Radion-Higgs Mixing in 2HDMs

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ABSTRACT: We study the custodial Randall-Sundrum model with two Higgs doublets localized in the TeV brane. The scalar potential is CP-conserving and has a softly broken $\mathbb{Z}_2$ symmetry. In the presence of a curvature-scalar mixing term $\xi_{ab} \mathcal{R} \Phi_a^\dagger \Phi_b$ the radion that stabilizes the extra dimension now mixes with the two CP-even neutral scalars $h$ and $H$. A goodness of fit of the LHC data on the properties of the light Higgs is performed on the parameter space of the type-I and type-II models. LHC direct searches for heavy scalars in different decay channels can help distinguish between the radion and a heavy Higgs. The most important signatures involve the ratio of heavy scalar decays into $b$ quark pairs to those into $Z$ pairs, as well as the decay of the scalar (pseudoscalar) into a $Z$ plus a pseudoscalar (scalar).
1 Introduction

The electroweak scale set by the vacuum expectation value (VEV) $v \approx 246$ GeV of the Higgs field is very sensitive to physics at high scales. This sensitivity appears in loop corrections to the Higgs mass and is known as the hierarchy problem. Randall and Sundrum [1] proposed a solution to this puzzle by considering an extra dimensional model with the extra dimension being spatial in nature and compactified into a $S_1/Z_2$ orbifold. In this model there are two 4D manifolds, called “3-branes”, separated by a distance $y_c = \pi r_c$ in the extra dimension where $r_c$ is the ”radius of compactification”. The brane at $y = y_c$ is called the TeV-brane or IR-brane and the brane at $y = 0$ is usually called the UV- or Planck brane. A fine tuning is required between the 5D cosmological constant and the brane tensions in order to achieve a static flat solution which corresponds to a vanishing effective 4D cosmological constant. The solution to Einstein equations gives the 5D metric

$$ds^2 = e^{-2A}g_{\mu\nu}dx^\mu dx^\nu - dy^2,$$

(1.1)

where $A = k|y|$ is the warp factor and $k$ is the $AdS$ curvature scale. This solution corresponds to a slice of $AdS_5$ space between the two branes. The result of their seminal work can
explain the hierarchy of scales by warping down the Planck scale \(^1\), i.e., \(M_{\text{TeV}} = M_{\text{Pl}} e^{-ky_c}\) to the TeV scale, therefore requiring that 
\(ky_c \approx 37\).

In the original Randall and Sundrum (RS) model, it was assumed that the SM fields live in the visible brane, and only gravity propagates in the bulk of the extra dimension. Shortly after the RS model appeared, several extensions with SM fields propagating in the bulk were found. Bulk gauge bosons were first considered in [2, 3] where the KK mass spectrum as well as their localization were derived. In [4] a complete analysis of the Higgs mechanism for bulk gauge bosons was done for both a bulk and a brane Higgs boson. Fermions in the bulk were introduced in [5]. The whole SM was placed in the bulk in [6]. In [7] bulk fields and supersymmetry were studied. Perhaps the most attractive reason to consider placing fermions in the bulk is that one can explain the mass hierarchy and flavor mixing with parameters of \(O(1)\) [4, 7]. Several works with bulk fermions have appeared [8–15].

One inconvenience in RS models with gauge and matter fields propagating in the bulk are large contributions to electroweak precision observables (EWPO) [16] that push the KK scale far beyond the reach of accelerators. A possible cure can be implemented by imposing a gauge \(SU(2)_L \times SU(2)_R \times U(1)_X\) symmetry in the bulk that is spontaneously broken to provide custodial protection [17] for the \(S\) and \(T\) parameters and this reduces the bound on the KK scale to \(m_{KK} \gtrsim 3\) TeV. This custodial protection also protects the \(Zb\bar{b}\) vertex from large corrections [18].

Scalar fluctuations in the RS metric give rise to a massless scalar field called the radion and in order to fix the size of the extra dimension, the radion needs to have a mass. Goldberger and Wise [19] were the first to consider a model with a scalar field propagating in the bulk of \(AdS_5\) and solved for its profile functions and KK masses. Later they showed in [20] that by choosing appropriate bulk and boundary potentials for the scalar one can generate an effective \(4D\) potential for the radion and therefore were able to stabilize it without requiring fine tuning of the parameters. This became known as the Goldberger Wise (GW) mechanism. However in the GW mechanism they used an ansatz for the metric perturbations that do not satisfy Einstein equations and did not include the radion wavefunction and the backreaction of the metric due to the stabilizing field. In the paper of Csaki et al [21] these effects were included by using the most general ansatz [22] and the superpotential method [23] to solve for the backreaction. Then they considered the small backreaction approximation to solve for the coupled scalar-metric perturbation system and found the radion mass to be \(m_r \sim l\) TeV where \(l\) parametrizes the backreaction and its value is model dependent on the specifics of the scalar VEV profile. Therefore the radion could have a mass of few hundred of GeV and is the lightest particle in the RS model.

Since the radion field emerges as the lightest new state the possibility of being experimentally accessible and its effects on physical phenomena must be investigated. In general, when a scalar is propagating on the brane one can include, by arguments of general covariance, in the four dimensional effective action terms involving the Ricci scalar \(\mathcal{L} \supset M R(g) \phi - \xi R(g) \phi^2\). In this way a scalar can couple non-minimally to gravity. If

\(^1\) We use the value \(M_{\text{Pl}} = 10^{19}\) GeV
the brane scalar is a Higgs boson, gauge invariance implies $M = 0$ and from dimensional analysis one expects $\xi$ to be an $O(1)$ number with unknown sign. Particular attention has been placed on the curvature-Higgs term $R \Phi^\dagger \Phi$ since after expanding out the radion field around its VEV this term induces kinetic mixing between the radion field and the Higgs, therefore requiring a non-unitary transformation to obtain the canonically normalized degrees of freedom. After diagonalization the physical fields become mixtures of the original non-mixed radion and Higgs boson. The phenomenological consequences of a non-zero mixing $\xi \neq 0$ have been studied extensively in the literature [21, 24–35].

The radion interacts with matter via the trace of the energy-momentum tensor and the form of these interactions is very similar to those of the SM Higgs boson but are multiplied by $v/\Lambda$ where $\Lambda \sim O(\text{TeV})$ is a normalization factor. In the case $\xi = 0$, there is no Higgs-radion mixing and the branching ratios of the radion become very similar to those of the SM in the heavy mass region, being dominated by vector bosons while for the low mass region the $gg$ mode is dominant. Due to its large, anomaly induced, coupling to two gluons a radion can be produced through gluon fusion.

