Bulk RS models, Electroweak Precision tests and the 125 GeV Higgs

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Abstract

We present upto date electroweak fits of various Randall Sundrum (RS) models. We consider the bulk RS model, deformed RS and the custodial RS models. For the bulk RS case we find the lightest Kaluza Klein (KK) mode of the gauge boson to be $\sim 8$ TeV while for the custodial case it is $\sim 3$ TeV. The deformed model is the least fine tuned of all which can give a good fit for KK masses $< 2$ TeV depending on the choice of the model parameters. We also comment on the fine tuning in each case.
I. INTRODUCTION

The discovery of the Higgs boson at \( \sim 125 \text{ GeV} \) has firmly established the status of the Standard Model (SM) as ‘the’ theory of electroweak scale physics. While an increase in experimental precision may serve as an indirect probe for new physics, they are also effective in constraining several new physics models. In this work we focus on the Randall-Sundrum (RS) model [1] and its variations. The RS model, in its original form, was based on a slice of AdS\(_5\) spacetime between two 3-branes, referred to as the UV and IR branes with a tiny inter-brane separation just somewhat larger than the Planck length. The RS geometry is warped and this helps provide a natural answer to the gauge-hierarchy problem by ”warping down” the fundamental five-dimensional scale which is of the order of \( M_P \) down to the weak scale.

In the original model, with a view to seeking a solution to the hierarchy problem, only gravity was allowed to propagate in the bulk while the Standard Model (SM) fields were on the brane. It was soon realized that allowing gauge bosons in the bulk facilitated unification of couplings [2]. But in this case, since the Higgs and the fermions were localized on the IR brane, they coupled maximally to the gauge KK states which are also localized near the IR brane. This ‘strong coupling’ lead to unacceptably large corrections to the oblique \( S \) and \( T \) parameters [3] resulting in bounds on the first KK mass in excess of 30 TeV [4–9]. Moving the fermions into the bulk served the following two purposes:

1) It offered an elegant solution to the Yukawa hierarchy puzzle achieved by localizing the fermions at different points in the bulk, resulting in interesting flavour phenomenology in the hadronic sector [10–14] and the leptonic sector [11, 15–27]. For a detailed description of RS phenomenology with bulk fields see [28, 29].

2) The constraints on the gauge KK states from the \( S \) parameter is significantly weakened as the all the light fermions except the top are localized away from the IR brane and Higgs.

The constraints from \( T \) parameter, however remain strong as the Higgs doublet is localized near the IR brane which is necessary for the solution to the hierarchy problem. In view of this the following extensions were proposed:

- Models with bulk custodial symmetry [30] The bulk gauge group in question \( SU(2)_L \times SU(2)_R \times U(1)_X \). In this case the additional corrections to the \( T \) parameter due to new KK gauge bosons cancel the volume enhanced contributions due to the KK states of the SM gauge bosons. The \( T \) parameter vanishes at tree level and the limits on the KK mass of the first gauge boson is mainly due to the \( S \) parameter. A straight forward estimation of
the $S$ parameter results in a lower bound on the first KK mass to be $\sim 4$ TeV for the point to lie inside the $3\sigma$ region in Figure 1. Taking into account the loop corrections to the $T$ parameter, (in scenarios with $Zbb$ protection ) it was found that one can lower the mass of the first KK gauge boson to around 3 TeV at $3 \sigma$. An additional is to consider custodial models with gauge-Higgs unification which can address the little hierarchy problem. [31]. A global fit to the precision observables in such models was performed in [32].

- Models with a deformed metric: [33–35] In this setup the bulk geometry is RS like (AdS) near the UV brane while there is a deviation from AdS geometry near the IR brane. Depending on the model parameters this often results in a smaller volume factor as compared to the original RS setup. In [36] the authors performed a fit to the data for different values of the Higgs mass and evaluated the fine tuning required to the fit that particular Higgs mass.

- Models with brane localized kinetic terms for the gauge bosons [37].

In this work we establish the current of the RS models taking into account the most recent measurements of precision electroweak observables We focus our attention on the scenarios a) and b) of the extensions to the RS model. In both the scenarios the brane mass parameter of the Higgs doublet plays an important role in determining constraint on the first KK mass of the gauge boson. Additionally with the precise measurement of the Higgs mass, the extent of fine tuning is related to the brane mass parameter $b$ [34–36]. We make explicit the interplay between the fine tuning required to fit the Higgs boson mass and the $b$ parameter which gives the best fit to the electroweak observables. Using the expansion formalism of [38] we determine the best fit points for the model parameters by using the Standard $\chi^2$ analysis. We find that the KK mass of the first gauge boson is lower than what was obtained in earlier analysis when naive bounds from evaluation the $S$ and $T$ parameters were taken into account.

The paper is organized as follows: In Section II and Section III we outline the formalism of [38] thus providing the necessary background for the analysis. In section IV we briefly review the bulk RS model and study the bulk RS model with no additional gauge symmetries. In Section V and Section VI we analyse RS model with deformed metric and custodial symmetry respectively. In Section VII we conclude.
II. EXPANSION FORMALISM AND THE SM

In this section we briefly review the expansion formalism of [38] which we use for our analysis. There are numerous observables in the Standard Model whose values have been very well measured. These observables are in general a function of the following lagrangian parameters:

\[ p_k' \equiv \{ g_i, y_t, v, \lambda \} \quad (1) \]

where \( g_i \) are the gauge couplings, \( y_t \) is top quark Yukawa coupling, \( v \) is vacuum expectation value (vev) and \( \lambda \) is the quartic coupling. These parameters are referred to as the ‘input parameters’. An \( i^{th} \) observable, \( \hat{O}^{SM}_i \) in the SM can expressed as a function of these parameters as

\[ \hat{O}^{SM}_i(p_k') = \hat{O}^{ref}_i + \sum_{k'} \frac{\partial \hat{O}^{SM}_i}{\partial p_{k'}} (p_{k'} - p^{ref}_{k'}) + \ldots \quad (2) \]

where the \ldots denote higher orders. \( \{ p^{ref}_{k'} \} \) is the set of lagrangian parameters chosen at a reference value \( p^{ref}_{k'} \) at which the evaluated expressions for the SM observables match closely with experiment and \( \hat{O}^{ref}_i = \hat{O}^{SM}_i(\{ p^{ref}_{k'} \}) \). Thus the expansion in Eq.(2) is about the reference values \( p^{ref}_{k'} \) and \( p_{k'} \) is the allowed deviation about the reference values.

