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Constraining minimal and non-minimal UED models with Higgs couplings

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Abstract

Early indications from the LHC for the observed scalar boson imply properties close to the Standard Model Higgs, putting considerable constraints on TeV scale new physics scenarios. In this letter we consider flat extra dimensional scenarios with the fifth spatial dimension compactified on an S^1/Z_2 orbifold. We find in the minimal model the experimentally preferred effective Higgs couplings to gluon and photon at 95% confidence level disfavor the New Physics scale below 1.3 TeV. We demonstrate that a generalization of these models to include brane localized kinetic terms can relieve the tension to accommodate scales as low as 0.4 TeV.

`Key Words: Extra dimension`

Introduction: Increasing evidence shows that the observed scalar boson at CERN as reported by ATLAS [1] and CMS [2] collaborations has properties that are close to the Standard Model (SM) Higgs. The era of precision measurements at LHC and future colliders that lie ahead will sit on judgment on the Higgs properties. However present trends indicate considerable constraints on New Physics (NP) scenarios that imply a modification of the Higgs production cross sections and branching ratios. Of particular importance are the loop driven Higgs couplings to the photon and the gluon that are susceptible to considerable corrections from TeV scale new physics. Indeed the impact of these couplings has been explored extensively in the recent past [3]. Aficionados of TeV scale new physics scenarios have been compelled to move to more general versions of specific models typically with a larger parameter space to accommodate these experimental constraints.

In this letter we focus on compactified extra dimensional models with a flat bulk profile [4]. These models have enjoyed considerable attention owing to their minimal nature, prediction of natural dark matter candidate [5] and ability to mimic supersymmetric extension of the SM at the collider [6], among other features. In their simplest avatar, the Universal Extra Dimension (UED) models extend the SM with two new parameters viz. (i) the size of the compactified extra dimension (R) and (ii) the cut-off of the effective theory (Λ). The effect of predicted new states, the so called Kaluza-Klein (KK) excitations of the SM fields, on the Higgs couplings have been studied in [7]. Corresponding studies in the context of warped 5d scenarios with explicit KK parity [8] or KK parity violating custodial models[3] [9] have also been made.

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[3]Note that in custodial Warped models, the sign of the contribution depends on the order in which the Higgs is regularized to the IR brane and the KK sum is made, see [10].

In this letter we discuss the constraints from the LHC Higgs results on the UED scenario, which we find to be quite severe. However generalizing to models with brane localized kinetic terms (BLKT) [11] that arise naturally due to the cut-off dependent radiative corrections, facilitates a considerable recovery of the constrained parameter space.

This letter is organized as follows. We begin by reviewing the formalism to study the Higgs couplings $gg \to h$ and $h \to \gamma\gamma$, including possible contributions from new states beyond the SM. Then we compute in turn the contributions from UED and BLKT scenarios and compare with present experimental values. Finally we conclude.

The loop driven Higgs couplings: At the LHC, the Higgs production chiefly proceeds through the gluon fusion process $gg \to h$, driven by the fermion triangle loop [12]. Within the SM the top loop contribution dominates. Similarly an important decay mode for a 125 GeV Higgs is the di-photon channel. This proceeds through a fermion and a W boson loop within SM. New states with correct quantum numbers can show up as virtual particles in these loops and may lead to a sizable correction of the effective couplings. The corresponding cross section and decay width, including possible contributions from new massive states are given below [14],

$$\begin{aligned} \Gamma_{h\to\gamma\gamma} &= \frac{G_F \alpha^2 m_H^3}{128\sqrt{2}\pi^3} \left| A_W(\tau_W) + 3\left(\frac{2}{3}\right)^2 A_t(\tau_t) \; + N_{c,NP} Q_{NP}^2 \, \mathcal{A}_{NP}(\tau_{NP}) \right|^2 , \\ \sigma_{gg\to h} &= \frac{G_F \alpha_s^2 m_H^3}{16\sqrt{2}\pi^3} \left| \frac{1}{2} A_t(\tau_t) \; + C(r_{NP}) \mathcal{A}_{NP}(\tau_{NP}) \right|^2 , \end{aligned} \tag{1}$$

where,

$$\mathcal{A}_{NP}(\tau_{NP}) = \sum_{NP} \frac{v}{m_{NP}} \frac{\partial m_{NP}}{\partial v} \; A_{F,V,S}(\tau_{NP}). \tag{2}$$