The parameter space coming from the curvature-Higgs mixing scenario consists of four parameters, viz., the bare mass terms $m_h$ and $m_r$, the mixing parameter $\xi$ and the normalization scale $\Lambda$. However in some of the above references, the Higgs boson had been discovered [36, 37] and their parameter space is reduced to $(m_r, \xi, \Lambda)$. The $\xi - m_r$ parameter space is very constrained by direct searches for additional scalars at the LHC [34] leaving only small experimentally and theoretically allowed windows for $\Lambda = 3 \text{ TeV}$ and these windows open up as one increases $\Lambda$. The bounds on the parameter $\Lambda$ are dependent the mass the first KK excitation $m_{K^+}$ and the curvature scale $k$ as was shown in [38].

Despite the model differences in the analyses that have appeared on Higgs-radion mixing, the overall conclusion is that there is possibility that the measured Higgs boson could be in fact a mixture of the radion with the Higgs doublet that is consistent with experimental data. However the constraints mentioned in the previous paragraph will be pushed further if a radion signal is not seen in the coming future and it would be interesting to look at possible ways to relax these constraints.

In addition to the RS model, several Beyond the Standard Model (BSM) scenarios have appeared in the last several decades as promising candidates for new physics. One of the most studied and simplest extensions is the Two-Higgs-Doublet Model (2HDM) where a second Higgs doublet is added to the electroweak sector. The 2HDM was primarily motivated by minimal supersymmetry [39] and it has also been studied in the context of axion models [40], the baryon asymmetry of the universe [41, 42], the muon $g - 2$ anomaly [43] and dark matter [44].

In this work we will study how some of the constraints on the minimal Higgs-radion mixing may be relaxed or modified by having curvature scalar couplings of the form $\mathcal{L} \supset \xi_{ab} R(g_{\text{ind}}) \Phi^\dagger_a \Phi_b$ where $a,b = 1, 2$ and a 2HDM is located on the TeV brane. The SM gauge bosons and fermions correspond to the zero modes of 5D bulk fields. In section 2 we introduce some notation and we briefly describe the custodial RS model in section 2.1. A review of the radion field emergence in the RS model together with its interactions with SM particles is done in section 2.2. The 2HDM is presented in subsection 2.3. The two-
Higgs-radion mixing Lagrangian is discussed in section 3. In section 4 the predictions of the model are presented including constraints from LHC data, collider signals and constraints and expectations from heavy Higgs searches. A summary of the interactions of the Higgs eigenstates and the radion with SM particles before mixing is given in appendix A.

2 Model Description

2.1 The Custodial RS Model

In our notation Latin letters denote 5D indices $M = (\mu, 5)$ and Greek letters denote 4D indices $\mu = 0, 1, 2, 3$. The background metric is that of equation (1.1) and we use the convention for the flat space Minkowski tensor $\eta_{\mu\nu} = \text{diag}(+1, -1, -1, -1)$. We will introduce fluctuations around the background later.

We first review the RS model with a custodial [17] gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_X \times P_{LR}$ in the bulk where $P_{LR}$ is a parity symmetry that makes left and right gauge groups equal to each other. The 5D action of the model is given by

$$S = \int d^5x \sqrt{g} \left[ -2M^3R(g) + L_\phi + L_{gauge} + L_{fermion} \right] + \int d^4x \sqrt{g_{\text{ind}}(y = y_c) \left[ L_H + L_Y - V_{IR}(\phi) \right]} - \int d^4x \sqrt{g_{\text{ind}}(y = 0) V_{UV}(\phi)}$$  \hspace{1cm} (2.1)

where the first term corresponds to the Einstein-Hilbert action where $M$ is the 5D Planck scale and $R$ the Ricci scalar. The stabilization mechanism is contained in $L_\phi$ together with its brane potentials $V_{IR}$ and $V_{UV}$. We do not discuss this sector and simply assume that stabilization is performed as in [21]. The gauge sector is given by

$$L_{gauge} = -g^{MO}g^{NP} \left[ \frac{1}{2} \text{Tr} \{ L_{MN} L_{OP} \} + \frac{1}{2} \text{Tr} \{ R_{MN} R_{OP} \} + \frac{1}{4} X_{MN} X_{OP} \right]$$  \hspace{1cm} (2.2)

where $L_{MN}, R_{MN}$ and $X_{MN}$ are the gauge bosons associated with $SU(2)_L, SU(2)_R$ and $U(1)_X$ respectively. In the Planck-brane the symmetry is broken $SU(2)_R \times U(1)_X \rightarrow U(1)_Y$ by appropriate BC’s of the gauge fields to generate the SM gauge group. This BC’s are given by [45]

$$\partial_5 L^a_\mu(x, 0) = 0, \quad a = 1, 2, 3,$$

$$R^a_i(x, 0) = \quad i = 1, 2$$

$$g_X \partial_5 R^3_\mu(x, 0) + g_R \partial_5 X_\mu(x, 0) = 0$$

$$-g_R R^3_\mu(x, 0) + g_X X_\mu(x, 0) = 0$$  \hspace{1cm} (2.3)

where $g_L, g_R$ and $g_X$ are the 5D gauge couplings associated with the gauge fields $L^a_\mu, R^a_\mu$ and $X_\mu$ respectively. The SM gauge bosons $W^\pm, Z$ and the photon are embedded into the 5D gauge bosons. Calculation of the spectrum and profiles was performed in Ref. [45, 46] with different KK basis.

Boundary mass terms are generated by the Higgs VEV’s

$$L_{mass} = \frac{v_1^2 + v_2^2}{8} (g_L L^a_\mu - g_R R^a_\mu)^2 \delta(y - y_c),$$  \hspace{1cm} (2.4)
therefore in the TeV brane the gauge symmetry is spontaneously broken down by the Higgs VEV’s to the diagonal group, i.e. $SU(2)_L \times SU(2)_R \rightarrow SU(2)_V$ so that $SU(2)_V$ generates custodial protection for the $T$ parameter. The extra parity symmetry $P_{LR} : SU(2)_L \leftrightarrow SU(2)_R$ was introduced to protect the $Zb_L\bar{b}_L$ vertex from non universal corrections [18].