The Lagrangian parameters, however are not measured directly, but are extracted from the measurements of certain observables. As a result it seems logical to re-express the SM observables in terms of a few accurately measured observables which will now serve as the input. One such list of input observables is

\[ \hat{O}_{k'} \equiv \{ m_Z, m_H, G_F, \alpha(m_Z), \alpha_s(m_Z), m_t(m_t) \} \quad (3) \]

In terms of the input observables, Eq.(2) can be re-expressed as follows:

\[ \hat{O}^{SM}_i(\{ \hat{O}_{k'} \}) = \hat{O}^{ref}_i + \sum_{k'} \frac{\partial \hat{O}^{SM}_i}{\partial \hat{O}_{k'}} (\hat{O}_{k'} - \hat{O}^{ref}_{k'}) + \ldots \quad (4) \]

where \( \hat{O}^{ref}_{k'} \) is the experimentally measured central value of the input observable and \( \hat{O}_{k'} \) quantifies the deviation from the central value. Thus the deviation in \( \hat{O}^{SM}_i \) can be expressed in terms of experimental deviation of the input observables from their central values. The relative deviation can be defined as

\[ \delta^{SM}_{\hat{O}_i}(\{ \hat{O}_{k'} \}) = \frac{\hat{O}^{SM}_i(\{ \hat{O}_{k'} \}) - \hat{O}^{ref}_i(\{ \hat{O}^{ref}_{k'} \})}{\hat{O}^{ref}_i} \quad (5) \]

\(^{1}\) A subset of these observables can be used to ‘determine’ the input parameters.
Defining
\[ c_{ik'} = \frac{\hat{O}_{k'}^{ref}}{\hat{O}_{i}^{ref}} \partial \hat{O}_{i}^{SM} \partial \hat{O}_{k'} \] (6)
we can express the relative deviation in Eq.(5) in the \( i \)th observable, due to deviation from the central value of the input observables as
\[ \bar{\delta}O_{SM}^{i} = \sum_{i'} c_{ii'} \delta \hat{O}_{SM}^{i'} \] (7)
Using this, Eq.(4) can be re-written as
\[ \hat{O}_{SM}^{i}(\{\hat{O}_{k'}\}) = \hat{O}_{i}^{ref}(1 + \bar{\delta}SM \hat{O}_{i}) \] (8)
The deviation of all the SM observables in can be quantified by constructing a \( \chi^2 \) statistic defined as
\[ \chi^2(\hat{O}_{k'}) = \sum_{i} \left[ \frac{\hat{O}_{SM}^{i}(\hat{O}_{k'}) - \hat{O}_{i}^{expt}}{\delta \hat{O}_{i}^{expt}} \right]^2 \] (9)
It should be noted that while constructing the \( \chi^2 \) we assume that there is no correlation \([39]\) between the output observables. However we note that taking into account correlation matrix for the output observables will not significantly change the results of our analysis. The central values and the allowed deviation for the input and output observables are given in Table[I]. Using the Z pole observables and the W mass as the output observables, we minimize the \( \chi^2 \) function in Eq.(9) by varying the input observables within the experimentally allowed deviation given in Table[I]. The minimization is performed using \([40]\). We obtain \( \chi^2_{min} = 24.54 \) with the corresponding best fit values for the input observables given in Table[II].

III. NEW PHYSICS

The expansion formalism presented in Section[II] can be extended to include new physics effects. Assuming the nature of new physics is such that it modifies mostly the oblique parameters, the new physics effects can be parametrized by the introduction of higher dimension operators in the lagrangian. These operators can give corrections to any of the input and the output observables. \[ In the presence of new physics Eq. [7] becomes \]
\[ \delta^{NP} \hat{O}_{i}^{th}(\{\hat{O}_{k'}\}, NP) = \bar{\delta} \hat{O}_{SM}^{i} + \zeta_i \] (10)
\[ In the calculations used for our analysis only tree level effective theory operators are considered. Things could get more stringent if one loop effective theory operators are considered [44–46]. \]
| Input observables | $m_Z$ | 91.1876(21) | $G_F$ | 1.1663787(6) $\times 10^{-5}$ |
|-------------------|-------|-------------|-------|-------------------------------|
| $\alpha(m_Z)$    | 7.81592(86) $\times 10^{-3}$ | $m_t(m_f)$ | 173.20(87) |
| $\alpha_s(m_Z)$  | 0.1185(6) | $m_H$ | 125.9(4) |

| Output observables | $m_W$ | 80.385(15) | $\Gamma_Z$ | 2.4952(23) |
|-------------------|-------|------------|-------------|-------------|
| $\sigma_{\text{had}}$ | 41.541(37) | $R_e$ | 20.804(50) |
| $R_\mu$ | 20.785(33) | $R_{\tau\mu}$ | 20.764(45) |
| $R_b$ | 0.21629(66) | $R_c$ | 0.1721(30) |
| $\sin^2\theta_c$ | 0.23153(16) | $\sin^2\theta_b$ | 0.281(16) |
| $\sin^2\theta_e$ | 0.2355(59) | $A^e_{FB}$ | 0.0145(25) |
| $A^b_{FB}$ | 0.0992(16) | $A^c_{FB}$ | 0.0707(35) |
| $A_b$ | 0.923(20) | $A_c$ | 0.670(27) |

| Input observables | $m_Z$ | 91.1914 | $G_F$ | 1.1663784 $\times 10^{-5}$ |
|-------------------|-------|---------|-------|----------------------------|
| $\alpha(m_Z)$    | 7.816714 $\times 10^{-3}$ | $m_t(m_f)$ | 173.34 |
| $\alpha_s(m_Z)$  | 0.119073 | $m_H$ | 125.405 |