Note that the contribution of a state in the loops is dependent on the spin statistics of the new states (F, V and S stand for fermions, vector bosons and scalars respectively). The various contributions are given as,

$$\begin{aligned} A_F(\tau) &= \frac{2}{\tau^2}\left(\tau + (\tau - 1) f(\tau)\right) , \;\; A_V(\tau) = -\frac{1}{\tau^2}\left(2\tau^2 + 3\tau + 3(2\tau - 1) f(\tau)\right) , \\ A_S(\tau) &= -\frac{1}{\tau^2}\left(\tau - f(\tau)\right) , \end{aligned} \tag{3}$$

where, $\tau_i = m_h^2/4m_i^2$, and $f(\tau) = \arcsin^2 \sqrt{\tau}$. In the light Higgs limit i.e., $m_h << 2m_i$, we get, $A_F \sim 4/3, A_V \sim -7, A_S \sim 1/3$.

It will be useful to define two dimensionless parameters, $C_{gg} = \sigma^{NP}_{gg\to h}/\sigma^{SM}_{gg\to h}$ and $C_{\gamma\gamma} = \Gamma^{NP}_{h\to\gamma\gamma}/\Gamma^{SM}_{h\to\gamma\gamma}$, to compare the correction induced by New Physics scenarios with experimentally allowed values. The best fit values for C_{gg} and $C_{\gamma\gamma}$ from the LHC Higgs results were computed in [3]. We are going to use the numerical values quoted in [13]: $\sqrt{C_{gg}} = 0.88 \pm 0.11$ and $\sqrt{C_{\gamma\gamma}} = 1.18 \pm 0.12$.

Universal Extra Dimension: At lowest order, flat extra dimensional theories like the UED [4] with the extra spatial dimension compactified over a S^1/Z_2 orbifold, is a simple extension of the SM with a single new parameter, the radius of compactification (R) of the extra spatial dimension added to the theory. As 5d quantum field theories are essentially non-renormalizable, physical parameters are sensitive to the finite cut-off Λ of the theory. The end points $(x_5 = \pm\pi R/2)$ of the compactified extra dimension are the locations of four dimensional hyper-surfaces called the 3-branes. The SM can be embedded in the 5d scenario, the relevant part of the action be schematically represented by the following,

$$S = \int_{-\pi R/2}^{\pi R/2} dx_5 \int d^4x \Bigg\{ -\frac{1}{2}\sum_{i=1}^{3} \mathrm{Tr}\left[F_{iMN}F_i^{MN}\right] + (D_M H)^\dagger D^M H + \mu^2|H|^2 - \frac{\lambda_5}{4}|H|^4$$
$$+ i\,\bar{Q}_3 \not{D} Q_3 + i\,\bar{u}_3 \not{D}\, u_3 - \left[i\lambda_5^u \overline{Q}_3 u_3 \tilde{H} + \mathrm{h.c.}\right] \Bigg\}\,, \tag{4}$$

where, the Latin index runs from 0 to 3 and 5 (representing the flat extra dimension) the covariant derivative is defined as $D_M = \partial_M - i\sum_{i=1}^{3} g_5^i\, T_i^a A_{i\,M}^a$ and $\Gamma^M = \{\gamma^\mu, i\gamma^5\}$. The leptons can be included analogous to the quarks, which we omit for brevity. The 4d effective action can be derived by performing a Fourier expansion of the fields as a function of the compactified fifth dimension and integrating out the fifth dimension term by term. This procedure elevates every five dimensional bulk field to an infinite tower of increasingly massive states. Depending on the boundary values at the locations of the orbifold fixed points, the lowest lying states can be massless (zero modes) and are identified with the SM fields. The respective KK excitations, corresponding to higher Fourier modes, form an infinite tower of massive states. The tree level masses of the n-th KK excitation are given by ,

$$(m^{(n)})^2 = (m^{(0)})^2 + n^2/R^2. \tag{5}$$

For all the SM fields the zero mode mass can be written in the form, $m^{(0)} \propto gv$, where g is the gauge coupling for massive vector bosons, the Yukawa coupling for fermions and the self coupling for the Higgs, and v is the vacuum expectation value of the zero mode Higgs. The integer n corresponds to quantized momentum along the fifth dimension. 5d Lorentz invariance implies that every vertex preserves the quantum number n. However, orbifolding required to obtain chiral fermions at the zero mode breaks this symmetry, leaving a residual reflection symmetry in the 5d geometry. An important aspect of the symmetric nature of the fifth dimension is that the KK excitations have identical couplings as their zero modes and at every vertex KK parity is preserved.