In the fermion sector all three generations are embedded in the same representation of the gauge group with the following transformation properties [46, 47]

\[
Q_L \sim (2, 2)_{2/3},
\]
\[
u_R \sim (1, 1)_{2/3},
\]
\[
d_R \sim (1, 3)_{2/3} \oplus (3, 1)_{2/3},
\]

and this choice guarantees custodial protection for the $Zbb$ coupling and for flavor violating couplings $Zd_i^L d_j^L$ as well. Using appropriate BC one can ensure that only the SM quarks appear in the low energy theory.

2.2 The Radion Field

For the background metric solution in the RS model, given by equation (1.1), any value of the radius dimension $y_c$ is equally acceptable. Therefore a mechanism is needed to fix the value $y_c \sim 37/k$ so that the EW hierarchy is explained and this must be accomplished without severe fine tuning of parameters. Here we simply assume that a GW bulk scalar is responsible for the stabilization and that the bulk and brane potentials are chosen by applying the method of the superpotential of Ref.[23]. This method has the advantage of reducing the coupled non-linear second order Einstein equations to simple ordinary differential equations for a simple choice of superpotential. The backreaction of the background metric due to the scalar can be solved directly using this method.

After the extra dimension is stabilized the radion field arises from the scalar fluctuations of the metric given by the general ansatz [21, 22]

\[
ds^2 = e^{-2A - 2F(x, y)} \eta_{\mu \nu} dx^\mu dx^\nu - (1 + G(x, y))^2 dy^2,
\]

and since the background VEV for the bulk scalar also depends on the extra dimension one also has to include the fluctuations in the GW scalar namely: $\phi(x, y) = \phi_0(y) + \varphi(x, y)$ where $\phi_0$ is the background VEV and $\varphi$ denotes the fluctuation. By evaluating the linearized Einstein equations one is able to derive $G = 2F$. To solve the system one linearizes the Einstein and scalar field equations to obtain coupled relations for $\varphi$ and $F$. In particular, by integrating the $(\mu, 5)$ component of the linearized Einstein equations $\delta R_{\mu 5} = \kappa^2 \delta T_{\mu 5}$ with $\kappa^2 = 1/2M^3$, one obtains

\[
\phi_0' \varphi = \frac{3}{\kappa^2} (F'' - 2A'F')
\]

where the prime indicates $d/dy$ and this equation implies that the fluctuations $\varphi$ and $F$ will have the same KK eigenstates but with different profiles. Using the Einstein equations together with (2.9) a single differential equation in the bulk for $F$ can be obtained [21]:

\[
F'' - 2A' F' - 4A'' F - 2 \frac{\phi''_0}{\phi_0} F' + 4A' \frac{\phi'_0}{\phi_0} F = e^{2A}\Box F
\]
supplemented by the boundary conditions

\[(F' - 2A'F)|_{y=0,y_c} = 0, \]  

(2.11)

where the boundary conditions are simplified in the limit of stiff boundary potentials of the bulk stabilizer $\partial^2 V_i / \partial \phi^2 \gg 1$ implying $\varphi|_{y=y_i} = 0$. In the system there are two integration constants and one mass eigenvalue $\Box F_n(x,y) = -m_n^2 F_n(x,y)$. One integration constant corresponds to an overall normalization while the other constant and the mass eigenvalue are determined by the boundary conditions. In Ref [21] this differential equation was solved in a perturbative approach in the limit of small backreaction of the metric due to the stabilizing scalar, and it was found to zero-order in the backreaction that the KK zero-mode can be approximated by

\[F_0(x,y) \approx e^{2ky_c} R(x) + O(l^2),\]  

(2.12)

where $R(x)$ is the radion field. Using the boundary conditions the radion mass is

\[m_r \approx lke^{-ky_c}.\]  

(2.13)

where $l^2 = \phi_P^2 / 4M^3$ is the backreaction and $\phi_P$ is the VEV of the bulk stabilizer field on the Planck brane. It should be noted that generically, the radion mass is always proportional to the backreaction independently of the stabilization mechanism. From the expression above, the radion mass is expected to be of $O$(TeV) scale. The canonical normalization of the radion comes from integrating out the extra dimension in the Einstein-Hilbert action

\[M^4 \int dy \sqrt{g} R(g) \geq 6M^3 \frac{k}{k} e^{2ky_c} (\partial_\mu R(x))^2\]  

(2.14)

therefore a canonically normalized radion is obtained by writing

\[R(x) = r(x) \frac{e^{-ky_c}}{\sqrt{6M_P}}.\]  

(2.15)

It is explicitly proved in [21] that the normalization is dominated by the gravitational contribution coming from the Einstein-Hilbert action against that coming from the kinetic term of the bulk stabilizer.

We now proceed to present the radion interactions with the SM fields. The induced metric on the TeV brane is given by

\[\bar{g}_{\mu\nu}^{ind}(x) = e^{-2A(y_c)} e^{-2ke^{2ky_c} R(x)} \eta_{\mu\nu} \equiv e^{-2ky_c} \Omega(r)^2 \eta_{\mu\nu},\]  

(2.16)

where we use $\bar{g}_{MN}$ to denote the metric with scalar perturbations included. After rescaling of the doublets $\Phi_a \rightarrow e^{ky_c} \Phi_a$, the radion couplings to the Higgs sector are obtained from (including the possibility of adding extra scalars in the sum)

\[S_H = \int d^4 x \left[ \sum_{a=1,2} \eta^{\mu\nu} \frac{1}{2} \text{Tr}[(D_\mu M_a)^\dagger D_\nu M_a] \Omega(r)^2 - V(M_1, M_2) \Omega(r)^4 \right],\]  

(2.17)
and all mass terms are redshifted accordingly. Expanding to linear order in the radion field \(\Omega(r) \approx 1 - r^2\), with \(\gamma \equiv v/\Lambda\) and \(\Lambda \equiv \sqrt{6M_P} e^{-ky_c}\), a straightforward calculation yields the coupling of the radion with the trace of the energy-momentum tensor
\[
\frac{\gamma}{v} \gamma T_\mu{}^\mu \supset - \sum \frac{\gamma}{v} r \left[ (\partial_\mu \phi)^2 - 2m_\phi^2 \phi^2 \right],
\]
with the sum performed over all physical scalars.