| Output observables | $m_W$ | 80.366 $\pm$ 0.005 | $\Gamma_Z$ | 2.4957 $\pm$ 0.0006 |
|-------------------|-------|----------------|-------------|-----------------|
| $\sigma_{\text{had}}$ | 41.472 $\pm$ 0.04 | $R_e$ | 20.7427 $\pm$ 0.03 |
| $R_\mu$ | 20.7428 $\pm$ 0.03 | $R_{\tau\mu}$ | 20.7897 $\pm$ 0.03 |
| $R_b$ | 0.215822 $\pm$ 0.00005 | $R_c$ | 0.17209 $\pm$ 0.000007 |
| $\sin^2\theta_c$ | 0.23161 $\pm$ 0.000002 | $\sin^2\theta_b$ | 0.2329 $\pm$ 0.00001 |
| $\sin^2\theta_e$ | 0.2355 $\pm$ 0.000002 | $A^e_{FB}$ | 0.0160 $\pm$ 0.00004 |
| $A^b_{FB}$ | 0.0992 $\pm$ 0.00002 | $A^c_{FB}$ | 0.0732 $\pm$ 0.00001 |
| $A_b$ | 0.923 $\pm$ 0.0005 | $A_c$ | 0.6675 $\pm$ 0.00008 |

TABLE I: Experimentally measured central values for the input and output observables along with the Standard Deviation [41–43]

TABLE II: Best fit values for the input and output observables for the SM fit with $\chi^2_{\text{min}} = 24.54$

where $\zeta_i(\{\hat{O}_k\}, NP) \equiv \frac{\delta^{NP}\hat{O}^{th}_i(\{\hat{O}_k\}, NP)}{\hat{O}^{th}_i(\{\hat{O}_k\}, NP)}$ parametrizes the relative contribution to the $i^{th}$ observable due to higher dimension operators. Using Eq. (7), Eq. (10) can be written as

$$
\delta^{NP}\hat{O}^{th}_i(\{\hat{O}_k\}, NP) = \sum_i c_{ii'}\delta\hat{O}^{SM}_{i'} + \zeta_i
$$

$$
= \sum_i c_{ii'}\delta\hat{O}^{th}_{i'} + \delta^{NP}\hat{O}_i
$$

(11)
where $\delta^{NP}_i \hat{O}_i = \zeta_i - \sum_{i'} c_{i'i} \zeta_{i'}$. Here the superscript 'th' denotes SM in addition to new physics. Note that from Eq.(6), the matrix of co-efficients $c_{ik}$ is a unit matrix for the input observables i.e. $c_{ik}^{i'} = \delta_{ik}$. Thus any new physics effects to the input observables are adjusted such that the net shift is zero, which is apparent in Eq.(11). This adjustment is however propagated in the evaluation of the output observables through Eq.(11).

In many scenarios, new physics is such that their dominant contribution to the various SM observables is only through the self energy corrections to the various gauge boson propagators given below:

$$\pi_{\alpha\beta} = \{\pi_{ZZ}, \pi'_{ZZ}, \pi_{\gamma Z}, \pi'_{\gamma}, \pi_{WW}, \pi^0_{WW}\} \quad (12)$$

The primed quantities denotes differentiation with respect to $q^2$, where $q$ is the four momentum. Note that the corrections to the fermion coupling to the gauge bosons are universal. In this case the new physics contribution to the input observables in Eq.(11) can be re-expressed as

$$\delta^{NP}_i \hat{O}^{th}_i = \sum_{\alpha\beta} b_{i,\alpha\beta} \delta^{NP}_i \pi_{\alpha\beta} \quad (13)$$

where it is understood that the sum extends over the list in Eq.(12) while the co-efficients $b_{\alpha\beta}$ are evaluated in [38].

In such models the corrections to the gauge boson propagators can be encoded in oblique parameters $S$ and $T$ [3, 47]. These oblique parameters are related to the new physics effects to the self energy correction as follows [48]:

$$\delta^{NP}_i \pi_{ZZ} = -\alpha(m_Z)T + \frac{\alpha(m_Z)}{2}S$$

$$\delta^{NP}_i \pi'_{ZZ} = \frac{\alpha(m_Z)}{2}S$$

$$\delta^{NP}_i \pi_{\gamma Z} = -\frac{\alpha(m_Z)}{4\sin^2\theta_W} \cos2\theta_W \tan\theta_W S$$

$$\delta^{NP}_i \pi'_{\gamma\gamma} = -\frac{\alpha(m_Z)}{2}S$$

We use Eq.(13) in Eq.(11) to construct the $\chi^2$ for the output observables at the Z peak along with the W mass. Using the results of the analysis in [38], the expression for the $\chi^2$ statistic defined in Eq.(9) is given as

$$\chi^2 = 25.0898 + 1102.39 \ S^2 + 28.746 \ S - 72.0085 \ T - 2256.69 \ ST + 1377.07 \ T^2 \quad (14)$$

\footnote{The contribution to $U$ is suppressed as only dimension 8 operators contribute to it}
The input observables were fixed to their experimentally measured central values while obtaining the above expression. Using this we obtain the S-T plot in Fig. 1 in which the 68%, 95% and 99% confidence level allowed regions are depicted by red, blue and orange regions respectively.

![Figure 1](image)

FIG. 1: The red, blue and orange regions denote the 68%, 95% and 99% confidence level allowed regions in the ST parameter space.

For a particular model of new physics, the oblique parameters $S, T$ depend on the model parameters. A given set of model parameters is valid only if the corresponding $S, T$ observables computed for that set lie at least within the orange ellipse in Fig. 1. Thus a very small contribution, for example to the $S$ parameter would necessitate $T$ to also be very small so as to lie within the bottom left portion of the ellipse. However an increasing $S$ can admit larger values of the $T$ parameter corresponding to moving towards the top right portion of the ellipse. Thus we can use Fig. 1 to constrain the model parameters. We now use this analysis to obtain constraints on various Randall-Sundrum models.