We are interested in the contribution of the KK excitations that may show up as virtual particles and modify the loop level coupling of the zero mode scalar field that is identified with SM Higgs field. The contribution from the entire KK top tower is a convergent sum and can be computed in closed form in the light Higgs approximation by using Eqs. 2 and 5 and is given by [14],

$$\mathcal{A}_t^{NP} \sim \frac{8}{3}\left(\frac{\pi m_t/R\coth(\pi m_t/R) - 1}{2}\right). \tag{6}$$

In the case of the KK version of the W boson loop, one needs to be careful to also incorporate the contribution from the additional scalar states present in the spectrum. These are states that are a linear combinations of KK excitations of the fifth component of the gauge bosons $W_5^{(n)}$ and the KK excitations of the Goldstone modes $G^{(n)}$, that are orthogonal to the longitudinal component of the massive $W_\mu^{(n)}$ boson. A systematic study of this within the light Higgs approximation gives us,

$$\mathcal{A}_W^{NP} \sim -\frac{20}{3}\left(\frac{\pi m_W/R\coth(\pi m_W/R) - 1}{2}\right). \tag{7}$$

Using these expressions in Eq. 1, we compute the correction to the Higgs couplings as a function of the radius of compactifications of the fifth spatial dimension. Figure 1 shows that the measurement of the Higgs couplings

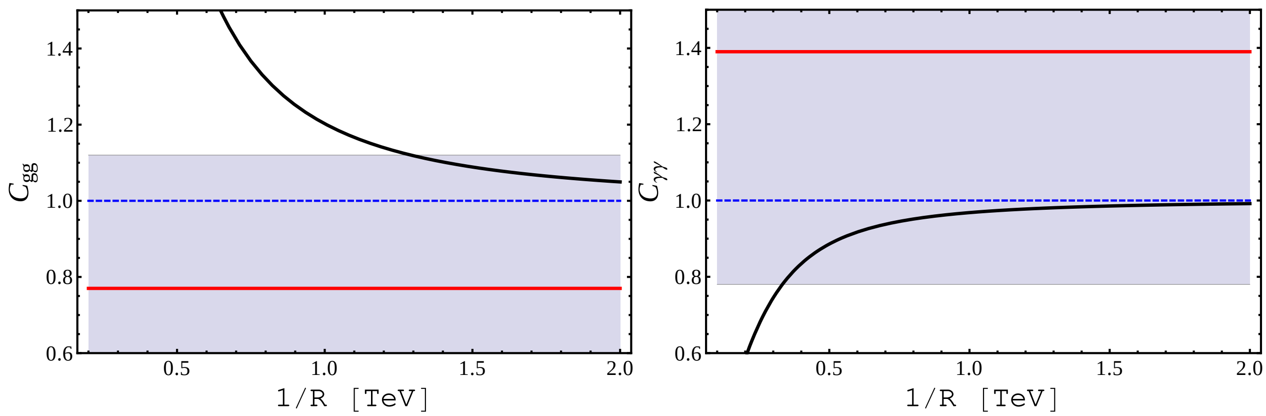

Figure 1: The ratios of the Higgs couplings in UED scenario to their SM values are plotted as functions of the inverse radius of compactification of the extra dimension (1/R). The gray (shaded) bands represent the 95% confidence level best fitted values for these ratios, with the red (light) horizontal lines representing the central values from LHC data. The blue (dashed) lines correspond to the SM points. The black (dark) curves are the UED predictions. We have assumed $m_h = 125$ GeV.