The couplings to the EW gauge sector are obtained from the kinetic terms of the Higgs doublets expanding to linear order in the perturbations
\[
S_H \supset - \int d^4x \frac{\gamma}{v} r(x) \gamma^{\mu\nu} \left\{ 2m_W^2 W^{(0)+}_\mu(x) W^{(0)-}_\nu(x) + m_Z^2 Z^{(0)}_\mu(x) Z^{(0)}_\nu(x) + \ldots \right\},
\]
where the dots represent higher KK excitations. In addition to the boundary terms there are tree level couplings of the radion coming from the kinetic term of the bulk gauge bosons \([9]\)
\[
S_{gauge} \supset - \int d^4x \frac{\gamma}{v} r(x) \left\{ \frac{1}{k y_c} \frac{1}{4} \eta^{\mu\nu} \eta^{\alpha\beta} V^{(0)}_\alpha(x) V^{(0)}_{\nu\beta}(x) + \frac{m_Y^2}{2k^2} e^{2ky_c} k y_c \eta^{\mu\nu} V^{(0)}_\mu(x) V^{(0)}_{\nu}(x) \right\},
\]
where \(V_{MN} = \partial_M V_N - \partial_N V_M\) is the usual field strength and \(V = \{\sqrt{2}W^\pm, Z, A\}\) and \(m_V = \{m_W, m_Z, 0\}\). The coupling to the field strengths above becomes significant for momentum transfer much larger than the EW scale and the second term constitutes a correction of about 20% to the dominant TeV-boundary coupling. In the case of the photon only the first term is present. A similar expression for gluons should be included.

Overall we can write
\[
\mathcal{L}_{WW,ZZ} r = \frac{\gamma}{v} \left\{ 2m_W^2 \left( 1 - \frac{3m_W^2 k y_c}{\Lambda^2} \right) W^{(0)+}_\mu W^{(0)-}_\nu + m_Z^2 \left( 1 - \frac{3m_Z^2 k y_c}{\Lambda^2} \right) Z^{(0)}_\mu Z^{(0)}_\nu \right\},
\]
For massless gauge bosons we have to include the contributions coming from the localized trace anomaly and from loop triangle diagrams in which the \(W\) gauge boson and fermions in the case of the photon and only fermions in case of the gluons that induce couplings to the radion.

All these contributions can be written as \([9, 24, 30, 34]\)
\[
\mathcal{L}_{tr}^{\gamma\gamma} = - \frac{\gamma}{4v} \left\{ \left( \frac{1}{k y_c} + \frac{\alpha_s b_{QCD}^r}{2\pi} \right) G_{\mu\nu} G^{\mu\nu} + \left( \frac{1}{k y_c} + \frac{\alpha_{EM} b_{EM}^r}{2\pi} \right) F_{\mu\nu} F^{\mu\nu} \right\},
\]
with \(\alpha_s(\alpha_{EM})\) being the strong (electroweak) coupling constant and
\[
b_{QCD}^r = 7 + F_f,
\]
\[
b_{EM}^r = - \frac{11}{3} + \frac{8}{3} F_f - F_W,
\]
\[
F_f = \tau_f (1 + (1 - \tau_f) f(\tau_f)),
\]
\[
F_W = 2 + 3\tau_W + 3\tau_W (2 - \tau_W) f(\tau_W),
\]
\[ f(\tau) = \arcsin\left( \frac{1}{\sqrt{\tau}} \right) \quad \tau \geq 1, \]  
(2.27)

\[ f(\tau) = -\frac{1}{4} \left( \log \frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}} - i\pi \right)^2, \quad \tau < 1, \]  
(2.28)

and \( \tau_i = \left( \frac{2m_i}{m_L} \right)^2 \), \( m_i \) is the mass of the particle going around the loop. An important property of the kinematic functions is their saturation \( F_f \to \frac{2}{3}, F_W \to 7, F_f(\tau) \to 1 \) for \( \tau > 1 \) and \( F_f, W \to 0 \) for \( \tau < 1 \).

In this paper we do not consider the corrections to the couplings coming from excited KK modes of the top and W boson in the loop and simply assume that the above contributions are dominant. However we leave this issue for future work.

Fermions propagating in the bulk are characterized by a bulk mass parameter \( c = m/k \) which specifies their location in the bulk. In addition, the boundary conditions of their profiles at the location of the branes force either the left- or the right-handed zero modes to be zero [5]. Therefore for each SM fermion we need to introduce two different bulk fermions, one with bulk mass parameter \( c_L \) and for which the right-handed zero mode vanishes and the other with a bulk mass parameter \( c_R \) and for which the left-handed zero mode vanishes.

The couplings of the radion to SM fermions can be simplified as [34]

\[ S \supset \int d^4x \sum_{f=u,d,e} \frac{\gamma}{v} r(x) m_f \bar{f} f \times \begin{cases} 1 & \text{Planck} \\ (c_L - c_R) & \text{TeV.} \end{cases} \]  
(2.29)

with the lower option if the zero-mode profile is peaked towards the TeV brane \( c_L < -1/2, c_R > 1/2 \) otherwise the localization is in the Planck brane and the upper option applies. Besides this couplings it seems that the boundary Yukawa couplings will have a direct contribution to the radion couplings to fermions. However, as shown in [9], these contributions get cancelled by induced wave function discontinuities obtained by carefully treating the boundary conditions.

### 2.3 The Two-Higgs Doublet Model

In this work we consider two Higgs doublets living in the visible brane. The most general parametrization for the scalar potential [41, 48] is given by

\[
V(\Phi_1, \Phi_2) = m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - \left( m_{12}^2 \Phi_1^\dagger \Phi_2 + H.c. \right) \\
+ \frac{\lambda_1}{2} (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1)(\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2)(\Phi_2^\dagger \Phi_1) \\
+ \left[ \frac{\lambda_5}{2} (\Phi_2^\dagger \Phi_2)^2 + \lambda_6 (\Phi_1^\dagger \Phi_1)(\Phi_2^\dagger \Phi_2) + \lambda_7 (\Phi_2^\dagger \Phi_2)(\Phi_1^\dagger \Phi_2) + H.c. \right],
\]  
(2.30)

where \( m_{11}^2, m_{22}^2, \) and \( \lambda_{1,2,3,4} \) are real by hermiticity and \( m_{12}^2 \) and \( \lambda_{5,6,7} \) are in general complex. In this expression there are fourteen parameters, however the freedom in the choice of basis can be used to reduce this number down to eleven degrees of freedom that are physical.
To provide custodial protection for the $T$ parameter we promote the Higgs fields to bi-doublets $M_i = (\tilde{\Phi}_i, \Phi_i)$ (with $\tilde{\Phi}_1 = i\sigma^2\Phi_1^*$) of the gauge group $SU(2)_L \times SU(2)_R$ that transform in the representation $(2, 2)_0$ [49]