IV. RANDALL-SUNDRUM MODELS

Randall-Sundrum model is a model of a single extra-dimension compactified on an $S_1/Z_2$ orbifold [1]. The five dimensional gravity theory is defined by the following line element:

$$ds^2 = e^{-2A(y)}\eta_{\mu\nu}dx^\mu dx^{\nu} - dy^2$$

(15)
Two opposite tension branes are located at the two fixed points of the orbifold. The space between the branes is endowed with a large negative bulk cosmological constant making it a slice of AdS. The presence of brane localized sources of energy results in zero cosmological constant being induced on the branes. In the original setup \( A(y) = ky \) where \( k \) is reduced Planck scale. Identifying the scale of physics on the \( y = 0 \) brane as \( M_{IR} \), the effective UV scale induced at the \( y = \pi R \) brane owing to geometry is given as

\[
M_{IR} = e^{-kR\pi} k
\]  

(16)

where \( R \) is the compactification radius. Choosing \( kR \sim 12 \), will result in \( M_{IR} \sim 200 \text{ GeV} \) owing to large exponential warping. Any radiative instability to the masses of fundamental scalars in the theory can be warped down to the electroweak scale thus solving the gauge hierarchy problem. In the original setup, with the exception of gravity all the SM fields were localized on the brane at \( y = y_1 = \pi R \) also referred to as the IR brane. Here we consider a generalization of the original setup where particles of all types of spin are allowed to propagate in the bulk.

A bulk field \( \Psi^s(x^\mu, y) \) with spin \( s \) can be expanded in the KK basis as follows:

\[
\Psi^s(x^\mu, y) = \frac{1}{\sqrt{\pi R}} \sum_{n=0}^{\infty} \psi_s^{(n)}(x^\mu) f_s^{(n)}(y)
\]  

(17)

The zero modes for the fields are identified as the SM fields. While the zero mode for the gauge bosons are flat at leading order, the ones for the scalars and the fermions are controlled by the brane and bulk mass terms respectively. They are parametrized as \( m_{brane}^s = bk \) and \( m_{bulk}^f = ck \) where \( b, c \) are dimensionless \( O(1) \) quantities. The normalized profiles for the fields are given as

\[
f^{(0)}_0(b, y) = \sqrt{\frac{2(b-1)}{e^{2(b-1)kR\pi} - 1}} e^{(b-1)ky}
\]

\[
f^{(0)}_{1/2}(c, y) = \sqrt{\frac{(1-2c)kR\pi}{e^{(1-2c)kR\pi} - 1}} e^{(0.5-c)ky}
\]

\[
f^{(0)}_1(y) = 1
\]  

(18)

where the normalization conditions are given as

\[
\frac{1}{\pi R} \int_0^{\pi R} (f^{(0)}_s(y))^2 dy = 1
\]  

(19)

\( c < 0.5 \) and \( b > 1 \) (\( c > 0.5 \) and \( b < 1 \)) correspond to the fields being localized towards the IR(UV) brane respectively. The KK modes of all fields are however localized near the IR brane. We note here that while the profiles of the gauge boson fields are flat at leading order, it receives corrections
due to the mixing of the KK mode with the zero mode. The mixing is proportional to the vacuum expectation value (vev) and is given as

\[ a_{01} = \frac{v^2}{(kR)M_{KK}^2} \int_{1}^{z_{IR}} z^2 f_0(0)(z)^2 f_1(0)(z)^2 f_1(1)(z) \]  

(20)

where \( z = e^{kRy} \) is the conformal co-ordinate. \( f_{0,1}^{(0)} \) are given in Eq.(18), while \( f_1^{(1)} \) is the profile of the first KK mode of the gauge boson. Detailed review about bulk RS models can be found in [29, 49]. Thus diagonalizing the mass matrix of KK modes and zero mode, will result in the lightest state (identified as the SM boson) having a small KK component proportional to Eq.(20).

As a result the coupling of the fermions to the SM boson will have a non-universal component which is a function of its localization parameter \( c \). The \( c \) parameters will be in general different for different fermionic generations to generate the required hierarchy in the Yukawa parameters. For the light fields with the exception of the top it is fair to assume \( c > 0.5 \) to reduce the overlap with the Higgs. For \( c > 0.5 \), non-universal component of the coupling is very small and can be neglected [29, 50]. This is enough to evade bounds from FCNC processes which can occur at tree level.

Solution of the hierarchy problem requires the Higgs zero mode to be localized very close to the IR brane. It corresponds to a choice \( b \geq 2 \) for the brane mass parameter [35, 51]. For a bulk scalar field with a massless zero mode the brane mass parameter \( b \) is related to the bulk mass parameter \( a \) as

\[ b = 2 \pm \sqrt{4 + a} \]  

(21)

Henceforth, we will drop the \( b = 2 - \sqrt{4 + a} \) solution as it will never lead to \( b \geq 2 \) necessary for the solution to the hierarchy problem. The zero mode increasingly becomes sharply localized near the IR brane as \( b \) increases. However an increase in \( b \) is only facilitated by the corresponding increase in \( |a| \). Depending on the value of \( \frac{k}{M_{Pl}} \), the bulk mass parameter cannot be increased indefinitely as the product \( ak \) will become greater than the 5D Planck scale. Fig.2 shows a plot of \( b \) as function of \( a \). Depending on the value of \( \frac{k}{M_{Pl}} \), the plot is terminated on the right at which \( ak = M_{Pl} \). For instance for \( \frac{k}{M_{Pl}} = 0.1 \), the plot (blue curve) is terminated at \( b = 5.74 \) while for \( \frac{k}{M_{Pl}} = 0.25 \) the plot (red curve) is terminated \( b = 4.82 \).

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4 Realization of EWSB also requires \( b > 2 \) [52].

5 Bulk scalar mass is parametrized as \( m_{bulk} = ak \)
Thus the case with brane localized Higgs will be treated separately and not as a limiting case where the bulk Higgs field tends to a brane localized one.

We now proceed to study the impact of EWPT in various RS models.