to gluons at 95% confidence disfavors $1/R < 1.3$ TeV. The constraints from the effective couplings to photons are smaller primarily owing to the partial cancellation of the contributions between the KK fermions and KK gauge-Goldstone bosons. This should be compared with other constraints which are typically strongest from the oblique electroweak corrections estimated to be around: $1/R > 0.8$ TeV [15]. The bounds from direct searches at colliders typically lead to sub-dominant or comparable constraints [16], owing to the relatively compressed mass spectrum within this class of models. Recent LHC Higgs mass bounds can constrain UED from the renormalization group running of the physical parameters [19]. For constraints from vacuum stability bounds see [20]. It may be noted that an exclusion limit at $1/R > 1.3$ TeV closes in on the overclosure limit from dark matter relic abundance in the minimal model [17]. Admittedly this limit may be avoided by reorganizing the mass spectrum in non-minimal models, as demonstrated in [18].

Thus we find that if the current trend of the Higgs data gets support in future with increased statistics, it will result in tighter bounds than from direct observations and pose a challenge to the parameter space of these models that is accessible to present and future collider experiments.

Brane Localized Kinetic Terms: In higher dimensional theories with lower dimensional defects, one expects localized divergent terms at the location of the defects. This arises due to the violation of translation invariance at these locations leading to UV sensitive radiative corrections that run with the scale and cannot be removed simultaneously at all scales. A way to quantify these corrections is to consider an effective theory at the cut-off scale, with higher dimensional operators that are allowed by the standard model gauge symmetries and Lorentz invariance to be added to the 5d theory. Such operators can be in the bulk or localized at the orbifold fixed points. Purely from the universality class that they belong to the brane localized kinetic terms, being relevant operators due to their mass dimension, lead to the most important corrections at low energy. The renormalization group running will amplify these operators at the infrared. From a phenomenological perspective these operators have considerable impact on the phenomenology of these models. They may lead to novel signals at colliders [21, 23], thus providing a way to probe these models at the LHC.

We introduce a minimal correction to the fermionic and Yukawa sector of UED Lagrangian given in Eq. 4, as given by,

$$\Delta S = \int_{-\pi R/2}^{\pi R/2} dx_5 \int d^4x \left\{ r_Q \Big(\delta(x_5 - \pi R/2) + \delta(x_5 + \pi R/2) \Big) \left[i\overline{Q}_3 \not{D}\, Q_3 + i\overline{u}_3 \not{D}\, u_3 \right] \right.$$
$$\left. - \Big(r_Y (\delta(x_5 - \pi R/2) + \delta(x_5 + \pi R/2)) \Big) \times \left[i\lambda_5^u \overline{Q}_3 u_3 \tilde{H} + \text{h.c.} \right] \right\}. \tag{8}$$

Though originating from radiative corrections, in this letter we will consider the size (r_Q, r_Y) of the brane operators at the weak scale to be free and independent parameters. Inclusion of these terms can lead to deviations in the bulk profile, couplings and masses of the bulk fields. Note that 5d gauge invariance imply universal BLKT parameters. The genesis of this lies in the observation that for general BLKT parameters the zero mode bulk profile of the massive gauge bosons are non-flat [11]. To avoid such complication we have taken same BLKT parameter for the doublet and the singlet states in Eq. 8.The KK masses obtained by performing the KK decomposition and solving the equations of motion with appropriate boundary conditions now shift from their UED predictions, and are given by the subsequent roots of the following transcendental equations,

$$r_Q m^{(n)} = \begin{cases} -\tan\left(\frac{m^{(n)}\pi R}{2}\right) & \text{for } n \text{ even,} \\ \cot\left(\frac{m^{(n)}\pi R}{2}\right) & \text{for } n \text{ odd,} \end{cases} \tag{9}$$

In the present context, the Yukawa sector that couples the fermions with non-trivial bulk profile with the zero mode of the Higgs that is assumed t o be essentially flat in the bulk, is of interest and can be written as,

$$S_{Yukawa} = -\frac{v}{\sqrt{2}}\lambda \int d^4x \left\{ \left(\overline{Q}_L^{(0)} u_R^{(0)} + r'_{Qnn} \overline{Q}_L^{(n)} u_R^{(n)} - R'_{Qnn} \overline{u}_L^{(n)} Q_R^{(n)} + \ldots \right) + \text{ h.c.} \right\}, \tag{10}$$