$$M_i \rightarrow U_LM_iU_R^\dagger, \quad i = 1, 2. \quad (2.31)$$

where

$$U_L \in SU(2)_L, \quad U_R \in SU(2)_R. \quad (2.32)$$

Using the three independent invariant quadratic forms $\text{Tr}[M_1^\dagger M_1]$, $\text{Tr}[M_2^\dagger M_2]$ and $\text{Tr}[M_1^\dagger M_2]^2$ the most general expression that has all possible combinations of traces invariants is given by

$$V(M_1M_2) = \frac{m_1^2}{2} \text{Tr}[M_1^\dagger M_1] + \frac{m_2^2}{2} \text{Tr}[M_2^\dagger M_2] - \bar{m}_1^2 \text{Tr}[M_1^\dagger M_2] + \frac{\lambda_1}{8} \text{Tr}[M_1^\dagger M_1]^2$$

$$+ \bar{\lambda}_2 \text{Tr}[M_2^\dagger M_2]^2 + \bar{\lambda}_3 \text{Tr}[M_1^\dagger M_1]\text{Tr}[M_1^\dagger M_2] + \frac{\lambda_4}{4} \text{Tr}[M_1^\dagger M_2]^2$$

$$+ \bar{\lambda}_5 \text{Tr}[M_1^\dagger M_1]\text{Tr}[M_2^\dagger M_2] + \frac{\lambda_6}{2} \text{Tr}[M_2^\dagger M_2]\text{Tr}[M_1^\dagger M_2] \quad (2.33)$$

where all the parameters are real and the correspondence with the potential of equation (2.30) is

$$\lambda'_4 \equiv \lambda_4, \quad \lambda_5' \equiv \lambda_6, \quad \lambda_7' \equiv \lambda_7. \quad (2.34)$$

Thus by imposing the gauge $SU(2)_L \times SU(2)_R$ symmetry one immediately reduces the number of free parameters in the scalar potential down to nine. Also a custodially protected 2HDM potential is automatically CP conserving.

The kinetic terms for the Higgs bi-doublets are given by

$$\mathcal{L}_H \supseteq \sum_{i=1,2} g^{\mu\nu}_{\text{ind}} \frac{1}{2} \text{Tr}[(D_\mu M_i)^\dagger D_\nu M_i] \quad (2.35)$$

where $g^{\mu\nu}_{\text{ind}}$ is the induced metric on the TeV brane and the covariant derivative is

$$D_\mu M_i = \partial_\mu M_i - ig_L L_\mu M_i + ig_R M_i R_\mu \quad (2.36)$$

and $L_\mu = L_\mu^a T^a_L$ is the gauge boson associated with $SU(2)_L$. Therefore under the custodial gauge symmetry the gauge bosons transform as

$$L_\mu \rightarrow U_L L_\mu U_L^\dagger - \frac{i}{g_L} \partial_\mu U_L U_L^\dagger. \quad (2.37)$$

$$R_\mu \rightarrow U_R R_\mu U_R^\dagger + \frac{i}{g_R} U_R \partial_\mu U_R^\dagger. \quad (2.38)$$

Of course one needs to also include the term corresponding to the gauge group $U(1)_X$ which violates the custodial symmetry.

In conventional 2HDM’s one can avoid the presence of potentially dangerous flavor changing neutral currents (FCNC) by imposing a discrete $Z_2$ symmetry $\Phi_1 \rightarrow \Phi_1, \Phi_2 \rightarrow \Phi_2$.

\[2\] For a basis independent treatment see Ref. [50]
since we still need to canonically normalize the Higgs doublets. The fermion mass in (2.29) is generated either by $\Phi_1$ or $\Phi_2$ since the discrete $Z_2$ symmetry is extended to the fermion sector. This results in four different types of Yukawa interactions [51]. In the type-I model all fermions couple to a single Higgs doublet, usually chosen to be $\Phi_2$. In the type-II model up-type quarks couple to $\Phi_2$ and d-type quarks and leptons couple to $\Phi_1$. In the lepton-specific model all leptons couple to $\Phi_1$ and all quarks couple to $\Phi_2$. Finally in the flipped model up-type quarks and leptons couple to $\Phi_2$ and d-type quarks couple to $\Phi_1$. In general, radion mediated FCNC can be present and this was analyzed in [52]. For simplicity we don’t consider flavor mixing in the bulk mass parameters, i.e., $c_{L,R}^{ij} = c_{L,R}^{ij}$ since we want to achieve minimal flavor violation [53] in the Yukawa sector.

In terms of bi-doublets this symmetry reads

$$M_1 \rightarrow M_1, \quad M_2 \rightarrow -M_2,$$

and implies $\lambda'_0 = \lambda'_0 = 0$ with $n_{i12}^2 \neq 0$ remaining as a soft-violating term. The Higgs doublets can be expressed as

$$\Phi_a = \left( \begin{array}{c} \phi_a^+ \\ \bar{\nu}_a + \rho_a^{ij} \bar{\nu}_b \end{array} \right), \quad a = 1, 2$$

where $\bar{\nu}_a$ are the VEV of the scalars. The VEV’s satisfy the relation $\bar{v}^2 = \bar{v}_1^2 + \bar{v}_2^2$ with $\bar{v}$ the localized Higgs VEV and should not be confused with the SM value $v = \bar{v}e^{-\frac{1}{2}M_{12}} = 246$ GeV since we still need to canonically normalize the Higgs doublets$^3$.