A. Bulk Higgs with no additional symmetries

This model is the same same setup discussed above. All the fermionic fields except the top are localized near the UV brane. This is sufficient to fit the masses of all fermions except the top. Due to localization of all the fermions near the UV brane the vertex corrections are very small and universal, thus the new physics effects can be parametrized in terms of the oblique operators $S$ and $T$ which are given as [34-36]

$$\alpha T = \frac{\sin^2 \theta_W m_Z^2 y_1 k^2 e^{-2k R \pi}}{\Lambda_{IR}^2} (\alpha_{hh} - 2\alpha_{hf} + \alpha_{ff})$$

$$\alpha S = \frac{8\sin^2 \theta_W \cos^2 \theta_W m_Z^2 y_1 k^2 e^{-2k R \pi}}{\Lambda_{IR}^4} (\alpha_{ff} - \alpha_{hf})$$

(22)

where $y_1$ denotes the position of the IR brane and $\alpha_{ij}$ are parameters involving the bulk propagators of the bulk gauge fields with $(++)$ boundary conditions, where $+$ denotes Neumann boundary conditions.
condition. They are given as

\[
\alpha_{hh} = \int e^{2A(y)} \left( \Omega_h - \frac{y}{y_1} \right)^2 \\
\alpha_{hf} = \int e^{2A(y)} \left( \Omega_h - \frac{y}{y_1} \right) \left( \Omega_f - \frac{y}{y_1} \right) \\
\alpha_{ff} = \int e^{2A(y)} \left( \Omega_f - \frac{y}{y_1} \right)^2
\]

where \( \Omega_{f,h}(y) = \frac{1}{y_1} \int_0^y dy f_{f,h}^2(y) \) and the profiles \( f's \) are given by Eq. (18) For the case where the fermions are localized on the UV brane \( \Omega_f = 1 \). These co-efficients are a function of the localization of the zero mode of the fermionic and the Higgs field. For a fixed KK scale, the co-efficients increase as the fields move closer to the IR brane due to larger overlap of the zero mode with the KK modes.

For the oblique \( T \) parameters, the co-efficient \( \alpha_{hh} \), also contributes in addition to \( \alpha_{hf,ff} \). Owing to the localization of the Higgs very close to the IR brane, \( \alpha_{hh} \) will be enhanced as compared to \( \alpha_{hf,ff} \), which is smaller as the fermions are closer to the UV brane. As a result in this scenario the contributions to the \( T \) parameter is large. In this case the oblique observables primarily depend on two parameters:

a) The localization parameter \( b \) for the bulk Higgs field.

b) First KK scale of the gauge boson.

To extract the parameter space of these two parameters which are consistent with the constraints on the \( S \) and \( T \) parameters, a scan is performed over the following ranges

\[
b \equiv [2, 5] \quad \Lambda_{IR} \equiv [1250, 10000]
\]

Fig. 3 shows the 3\( \sigma \) region in the \( b - \Lambda_{IR} \) plane. The first KK mass of the gauge boson is related to the IR scale as \( m_{KK}^{(1)} \sim 2.44 \Lambda_{IR} \). We see \( \Lambda_{IR} \) is lowered as \( b \) approaches 2 corresponding to
shifting of the Higgs away from the IR brane. However $b \geq 2$ must be maintained for the model to serve as solution to the hierarchy problem [35, 51]. Table III gives the fit values when all input observables along with $b$ and $\Lambda_{IR}$ are varied simultaneously to minimize the $\chi^2$ in Eq.(9). From the plot in Fig. 3 we find that the lowest value of $\Lambda_{IR}$ possible is around 3.4 TeV corresponding a first KK mass for the gauge boson to be around 8 TeV. The plot is highly concentrated around $b = 2$ since the coupling of the SM fields to the KK states is small as compared to higher values of $b$. This point corresponds to the case where the mass of the first KK gauge boson is minimum.

| Input observables | $m_Z$  | $\alpha(m_Z)$ | $G_F$ | $\alpha_s(m_Z)$ | $\alpha_s(m_Z)$ | $m_t(m_t)$ | $m_H$ |
|-------------------|--------|---------------|-------|----------------|----------------|------------|-------|
|                   | 91.1813 | 7.81611 x 10^{-3} | 1.1663784 x 10^{-5} | 0.119101 | 173.05 | 126.3 |

| Output observables | $m_W$ | $\Gamma_Z$ | $R_\mu$ | $R_{\tau\mu}$ | $R_{c\tau}$ | $R_e$ | $\sin^2 \theta_e$ | $\sin^2 \theta_b$ | $\sin^2 \theta_c$ |
|-------------------|--------|------------|---------|-------------|------------|------|----------------|----------------|----------------|
|                   | 80.411 | 2.4983     | 20.7473 | 20.7941     | 0.1712     |      | 0.2313         | 0.2327         | 0.2312         |

| Model Parameters | $b$ | $m_{kk}^1$ | $8.3$ TeV |
|------------------|-----|-----------|-----------|
|                  | 2.00 |           |           |

TABLE III: Fit values for the input and output observables with a bulk Higgs. Input observables are free in the fit and are varied within their 1 $\sigma$ allowed experimental deviation. $b$=2.00 and $m_{kk}^1$ = 8.3 TeV is obtained for the fit.

We finally note that for the brane localized case, a minimum KK mass of 13.6 TeV is required for the model to be consistent with the data. The fit values for the input and the output observables are given in Table IV.

It is to be noted that KK scales in excess on 20 TeV is required when constraints from FCNC like $\mu \rightarrow e\gamma$ are taken into account. Implementation of bulk flavour symmetries with the imposition of the Minimal Flavour Violation (MFV) ansatz helps in substantially reducing the KK mass to around $\sim 3$ TeV [10, 21, 53, 54].

Fine Tuning: It is well known that the requirement of massless zero mode for bulk scalar field requires the bulk mass $m_{bulk} = ak$ and the brane mass $m_{brane} = bk$ be related by the following
\[
\begin{array}{|c|c|c|}
\hline
\text{Input observables} & m_Z & 91.1856 \\
\alpha(m_Z) & 7.816649 \times 10^{-3} \\
\alpha_s(m_Z) & 0.118657 \\
\hline
\text{Output observables} & m_W & 80.402 \\
\sigma_{\text{had}} & 41.477 \\
R_\mu & 20.744 \\
R_b & 0.2158 \\
\sin^2 \theta_c & 0.2314 \\
\sin^2 \theta_e & 0.2313 \\
A^b_{FB} & 0.1036 \\
A_b & 0.9347 \\
\hline
\end{array}
\]

TABLE IV: Fit values for the input and output observables. Input observables are free in the fit and are varied within their 1 \(\sigma\) allowed experimental deviation. \(m_{kk} = 13.6 \text{ TeV}\) is obtained for the fit.

relation\(^6\):

\[m_{\text{brane}} = 2k + \sqrt{4k^2 + m_{\text{bulk}}^2}\] (25)