where, $r'_{Qnn}(r_Q, r_Y, m^{(n)})$ and $R'_{Qnn}(r_Q, r_Y, m^{(n)})$ are overlap integrals obtained by introducing the bulk profile of the fermions into Eq. 8 and integrating the fifth dimension. The effects of possible KK level mixing terms in Eq. 10 will be discussed seperately. The overlap inegrals can be adopted from the expression given in [21] and are given by,

$$r'_{Qnn}(\text{e/o}) = \frac{2r_Q + \pi R}{2r_Y + \pi R} \left(\frac{2r_Y + \frac{1}{A^2_{Q^{(n)}}} \left[\frac{\pi R}{2} \pm \frac{1}{2m_Q^{(n)}} \sin(m_Q^{(n)} \pi R) \right]}{2r_Q + \frac{1}{A^2_{Q^{(n)}}} \left[\frac{\pi R}{2} \pm \frac{1}{2m_Q^{(n)}} \sin(m_Q^{(n)} \pi R) \right]} \right),$$
$$R'_{Qnn}(\text{e/o}) = \frac{2r_Q + \pi R}{2r_Y + \pi R} \left(\frac{2r_Y (B_{Q^{(n)}})^2 + \frac{1}{A^2_{Q^{(n)}}} \left[\frac{\pi R}{2} \mp \frac{1}{2m_Q^{(n)}} \sin(m_Q^{(n)} \pi R) \right]}{\frac{1}{A^2_{Q^{(n)}}} \left[\frac{\pi R}{2} \mp \frac{1}{2m_Q^{(n)}} \sin(m_Q^{(n)} \pi R) \right]} \right), \tag{11}$$

where, the subscript (e/o) represents whether n is even or odd. $A_{Q^{(n)}} = \sin\left(m_Q^{(n)} \pi R/2\right)$, for n odd and $\cos\left(m_Q^{(n)} \pi R/2\right)$, for n even, similarly $B_{Q^{(n)}} = \cot\left(m_Q^{(n)} \pi R/2\right)$, for n odd and $\tan\left(m_Q^{(n)} \pi R/2\right)$, for n even and $m_Q^{(n)}$ is the n-th root of Eq. 9. The mass matrix for the n-th KK excitation of the top quark can be written as,

$$S_{t^{(n)}} = -\int d^4x \left\{ \left[\overline{Q}_3^{(n)},\ \overline{u}_3^{(n)} \right]_L \begin{bmatrix} m_{Q_3}^{(n)} & r'_{Qnn} \frac{v}{\sqrt{2}} \lambda_t \\ -R'_{Qnn} \frac{v}{\sqrt{2}} \lambda_t & m_{u_3}^{(n)} \end{bmatrix} \begin{bmatrix} Q_3^{(n)} \\ u_3^{(n)} \end{bmatrix}_R + \text{ h.c.} \right\}. \tag{12}$$

In the above expressions we have assumed that in the 4d effective theory obtained by integrating the fifth dimension, the bulk fields $Q_L(x,x_5)$ and $u_R(x,x_5)$ now split into a massless zero mode and an infinite tower of massive 4d states given by $(Q_L^{(0)}, Q_L^{(n)}, Q_R^{(n)})$ and $(u_R^{(0)}, u_L^{(n)}, u_R^{(n)})$ respectively.

Note that introduction of the BLKT parameters lead to KK level mixing at the leading order in the Yukawa interactions. However the symmetric nature of the BLKT parameters as introduced in Eq. 8 confines the mixing within the odd or even modes, due to a residual unbroken KK parity in the theory. The mixing angle between the n-th KK level and $(n+2l)$-th KK level can be estimated as $\theta_{mix} \sim f(r_Y, r_Q, n, l) m_t R/2l$, where $f(r_Y, r_Q, n, l)$ schematically represents the corresponding overlap integrals. Though the mixing angle is supressed by the new physics scale ($1/R$), significant level mixing between zero mode and $n=2$ KK states are possible in certain regions of the parameters space specially for smaller values of $1/R$ [21]. For the collider implications of this mixing in the top sector see [22]. We now turn our attention to make a careful study of this phenomenon. The relevant part of the action can be written as,