The fields appearing in the expression of the Higgs doublets (2.40) are not the physical scalars. To obtain the physical eigenstates one has to diagonalize the mass matrices that are constructed using equation (2.33) with the appropriate imposed symmetries. For a custodial and $Z_2$ symmetric scalar potential the mass matrix for the CP-odd state and for the charged Higgs fields are equal

$$\begin{pmatrix} \tilde{m}_{11}^2 + \frac{v_1^2 \lambda_1 + v_2^2 \lambda_3}{2} & -\tilde{m}_{12} + \bar{v}_1 \bar{v}_2 \lambda'_4 \\ -\tilde{m}_{12} + \bar{v}_1 \bar{v}_2 \lambda'_4 & \tilde{m}_{22}^2 + \frac{v_3^2 \lambda_2 + v_1^2 \lambda_3}{2} \end{pmatrix} = \begin{pmatrix} \tilde{m}_{12}^2 \bar{v}_1 - \lambda'_1 \bar{v}_2^2 & -\tilde{m}_{12}^2 + \bar{v}_1 \bar{v}_2 \lambda'_4 \\ -\tilde{m}_{12}^2 + \bar{v}_1 \bar{v}_2 \lambda'_4 & \tilde{m}_{22}^2 \bar{v}_1 - \lambda'_1 \bar{v}_2^2 \end{pmatrix}$$

where in the last equality $\tilde{m}_{11}^2$ and $\tilde{m}_{22}^2$ were eliminated using the minimization conditions of the potential. The matrix above has a zero eigenvalue corresponding to the Goldstone bosons $G^0$ and $G^\pm$ and the nonzero mass eigenvalue is given by

$$\tilde{m}_{A}^2 = \tilde{m}_{H^\pm}^2 = \tilde{m}_{12}^2 \frac{\bar{v}^2}{\bar{v}_1 \bar{v}_2} - \lambda'_1 \bar{v}^2.$$

The fact that the CP-odd field mass is degenerate with the charged Higgs bosons is a direct consequence of imposing a custodial symmetry in the scalar potential however this symmetry is not respected by the hypercharge gauge and Yukawa interactions, so we can only expect the masses to be approximately degenerate. The diagonalization of the CP odd fields (as well as the charged scalars) is carried out by the orthogonal transformation

$$\begin{pmatrix} \eta_1 \\ \eta_2 \end{pmatrix} = \begin{pmatrix} c_\beta & -s_\beta \\ s_\beta & c_\beta \end{pmatrix} \begin{pmatrix} G^0 \\ A \end{pmatrix}$$

$^3$We put a bar on mass parameters that are not yet redshifted down to the EW scale.
\[ c_\beta = \cos \beta, \quad s_\beta = \sin \beta \quad \text{and} \quad \tan \beta = v_2/v_1. \quad G^0 \quad \text{is the neutral Goldstone boson and} \quad A \quad \text{is the physical pseudoscalar.} \]

The physical CP even scalars are obtained by the rotation
\[
\begin{pmatrix}
\rho_1 \\
\rho_2
\end{pmatrix} = \begin{pmatrix}
c_\alpha & -s_\alpha \\
s_\alpha & c_\alpha
\end{pmatrix} \begin{pmatrix}
H \\
h
\end{pmatrix}
\]
\[ (2.44) \]

where \( h(H) \) corresponds to the lighter (heavier) scalar.

Notice that there were 7 real parameters in the Higgs potential to start with, namely \({\tilde{m}}^2_{11}, {\tilde{m}}^2_{22}, {\tilde{m}}^2_{12}, \lambda'_1, \lambda'_2, \lambda'_3, \lambda'_4 \). Using the two minimization conditions we can trade \( {\tilde{m}}^2_{11} \) and \( {\tilde{m}}^2_{22} \) for \( v_1 \) and \( v_2 \) and then use the relations \( v^2 = v_1^2 + v_2^2 \) and \( \tan \beta = v_2/v_1 \) to trade \( v_1 \) and \( v_2 \) for \( v \) and \( \beta \).

Finally we can trade the soft breaking parameter and three lambdas for the three scalar masses and \( \alpha \) ending up with the set \( \{\beta, \alpha, m_h, m_H, m_A, \lambda_4\} \) (notice that \( \lambda_4 = \lambda'_4 \)) where we fixed \( v = 246 \) GeV therefore we only have to specify 6 parameters.

| Type-I | \( \xi^u_h \) | \( \xi^d_h \) | \( \xi^u_{H} \) | \( \xi^d_{H} \) | \( \xi^u_{A} \) | \( \xi^d_{A} \) | \( \xi^u_{A} \) | \( \xi^d_{A} \) |
|--------|-------------|-------------|-------------|-------------|-------------|-------------|-------------|-------------|
| Type-II | \( c_\alpha/s_\beta \) | \( -s_\alpha/c_\beta \) | \( c_\alpha/s_\beta \) | \( -s_\alpha/c_\beta \) | \( s_\alpha/s_\beta \) | \( c_\alpha/s_\beta \) | \( -c_\alpha/c_\beta \) | \( -s_\alpha/c_\beta \) |

Table 1: Scalar couplings to pairs of fermions.

The couplings of the scalars with the fermion fields can be written as [51]
\[
\mathcal{L}^f = \sum_{f=u,d,l} \frac{m_f}{v} \left( \xi^f_h \bar{f} h + \xi^f_H \bar{f} H - i \xi^f_A \bar{f} \gamma_5 A \right) + \left\{ \sqrt{2} V_{ud}/v (m_u \xi^u_A P_L + m_d \xi^d_A P_R) d H^+ + \frac{\sqrt{2} m_l \xi^l_A}{v} \bar{u} l \gamma_5 H^+ + h.c. \right\},
\]
\[ (2.45) \]

where the mixing factors are summarized in Table 1.

The couplings of the scalars to a pair of gauge bosons are given by
\[
\mathcal{L}^{W,W,Z} = (h \sin (\beta - \alpha) + H \cos (\beta - \alpha)) \left( \frac{2m^2_W}{v} W^+ W^- + \frac{m^2_Z}{v} Z^+ Z^- \right),
\]
\[ (2.46) \]
\[
\mathcal{L}^{\gamma\gamma} = \sum_{\phi=h,H,A} -\frac{\phi}{4v} \left\{ \frac{\alpha_s}{2\pi} b^{\phi}_{QCQ} G_{\mu\nu}^a G^{a\mu\nu} + \frac{\alpha_{EM}}{2\pi} \frac{b^\phi_{EM}}{f^2} F_{\mu\nu} F^{\mu\nu} \right\},
\]
\[ (2.47) \]

where
\[
b^{\phi}_{QCQ} = \xi^\phi \times \left\{ \begin{array}{l}
F_f, \quad \phi = h, H, \\
F_{(\tau_i)\tau_i}, \quad \phi = A,
\end{array} \right.
\]
\[ (2.48) \]
\[
b^h_{EM} = \frac{8}{3} \xi^h F_f - \sin(\beta - \alpha) F_W + g_h F_H
\]
\[ (2.49) \]
\[
b^H_{EM} = \frac{8}{3} \xi^H F_f - \cos(\beta - \alpha) F_W + g_H F_H
\]
\[ (2.50) \]
The form factor for the charged Higgs in the loop is \[ F_H = -\tau_H (1 - \tau_H f(\tau_H)) \]
and has limiting behaviors \( F_H \to 1/3 \) for \( \tau > 1 \) and \( F_H \to 0 \) for \( \tau < 1 \). The couplings multiplying the form factor are given by \( g_\phi = -\frac{m_W}{g^{\tau H}_H} g_{\phi H} \) with \( g_{\phi H} \) the tree level coupling that arises from the 2HDM potential.