Any misalignment between the brane and the bulk masses will result in a non-zero mass for the zero mode. In a realistic model with electroweak symmetry breaking, the Higgs boson is massive and is related to the misalignment as\(^3\)\(^4\)\(^5\)\(^5\)\(^5\)\(^5\):

\[m^2_H = 4(b - 1)(m_{\text{brane}} - m'_{\text{brane}}) \frac{\Lambda^2_{IR}}{k}\] (26)

where \(m'_{\text{brane}} = bk\) is value of the brane mass when the zero mode is massless. As given in Table\(^\text{III}\)\(^\text{III}\), \(\Lambda_{IR} = 2.41 \text{ TeV}\) for the model to consistent with the electroweak precision data. As a result for \(b = 2\) a cancellation upto the fourth decimal between \(m'_{\text{brane}}\) and \(m_{\text{brane}}\) is required to fit a Higgs mass of 126 GeV. The level of tuning increases as the Higgs field is pushed further near the IR brane corresponding to an increase in \(b\). Due to a direct dependence on the brane mass parameter \(b\)\(^\text{34}\)\(^\text{35}\)\(^\text{56}\), it is fair to expect that the Higgs boson mass is best fit by \(b \sim 2\). In the dual theory this corresponds to the Higgs field is a partial composite state with relevant coupling between the source and the CFT (conformal field theory) sectors. As \(b\) increases, this coupling increases and the state is fully composite of the CFT, thus recovering the original RS setup.

\(^6\) The relation \(m_{\text{brane}} = 2k - \sqrt{4k^2 + m_{\text{bulk}}^2}\) is relevant for a Higgs field localized away from the IR brane and is not relevant to the discussion here.
V. DEFORMED RS MODEL

The expression for the $S$ and $T$ parameters in Eq. (22) can be re-expressed as

$$\alpha S = 8\cos^2\theta_W \sin^2\theta_W \frac{m_Z^2}{A_{IR}^2} \frac{1}{Z} I$$

$$\alpha T = \sin^2\theta_W \frac{m_Z^2}{A_{IR}^2} \frac{ky_1}{Z^2} I$$

(27)

where $y_1$ is the position of the IR brane and the dimensionless integral $I$ and $Z$ are defined as

$$I = k \int_{0}^{y_1} \left( (k (y_1 - y))^2 \right) e^{2A(y)-2A(y_1)} dy$$

$$Z = k \int_{0}^{y_1} dy \frac{h^2(y)}{h^2(y_1)} e^{-2A(y)+2A(y_1)}$$

(28)

$h(y)$ is the profile of the vacuum expectation value and is given as

$$h(y) = h(y_1)e^{bk(y-y_1)}$$

(29)

We find the $T$ parameter is enhanced by the volume factor in addition to being suppressed by two powers of $Z$. In RS models where $A(y) = ky$, $Z = 0.5$ for $b = 2$ and becomes smaller as the Higgs field approaches the IR brane ($b \to \infty$). This results in the enhancement of the $T$ parameter leading to stringent constraints in the KK scale. As a result, the authors in [34–36] considered an alternative solution by considering modification of the line element in Eq. (15) where $A(y)$ is now given as

$$A(y) = ky - \frac{1}{\nu^2} \log(1 - \frac{y}{y_s})$$

(30)

Note that $\nu \to \infty$ results in RS limit. A consequence of this metric is that the singularity at the IR brane is shifted outside the patch between IR and UV brane at $y_s = y_1 + \Delta$. $\Delta$ is the distance of the singularity from the IR brane. For the case where the hierarchy problem is solved i.e. $A(y_1) \sim 36$, the position of the IR brane in the bulk $y_1$ is a function of $\nu, \Delta$. Smaller $\nu$ will in general result in a smaller volume factor $y_1$ and helps in ameliorating the constraints on the KK mass from the $T$ parameter. Additionally as noted in [34–36] this setup results in large values of $Z$ for certain choices of parameters $\nu, \Delta, b$ which help in reducing the KK scales so as to be within the reach of LHC.

As before we perform an analysis to determine the parameter space of the $b - \Lambda_{IR}$ plane. We choose two sets of $(\nu, \Delta)$ as follows:

a) $\nu = 0.8$ and $\Delta = 1$. This corresponds to $y_1 = 30.60/k$ so that $A(y_1) \equiv 36$

b) $\nu = 1$ and $\Delta = 0.1$. This corresponds to $y_1 = 30.28/k$ so that $A(y_1) \equiv 36$
FIG. 4: Allowed parameter space in the $b - \Lambda_{IR}$ plane for deformed metric. $\Lambda_{IR}$ is in GeV. The left panel corresponds to $\nu = 0.8$ and $\Delta = 1$ while the right panel corresponds to $\nu = 1$ and $\Delta = 0.1$.

For the deformed metric, the $\Lambda_{IR}$ is related to the first $m_{kk}$ scale from the following relation\textsuperscript{[34, 35]}

$$m_{kk}^1 \sim j_{0,1} \frac{A'(y_1)}{k} \Lambda_{IR}$$

(31)

where $j_{0,1}$ is the first zero of Bessel function $J_0(x)$. We scan the $b$ parameter from 2 to 5 and $\Lambda_{IR}$ is scanned from 50 to 587 GeV for case a) while it is scanned from 50 to 120 GeV for case b). The upper limit on $\Lambda_{IR}$ corresponds to a KK mass of $\sim 3$ TeV. From Fig.(4), for the left panel, a lowest value of $\Lambda_{IR} = 472.6$ GeV is obtained for $b=2$ which corresponds to a first KK mass of about 2.3 TeV. While for case b) depicted in the right panel of Fig.(4), a lowest value of $\Lambda_{IR} = 71.10$ GeV is obtained again for $b = 2$. This corresponds to a first KK mass of about 1.7 TeV for the gauge boson. Thus we see that for certain choices of the metric depending on the values ($\nu, \Delta$), the first KK mass of the gauge boson can be below 2 TeV. The fit values for the input and output observables are given in Table\textsuperscript{V}. Case b) offers an advantage over Case a) in terms of being a less fine tuned model since the $\Lambda_{IR}$ for the fit is small. The analysis can be repeated for different values of $\nu$ and $\Delta$. For our analysis we fit the top quark mass by $c_{Q_3} \sim 0.475$ and $c_t \sim -1$ with a choice of $O(1)$ Yukawa $\sim 4$.