$$S_{t^{(0-2)}} = -\int d^4x \Bigg\{ \left[\overline{Q}_3^{(0)}, \overline{Q}_3^{(2)},\, \overline{u}_3^{(2)}\right]_L \begin{bmatrix} \frac{v}{\sqrt{2}}\lambda_t & 0 & r'_{Q02}\frac{v}{\sqrt{2}}\lambda_t \\ r'_{Q20}\frac{v}{\sqrt{2}}\lambda_t & m_{Q_3}^{(2)} & r'_{Q22}\frac{v}{\sqrt{2}}\lambda_t \\ 0 & -R'_{Q22}\frac{v}{\sqrt{2}}\lambda_t & m_{u_3}^{(2)} \end{bmatrix} \begin{bmatrix} u_3^{(0)} \\ Q_3^{(2)} \\ u_3^{(2)} \end{bmatrix}_R + \text{ h.c.} \Bigg\}. \quad (13)$$

where,

$$r'_{Q20} = \frac{2r_Q + \pi R}{2r_Y + \pi R} \left(\frac{1}{\sqrt{2r_Q + \pi R}} \frac{2\,(r_Y - r_Q)}{\sqrt{2r_Q + \frac{1}{A^2_{Q^{(2)}}}\left[\frac{\pi R}{2} + \frac{\sin\left(m_Q^{(2)}\pi R\right)}{2m_Q^{(2)}}\right]}} \right) \quad (14)$$

The overlap integral between the zero and second level, r'_{Q20} is symmetric in its last two indices. In Eq 13, the identification of the lightest eigenvalue with the SM top with mass in the range, $m_t = 173.07 \pm 0.52 \pm 0.72$ GeV [25] puts a severe constraint on the allowed BLKT parameters and consequently constrains the Higgs couplings that we will now discuss.

In order to obtain the Higgs couplings, we diagonalize the mass matrix in Eq. 12, that gives us the physical mass of the n-th KK excitations. One can use this to compute the Higgs cross section and decay width in the BLKT scenario by using Eqs. 1 and 2. We perform the KK summation numerically and terminate the procedure at $n = 20$, as the contribution decouples with the KK number and becomes numerically insignificant beyond this. Note that the gauge and scalar parts of the Lagrangian are unaffected by the introduction of the BLKT action given in Eq. 8. For the gauge-Goldstone sector one can adopt the analytic expression in Eq. 7. As indicated above we take care to include the constraint from the mixing effect on the parameter space of the theory. We find that once the constraint on the top mass is imposed the numerical effect of the mixing on the Higgs coupling is insignificant. However we include the leading order contribution from the mixing effect by considering the 0-2 KK level mixing. This can be consistently included into the calculation by replacing the contribution of the second KK top to the sum in Eq. 2 by the following expression,

$$\mathcal{A}_{NP}^{(2)} = \frac{4}{3}\left[\sum_{j=1}^{3} \frac{v}{m_j}\frac{\partial m_j}{\partial v} - 1\right], \quad (15)$$

where, m_j-s are the three eigenvalues of the mass matrix in Eq. 13 with lowest eigenvalue identified with SM top mass.