\section{Two Higgs-radion Mixing}

The most general term that will give rise to kinetic mixing between the Higgs doublets and the radion field is given by

\[ L_\xi = \sqrt{g} \xi_{a b} \mathcal{R}(g_{a b}) \frac{1}{2} \text{Tr}[M_a^\dagger M_b] \]

where the indices \( a, b = 1, 2 \) are summed so that we have, in principle, four different mixing parameters. However the assumption of CP invariance forces \( \xi_{12} = \xi_{21} \) and thus the pseudoscalar does not mix with the radion. Evaluation of the Ricci scalar is straightforward and yields the following expression \[ L_\xi = -6 \xi_{a b} \Omega^2 \left[ \Box \ln \Omega + (\nabla \ln \Omega)^2 \right] \frac{1}{2} \text{Tr}[M_a^\dagger M_b] \]

The warp factor disappears after we make the rescaling of the Higgs doublets. Using the expression for the Higgs mass eigenstates (2.44) and expanding to linear order in the fields we can write

\[ L_\xi \supset -6 \left[ -\frac{\gamma}{v} \Box r + \frac{\gamma^2}{v^2} r \Box r \right] \left[ \frac{v^2}{2} K_r + \frac{v}{2} K_h h + \frac{v}{2} K_H H \right], \]

where \( \gamma \equiv v/\Lambda \) and we define the mixing parameters by

\[ K_r = \xi_{11} c^2_\beta + \xi_{22} s^2_\beta + 2 \xi_{21} s_\beta c_\beta, \]

\[ K_h = 2(\xi_{22} s_\beta c_\alpha - \xi_{11} c_\beta s_\alpha) + 2 \xi_{12} \cos(\alpha + \beta), \]

\[ K_H = 2(\xi_{11} c_\beta c_\alpha + \xi_{22} s_\beta s_\alpha) + 2 \xi_{12} \sin(\alpha + \beta). \]

Adding the kinetic and mass terms of each field, the mixing Lagrangian can be expressed as

\[ L = -\frac{1}{2} (1 + 6 \gamma^2 K_r) r \Box r - \frac{1}{2} m_r^2 r^2 + \sum_{\phi = h, H} \left\{ 3 \gamma K_\phi \Box r - \frac{1}{2} \phi (\Box + m_\phi^2) \phi \right\} \]

The kinetic terms can be diagonalized by performing the transformation

\[ r \to r', \quad \phi \to \phi' + \frac{3 \gamma K_\phi}{Z} r' \]

with \( \phi = h, H \) and

\[ Z^2 = 1 + 6 \gamma^2 K_r - 9 \gamma^2 (K_h^2 + K_H^2), \]
is the determinant of the kinetic mixing matrix and therefore should always satisfy $Z^2 > 0$ to avoid the presence of ghosts fields. This condition allows us to impose our first theoretical constraint on the mixing parameters after choosing appropriate values for $\alpha$, $\beta$ and $\gamma$. This transformation induces mixing in the mass terms. The mass matrix obtained can be written as

$$M = \begin{pmatrix}
\omega^2_{rr} & \omega^2_{rh} & \omega^2_{rH} \\
\omega^2_{rh} & m^2_h & 0 \\
\omega^2_{rH} & 0 & m^2_H
\end{pmatrix},$$

(3.10)

where

$$\omega^2_{rr} = \frac{m^2_r}{Z^2} + \frac{9\gamma^2}{Z^2} \left(K^2 m^2_h + K^2 m^2_H\right),$$

(3.11)

$$\omega^2_{r\phi} = \frac{3\gamma}{Z} K_h m^2_{\phi}.$$  

(3.12)

The physical eigenstates are obtained by performing a three dimensional rotation

$$\begin{pmatrix}
r' \\
h' \\
H'
\end{pmatrix} = U \begin{pmatrix}
r_D \\
h_D \\
H_D
\end{pmatrix},$$

(3.13)

The Higgs scalars-radion system is determined by the three mixing parameters of equation (3.1), the two mixing angles of the Higgs sector, the scale $\gamma$ and the three scalar masses, giving a total of nine parameters. However one of the physical masses will be set to the Higgs mass value and only the set $(\xi_{11}, \xi_{12}, \xi_{22}, \alpha, \beta, \gamma, \lambda_r, \lambda_H)$ needs to be specified.

Another important parameter in the study of RS models with bulk gauge bosons is the KK scale defined to be the mass of the first excited state of the gauge bosons. Recall that this parameter is independent of the gauge symmetry and gauge couplings and is universal for all gauge bosons that satisfy the same BCs. In particular, for gauge bosons satisfying Neumann BCs at both branes it is given by [26]

$$m_{KK} = \frac{k}{\sqrt{6M_{Pl}}} \Lambda,$$  

(3.14)

so any bound on the KK scale will directly affect the allowed values of the curvature scale $k$ and $\Lambda$.

In Higgs-radion mixing scenarios there is a particular point in the parameter space called the “conformal point” [24, 34, 35], usually around $\xi = 1/6$ where the conformal symmetry is minimally violated by the Higgs VEV. At this point the tree-level couplings of the radion to the massive fermions and gauge bosons are very suppressed and the $gg$ decay mode dominates even in the large radion mass limit. In this work we do not attempt to calculate a conformal point due to the large number of parameters.

In what follows we will reduce the parameter space by assuming that the diagonal elements of the curvature-scalar mixing matrix are equal to each other, $\xi_{11} = \xi_{22} = \xi_1$ and for simplicity we will refer to the off diagonal as $\xi_{12} = \xi_2$. Relaxing this constraint will not radically alter the numerical results in the following sections.

From now on we will drop the subindex $D$ for the diagonal eigenstates and simply write them as $r, h$ and $H$. Whenever we need to distinguish between the non-diagonal and physical states a clarification will be made.
4 Model Predictions

4.1 Constraints From Current LHC Higgs Data

In the 2HDM the interactions of all the scalars to the SM fields are completely determined by the two mixing angles of the scalar sector $\beta$ and $\alpha$. In addition, the alignment limit is defined to be the limit in which one of the CP-even scalars has exactly the same interactions as the SM Higgs and corresponds to $\cos(\beta - \alpha) = 0$.