The localization of the top doublet relatively near the UV brane is to minimize the correction to the $Zbb$ vertex.

Fine Tuning: Due to the deformation in the metric, the Higgs mass in Eq.(26) can be generalized to \textsuperscript{[34, 35, 55]}

$$m_H^2 = \frac{2}{Z} (m_{brane} - m'_{brane}) \frac{A_{IR}^2}{k}$$

where $$Z = \int_0^{y_1} \frac{h^2(y)}{h^2(y_1)} e^{-2A(y) + A(y_1)}$$

(32)
\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|}
\hline
Input observables & $m_Z$ & 91.1813 & $G_F$ \hline
\alpha(m_Z) & $7.81663 \times 10^{-3}$ & & $1.1663785 \times 10^{-5}$ \hline
\alpha_s(m_Z) & 0.119118 & & $m_H$ \hline
Output observables & $m_W$ & 80.419 & $\Gamma_Z$ \hline
& $\sigma_{had}$ & 41.486 & $R_e$ \hline
& $R_\mu$ & 20.7382 & $R_{\tau\mu}$ \hline
& $R_b$ & 0.215 & $R_c$ \hline
& $\sin^2\theta_c$ & 0.2314 & $\sin^2\theta_b$ \hline
& $\sin^2\theta_e$ & 0.2313 & $A_{FB}^c$ \hline
& $A_{FB}^b$ & 0.1032 & $A_{FB}^e$ \hline
& $A_b$ & 0.9346 & $A_c$ \hline
Model Parameters & $b$ & 2.00 & $m_{1k}^1$ \hline
& & & 2.3 TeV \hline
\end{tabular}
\caption{Fit values for the input and output observables for the deformed RS case. $b = 2.0006$ and $m_{1k}^1 = 2.3$ TeV is obtained for the fit. $\nu = 0.8$ and $\Delta = \frac{1}{2}$ are chosen for the fit.}
\end{table}

For the normal RS case $A(y) = ky$ and $Z = \frac{1}{2(a-1)}$, thus reducing to Eq. (26). In comparison to RS where $Z < 1$, certain choices of $\nu$ and $\delta$ result in $Z > 1$ which not only lowers the contribution to the $T$ parameter in Eq. (27) but also helps in reducing the fine tuning to obtain the Higgs mass. For instance for the parameters in Table. V, $Z = 2.6$, the tuning reduces to 0.018.

\section{VI. CUSTODIAL RS}

The custodial Randall Sundrum set up \cite{30} contains an enlarged bulk gauged symmetry given by

\begin{equation}
SU(2)_L \times SU(2)_R \times U(1)_X
\end{equation}

which restores the custodial symmetry in the RS setup for the Higgs potential. The corresponding gauge bosons are denoted by $W^{1,2,3}_{L\mu}$, $W^{1,2,3}_{R\mu}$, $X_\mu$ with $g_{5L,5R,5X}$ denoting the corresponding five dimensional gauge couplings. In updating the electroweak constraints in this setup we follow the notation of \cite{31}.

The bulk symmetry is broken down to the Standard Model by considering the following boundary conditions for the gauge fields

\begin{equation}
W^{1,2,3}_{L\mu}(++) B_\mu(++) W^{1,2}_{R\mu}(-+) Z'_\mu(-+)
\end{equation}
with \(+(-)\) denoting Neumann(Dirichlet) boundary conditions as before. The gauge fields $B_\mu$ and $Z'_\mu$ are defined as

$$B_\mu = \frac{g_5 X W^2_{R\mu} + g_5 X R_\mu}{\sqrt{g_5^2 + g_5^2}}; \quad Z'_\mu = \frac{g_5 R W^3_{R\mu} - g_5 X X_\mu}{\sqrt{g_5^2 + g_5^2}}$$ (35)

The the $W^1_{L\mu}$ and $B_\mu$ possess zero modes corresponding to the SM $SU(2)_L$ and the $U(1)_Y$ gauge boson respectively. The hypercharge coupling is given by $Y = 2(T^3_R + Q_X)$. After electroweak symmetry breaking, the electromagnetic charge is given by $Q_{em} = T^3_R + T^3_L + Q_X$. On the other hand, owing to the mixed boundary conditions of $W^1_{L\mu}$ and $Z'_\mu$, they do not possess a zero mode.

The presence of new gauge bosons induce additional corrections to the $T$ parameter but the $S$ parameter remains unchanged. It is given as [9, 13, 35, 57, 58]

$$\alpha T = \frac{\sin^2 \theta_W m_{Z'}^2 y_1 k e^{-2 k R \pi}}{\Lambda^2_{IR}} \left(\alpha_{hh} - \alpha'_{hh} - 2 \alpha_{hf} + \alpha_{ff}\right)$$ (36)

where the $\alpha'_{hh}$ is the bulk propagator for the bosons with \((-+)\) boundary conditions and is given as

$$\alpha'_{hh} = \int_0^{y_1} e^{2 A(y)} \left(1 - \Omega_h\right)^2$$ (37)

The contribution to the $T$ parameter due to the KK states of the SM as well as the new gauge bosons are very similar in magnitude. Recalling that the dominant contribution to the $T$ parameter is due to $\alpha_{hh}$, the presence of $\alpha'_{hh}$ nearly cancels this contribution, thus significantly lowering the constraint on the first KK scale of the gauge boson. As a result, the dominant constraint to the KK scale is due to the $S$ parameter. To see the effects of a negligible $T$ parameter on the fits we assume UV localized fermions to begin with. To obtain the plot for $b - \Lambda_{IR}$ parameter space in the presence of custodial symmetry, we first evaluate the constraints at tree level. From the $S - T$ plot in Fig.1, the region around $T \sim 0$ would also necessitate the $S$ parameter to be small thereby pushing the KK scale up. Indeed, as is noted in the left panel of Fig.5, a lowest value of $\Lambda_{IR} = 1659$ GeV is obtained which translates into a lowest KK mass of around 4 TeV. While this case does better than the normal bulk RS scenario the first KK mass is still out of reach of LHC.