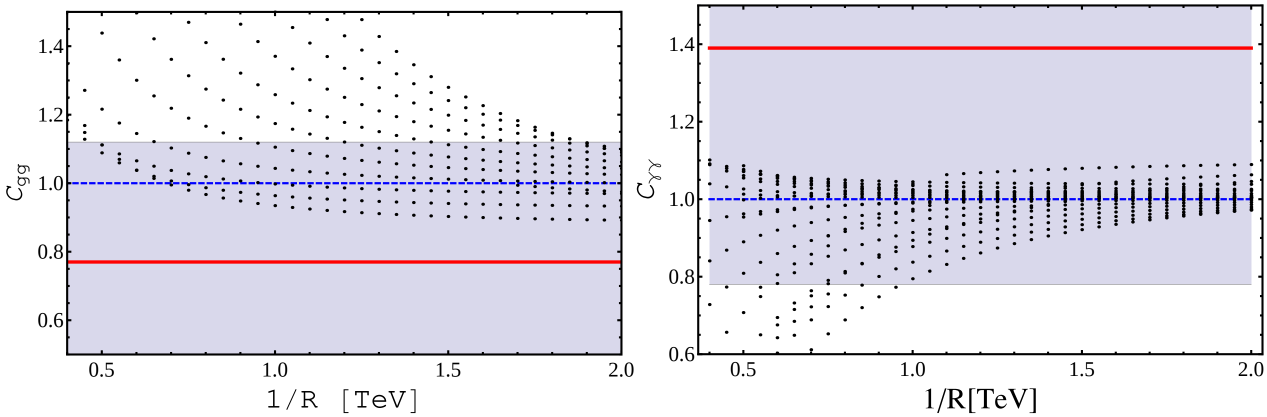

Figure 2: The ratios the Higgs couplings in BLKT scenario to their SM values are plotted as functions of the inverse radius of compactification of the extra dimension (1/R). The gray (shaded) bands represent the 95% confidence level best fitted values for these ratios, with the red (light) horizontal lines representing the central values from LHC data. The blue (dashed) lines correspond to the SM points. The black (dark) points represent the BLKT results. We have assumed $m_h = 125$ GeV. The BLKT parameters r_Q and r_Y are varied within the range $[-\pi R/2, \pi R/2]$.

After the KK sum is done the ratios C_{gg} and $C_{\gamma\gamma}$ still remain functions of the BLKT parameters (r_Q, r_Y) and the inverse radius of compactification $(1/R)$. We vary the BLKT parameters within the range $[-\pi R/2, \pi R/2]$ and obtain the corresponding scatter plots for C_{gg} and $C_{\gamma\gamma}$ as functions of $1/R$ in Figure 2. Expectedly the points form a band around the minimal UED prediction that corresponds to $r_Q = r_Y = 0$.

Crucially we find that in certain regions of the parameter space, the contribution from the KK fermions can change sign relative to the zero mode (SM) contribution. As can be seen from the plot this reduces constraint from the coupling, which was at 1.3 TeV for minimal UED models, now becoming 0.4 TeV. Significantly we find that the UED couplings are always disfavored over the SM predictions in all regions of the parameter space by the experimental data. In the extended scenario with the BLKT parameters, we find that in certain regions of the parameter space the couplings fit better than the SM. The limits on $1/R$, from dark matter relic density measurements and their direct detection at experiments within the BLKT framework, are rather model dependent and can lead to considerable relaxation over the minimal UED bound [24, 26]. The corresponding constraints on this class of models from electroweak precision measurements can be found in [11, 24, 27].

Conclusion: In this paper we have studied the impact of Higgs couplings as measured at the LHC on Universal Extra Dimension models. We find that the minimal models are particularly constrained from the Higgs coupling to the gluon. We make a simple extension of the model by introducing relevant brane localized kinetic terms. This leads to non-trivial 5d profiles for the bulk fields. The interactions are modified by the corresponding overlap integrals. In certain regions of the parameter space this can lead to better fitting of the experimental data implying a considerable relaxation of the experimental constraints.

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References

[1] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716** (2012) 1 [arXiv:1207.7214 [hep-ex]].

[2] S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716** (2012) 30 [arXiv:1207.7235 [hep-ex]], S. Chatrchyan *et al.* [CMS Collaboration], arXiv:1303.4571 [hep-ex].

[3] M. R. Buckley and D. Hooper, Phys. Rev. D **86** (2012) 075008 [arXiv:1207.1445 [hep-ph]], P. P. Giardino, K. Kannike, M. Raidal and A. Strumia, JHEP **1206** (2012) 117 [arXiv:1203.4254 [hep-ph]], S. Banerjee, S. Mukhopadhyay and B. Mukhopadhyaya, JHEP **1210** (2012) 062 [arXiv:1207.3588 [hep-ph]], G. Cacciapaglia, A. Deandrea, G. D. La Rochelle and J. -B. Flament, JHEP **1303** (2013) 029 [arXiv:1210.8120 [hep-ph]], P. P. Giardino, K. Kannike, I. Masina, M. Raidal and A. Strumia, arXiv:1303.3570 [hep-ph].

[4] I. Antoniadis, Phys. Lett. B **246** (1990) 377. T. Appelquist, H. C. Cheng and B. A. Dobrescu, Phys. Rev. D **64** 035002 (2001) [hep-ph/0012100], D. Hooper and S. Profumo, Phys. Rept. **453** (2007) 29 [hep-ph/0701197].

