \to h} \Gamma_{h \to \gamma\gamma}$. In Fig. 3, we have studied variation of $(R_{RS} - R_{SM})/R_{SM}$ with m_h. For $m_{KK} = 1.0$, 1.5, 2.0 and 3.0 TeV, the fractional changes in R are 30%, 14%, 8% and 4%, respectively, for $m_h < 150$ GeV. The comparison shown in the inset shows that RS wins over UED roughly by a factor of 2 for identical KK scale for $m_h < 150$ GeV. Incidentally, our UED plots in the insets of all the three figures are in complete agreement with [8]. See also [9] for a numerical simulation of the Higgs signal at LHC in the UED context.

5. Conclusions

In conclusion, we highlight the core issues. In the RS scenario, the brane-bound Higgs can have order one Yukawa coupling with the KK fermions of all flavors. Such large KK Yukawa couplings can sizably enhance the Higgs production via gluon fusion and alter the Higgs decay width into two photons, provided the KK masses are in a regime accessible to the LHC. Because of the proactive involvement of more flavors inside the loop the effect in RS is significantly stronger (typically, by a factor of 2 to 2.5) than in UED for similar KK masses. Admittedly, this advantage in RS is somewhat offset by the fact that the lightest KK mass in UED can be as low as 500 GeV thanks to the KK-parity, while in RS a KK mass below 1.5 TeV would be difficult to accommodate (see below). However, attempts have been made to impose KK parity in warped cases as well [10].

Electroweak precision tests put a severe lower bound on m_{KK} (~ 10 TeV) [11]. To suppress excessive contribution to T and S parameters the gauge symmetry in the bulk is extended to $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, and then m_{KK} as low as 3 TeV can be allowed [12,13]. A further discrete symmetry $L \to R$ helps to suppress $Zb_L \bar{b}_L$ vertex correction and admits an even lower $m_{KK} \sim 1.5$ TeV [14]. If some other new physics (e.g. supersymmetrization of RS) can create further room in T and S by partial cancellation, $m_{KK} \sim 1$ TeV can also be accommodated. In our analysis, values of m_{KK} in the range of 1–3 TeV chosen for illustration may arise in the backdrop of such extended symmetries. Furthermore, if the b' quark, present in the case of left–right gauge symmetry, weighs around 1 TeV, one obtains an *additional* $\sim 10\%$ contribution to $\sigma(gg \to h)$ [15].

A very recent paper [16] lists the relative contribution of different scenarios (supersymmetry, flat and warped extra dimension, little Higgs, gauge–Higgs unification, fourth generation, etc.) to $gg \to h$ and $h \to \gamma\gamma$ for some benchmark values. A comparison between their work and ours in order. As regards the RS scenario, the authors of [16] consider the region of parameters where the zero mode quarks mix with their KK partners. Additionally, their choice of c_L is substantially different from c_R, where they observe large destructive interference in the effective ggh coupling. On the other hand, our working hypothesis is based on: $c \equiv c_L = c_R$ (see Eq. (6)), and we assume KK number conservation at the Higgs vertex. We observe that the Higgs coupling to KK quarks is large for any flavor (see Eq. (7)), and the (direct) loop effects of the KK quarks (which carry the same quantum numbers as their zero modes) do enhance the effective ggh vertex (like the *enhancement* observed for the fourth family contribution [16], or the b' quark contribution [15], or the UED contribution [8,16]), and the magnitude is rather insensitive to the value of c as long as $|c| \gtrsim 0.5$. The authors of [17] also calculate the KK-induced effective ggh vertex, but they rely on the gauge-Higgs unification set-up, and hence an efficient numerical comparison of their work with ours is not possible.

Acknowledgements

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