The assumption of UV localized fermions is not sufficient to fit the top quark mass as it would result in large $O(1)$ Yukawa coupling. As a result the zero mode top doublet and the singlet must be moved closer to the IR brane ($c < 0.5$) to increase overlap with the Higgs. This results in the shift of the coupling of $b_L$ to the Z boson. The relative shift to the $b_L$ is given as [9, 31, 59]

$$\frac{\delta g_{b_L}}{g_{b_L}} = - v^2 \left(\frac{g_5^2 T^3_L - g_5^2 T^3_R}{g_5^2 T^3_L - (g')^2 Y} \right) \left(\alpha_{hf} - \alpha'_{hf} - \alpha'_{hf} \right)$$ (38)
The first term involving to $\alpha_{h}^{l}$ will be significant in this case as the third generation doublet is localized closer to the IR brane to fit the top quark mass. In the second term however the presence of $\alpha_{h}^{l}$ and $\alpha_{h}^{r}$ with a relative minus sign softens the impact of localization of the third generation on $\delta_{\text{GBL}}$. The current constraints on the corrections to the $Zb_{L}b_{L}$ coupling pushes the limit obtained in the left panel of Fig. 5 beyond 5 TeV. However it was observed in [60], that the dominant contribution due to the first term in Eq.(38), can be removed by assuming $T_{L}^{3} = T_{R}^{3}$ and $g_{L} = g_{R}$ thus significantly softening the constraints on the KK mass from corrections to the $Zb\bar{b}$
vertex. This implies that the left handed bottom must belong to bi-doublets of $SU(2)_L \times SU(2)_R$. The bi-doublets induce large negative contributions to the $T$ parameter in most regions of the parameter space. The contribution is a function of $c_{Q,t}$, which are the localization parameters for the bi-doublet and the singlet $t_R$. It was noted in \cite{31} that the negative contribution decreases as the doublet and/or singlet are localized away from the IR brane. However, the top quarks mass as a function of $c_{Q,t}$ is given as

$$m_{\text{top}} = \frac{Y_{u,3} \sqrt{v}}{\sqrt{2}(\pi R)^{3/2}} \int_0^{\pi R} f_0^{(0)}(b, y) f_{1/2}^{(0)}(c_{Q,t}, y) f_{1/2}^{(0)}(c_t, y) \frac{1}{\sqrt{2}}$$

where $Y_{u,3}$ is dimensionless $O(1)$ parameter. Choosing the bi-doublet and the singlet to be localized away from the IR brane, will result in the choice of large $Y_{u,3} > 10$ to fit the top quark mass and is not feasible. As result the combination of $c_{Q,t}$ which induces a positive contribution to the $T$ parameter is when $c_Q \leq 0$ and $c_t \sim [0, 0.4, 0.5]$. As shown by \cite{31}, this choice of $c$ parameters not only fits the top quark mass but also gives non-negative contribution to the $T$ parameter.

A non-zero positive contribution to the $T$ parameter would correspond to moving vertically up in the $S - T$ plane in Fig.\[1\]. Owing the tilted orientation of the ellipse, it makes it possible to accommodate larger values of $S$ thereby helping in reducing the lower bound on the first KK mass. The right panel of Fig.\[5\] corresponds to the case where the loop level contributions to the $T$ parameter \cite{31} have been turned on. In the figure we have assumed $\Delta T_{\text{loop}} \leq 0.1$. We find that a minimum of $\Delta T_{\text{loop}} \sim 0.06$ for $b = 2$ is required to lower the scale of the first KK gauge boson below 3 TeV. As $b$ increases corresponding to Higgs moving further towards the IR brane, one can expect minimum $\Delta T_{\text{loop}}$ required to keep the KK scale below 3 TeV to increase. The fit values for the input and output observables is given in Table.\[VI\].

Fine tuning: In this case too we find that the best fit to the precision data is when the brane mass parameter $b = 2$. The KK mass is lowered to 3 TeV thus reducing the fine tuning by an order of magnitude. As a result the cancellation between $m_{\text{brane}}$ and $m_{\text{brane}}'$ is of the order of 0.002

VII. CONCLUSIONS

There is a perception that in order that the Randall-Sundrum model successfully address the gauge-hierarchy problem the Higgs ought to be localised on the IR brane. It has been noted earlier \cite{51, 55} that this is not the case and our results bear this out. In fact, we find that even if we move the Higgs field off the IR brane, a solution to the gauge-hierarchy problem is obtained as long as we have $b \geq 2$. Further, electroweak fits and fine tuning argument seem to be preferring a
In the dual CFT terminology, the Higgs field is a partially composite state \([61, 62]\). This has to do with the exponential form of the scalar profiles which get pushed close to the IR brane for values of \(b\) greater than 2, so that from the point of view of the gauge-hierarchy the Higgs is essentially IR-localized. However, such a bulk Higgs differs from the brane-localised Higgs in the freedom that it offers in exploring the parameter space of the model when confronted with electroweak precision constraints.

A few remarks about the collider implications of bulk RS models are in order. Generically, in these models the gauge boson KK modes provide the most interesting signals and the KK gluon is, of these, the most important \([63–65]\). The production cross-section of the KK gauge boson modes is very small partly because of the couplings of these modes to the SM particles but also because of the strong constraints on the masses of the KK modes coming from electroweak and flavour constraints. The cross-sections for other KK modes, like those of the fermions, are even smaller than that of the gauge boson KK modes (except in some versions of the RS model where the Higgs is treated as a pseudo-Nambu Goldstone boson). The collider tests of the bulk RS models are therefore difficult and several studies which propose probing alternative production channels have been presented \([66, 67]\) but the range of masses probed by these processes is just marginally larger than that allowed by precision constraints. In view of this, our results of the global fit for the deformed metric case are very encouraging. Unlike the custodial symmetry case for which the global fits yield a bound on the mass of the first KK mode of about 2.9 TeV, one gets a lower bound of around 2.3 TeV at 3 \(\sigma\) for the case of the deformed metric for \(\nu = 0.8\) and \(\Delta = 1\). This bound reduces to about 1.7 TeV \(\nu = 1\) and \(\Delta = 0.1\) Table\([VII]\) gives a summary of the results obtained. The deformed metric model then is testable at the LHC at a statistically significant level and a more detailed study of the collider implications of this model is called for.

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