[5] G. Servant and T. M. P. Tait, Nucl. Phys. B **650** (2003) 391 [hep-ph/0206071].

[6] A. Datta, K. Kong and K. T. Matchev, Phys. Rev. D **72** (2005) 096006 [Erratum-ibid. D **72** (2005) 119901] [hep-ph/0509246].

[7] F. J. Petriello, JHEP **0205** (2002) 003 [arXiv:hep-ph/0204067], G. Belanger, A. Belyaev, M. Brown, M. Kakizaki and A. Pukhov, Phys. Rev. D **87** (2013) 016008 [arXiv:1207.0798 [hep-ph]], T. Kakuda, K. Nishiwaki, K. -y. Oda, N. Okuda and R. Watanabe, arXiv:1304.6362 [hep-ph], A. Datta and S. K. Rai, Int. J. Mod. Phys. A **23** (2008) 519 [hep-ph/0509277].

[8] G. Bhattacharyya and T. S. Ray, Phys. Lett. B **675** (2009) 222 [arXiv:0902.1893 [hep-ph]].

[9] A. Azatov, M. Toharia and L. Zhu, Phys. Rev. D **82** (2010) 056004 [arXiv:1006.5939 [hep-ph]]. F. Goertz, U. Haisch and M. Neubert, Phys. Lett. B **713** (2012) 23 [arXiv:1112.5099 [hep-ph]].

[10] M. Carena, S. Casagrande, F. Goertz, U. Haisch and M. Neubert, JHEP **1208** (2012) 156 [arXiv:1204.0008 [hep-ph]].

[11] T. Flacke, A. Menon and D. J. Phalen, Phys. Rev. D **79** (2009) 056009 [arXiv:0811.1598 [hep-ph]].

[12] A. Djouadi, Phys. Rept. **457** (2008) 1 [hep-ph/0503172].

[13] J. Ellis and T. You, arXiv:1303.3879 [hep-ph].

[14] G. Cacciapaglia, A. Deandrea and J. Llodra-Perez, JHEP **0906** (2009) 054 [arXiv:0901.0927 [hep-ph]].

[15] A. Belyaev, M. Brown, J. Moreno and C. Papineau, arXiv:1212.4858 [hep-ph].

[16] L. Edelhauser, T. Flacke and M. Kramer, arXiv:1302.6076 [hep-ph], T. Flacke, A. Menon and Z. Sullivan, Phys. Rev. D **86** (2012) 093006 [arXiv:1207.4472 [hep-ph]], A. Datta, A. Datta and S. Poddar, Phys. Lett. B **712** (2012) 219 [arXiv:1111.2912 [hep-ph]].

[17] G. Belanger, M. Kakizaki and A. Pukhov, JCAP **1102** (2011) 009 [arXiv:1012.2577 [hep-ph]].

[18] K. Kong and K. T. Matchev, JHEP **0601** (2006) 038 [hep-ph/0509119].

[19] M. Blennow, H. Melbeus, T. Ohlsson and H. Zhang, Phys. Lett. B **712** (2012) 419 [arXiv:1112.5339 [hep-ph]].

[20] A. Datta and S. Raychaudhuri, Phys. Rev. D **87** (2013) 035018 [arXiv:1207.0476 [hep-ph]].

[21] A. Datta, K. Nishiwaki and S. Niyogi, JHEP **1211** (2012) 154 [arXiv:1206.3987 [hep-ph]].

[22] A. Datta, K. Nishiwaki and S. Niyogi, in preparation.

[23] A. Datta, U. K. Dey, A. Shaw and A. Raychaudhuri, Phys. Rev. D **87** (2013) 076002 [arXiv:1205.4334 [hep-ph]].

[24] T. Flacke, K. Kong and S. C. Park, arXiv:1303.0872 [hep-ph].

[25] J. Beringer *et al.* [Particle Data Group Collaboration], Phys. Rev. D **86** (2012) 010001.

[26] A. Datta, U. K. Dey, A. Raychaudhuri and A. Shaw, Phys. Rev. D **88** (2013) 016011 [arXiv:1305.4507 [hep-ph]].

[27] F. del Aguila, M. Perez-Victoria and J. Santiago, Acta Phys. Polon. B **34** (2003) 5511 [hep-ph/0310353].