In this section we perform an analysis on the effects Higgs-radion mixing has on the 2HDM parameter space, $\cos(\beta - \alpha)$ and $\tan \beta$. We use a chi-square test to fit the model to the data presented in Appendix B and find the region in the 2HDM parameter space allowed by current LHC data on the SM-like Higgs boson, $h$. By definition the chi-square function to be minimized is written as

$$\chi^2 = \sum_i \frac{(R^p_i - R^m_i)^2}{(\sigma_i)^2},$$

where $R^p_i$ is the signal strength predicted by the model, $R^m_i$ is the measured signal strength and $\sigma_i$ is the corresponding standard deviation of the measured signal strength. Asymmetric uncertainties are averaged in quadrature $\sigma = \sqrt{\sigma^2 + \sigma'^2}$. The expected signal strengths are defined as the production cross section times branching ratio of a particular decay channel $ff$ normalized to the standard model prediction, i.e.,

$$R^p_f \equiv \frac{\sigma(pp \rightarrow h)BR(h \rightarrow ff)}{\sigma(pp \rightarrow h_{SM})BR(h_{SM} \rightarrow ff)}.$$  (4.2)

Directly obtaining analytical expressions for the mass eigenstates is challenging therefore we resort to numerical techniques. The analysis was carried out using two benchmarks for the radion vev, $\Lambda = 3, 5 \text{ TeV}$. We generated random values for 2HDM mixing angles, $(\alpha, \beta)$, the curvature scalar couplings $(\xi_1, \xi_2)$ and the scalar mass parameters before radion mixing $(m_h, m_H, m_r)$ amounting to seven degrees of freedom. By imposing the field $h$ has a mass of $125.09 \pm 0.5 \text{ GeV}$ one degree of freedom is removed leaving us with six degrees of freedom in our chi-square analysis. We also imposed a lower bound of 150 GeV on the mass for the scalar mostly aligned with the heavy neutral scalar of the 2HDM. We plot the points allowed by the LHC data in Fig. 1 at a 95% confidence level for the type-I and type-II models.

No significant difference can be observed between the $\Lambda = 3 \text{ TeV}$ and $\Lambda = 5 \text{ TeV}$ plots for each type of model. Therefore it seems that a curvature-scalar mixing has no significant effect on the 2HDM parameter space. One can understand this by looking at the off-diagonal elements of the mass matrix, equation (3.10), which are $3\gamma K_{\phi}/Z \sim 1/1000$ times the diagonal elements. This is a reasonable approximation since we assume natural values for the curvature-scalar mixing parameters, $\xi \sim \mathcal{O}(1)$ and therefore the unitary matrix that diagonalizes (3.10) is nearly diagonal which implies that the couplings of the SM-like Higgs to a pair of gauge bosons and fermions receive very small corrections and are nearly given by the corresponding couplings in the 2HDM, i.e.,

$$g_{hVV} = U_{22} \sin(\beta - \alpha) + U_{32} \cos(\beta - \alpha) + U_{12} \gamma (1 - \frac{3m_r^2k_{yr}}{\Lambda^2}) \approx \sin(\beta - \alpha),$$  (4.3)
Figure 1: The top plots show the allowed regions for the type-I model and the bottom plots show allowed region in the type-II model. The blue (red, black) points shown are used for the $\Lambda = 3(5,100)$ TeV cases. Values of the curvature scalar couplings, $\xi_1, \xi_2$ were allowed to range between $[-4,4]$

\[ g_{hff} = U_{22}^2 \xi_h^f + U_{32}^2 \xi_h^f + U_{12} \gamma (c_L - c_R) \approx \xi_h^f, \]  

(4.4)

where $U_{ij}$ are the elements of the unitary matrix. The general shape of the regions is understood by looking at the behavior of the couplings. In the type-I model $\xi_h^f = \cos \alpha / \sin \beta$ and in the large $\tan \beta$ limit the production cross section is suppressed, allowing the parameter space to grow. For type-II model the coupling to a pair of $b$ quarks is $\xi_h^b = -\sin \alpha / \cos \beta$ and therefore the production cross section is enhanced by the $b$ quark loop squeezing the parameter space.

The region of the curvature-scalar parameter space allowed by the chi-square test is shown in Fig. 2. The region shrinks by reducing the value of $\Lambda$.

5.2 Collider Signals

Let us now consider some predictions of this model accessible to the LHC and how one may distinguish this model from some other multi-Higgs model. One feature of a multi-Higgs model is that the sum of the CP-even scalar couplings to $Z$ bosons in quadrature should total to the square of the SM Higgs coupling to the $Z$ bosons, namely

\[ g_{hSMZZ}^2 \sum_i g_{\phi_i ZZ}^2 = 1. \]  

(4.5)

Due to the bulk couplings of the radion to the bulk gauge bosons we find that the sum of the neutral scalar couplings in quadrature normalized to the $h_{SMZZ}$ coupling gives $1 + \gamma^2 (1 - 3m_Z^2 k y_c / \Lambda^2)$ being bounded from below by 1 and setting it apart from other
multi-Higgs models. However, this deviation from unity may be quite small. For $\Lambda_\phi = 3$ TeV one finds Eq. 4.5 gives 1.0054 and the deviation from unity vanishes in the limit $\Lambda_\phi \to \infty$. It is unlikely that the LHC will be able to measure such a small deviation, but such a measurement may be possible at the future ILC.

Another strategy to distinguish this model from other multi-Higgs models is to measure the ratio of the widths of the heavy scalars to $b\bar{b}$ and $ZZ$ pairs,

$$R_{bb/ZZ}^\Phi \equiv \frac{\Gamma(\Phi \to b\bar{b})}{\Gamma(\Phi \to ZZ)}, \text{ for } \Phi = r, H.$$  

The mass eigenstates, $H$ and $r$ are primarily aligned with the unmixed states. This means that couplings of $H$ to the $Z$ boson and $b$ quark should be dominated by the corresponding expressions in a 2HDM. Then for $H$, $R_{bb/ZZ}^H$ should mostly scale like $\left(\frac{\sin \alpha}{\sin \beta \cos(\beta - \alpha)}\right)^2$ for the type-I model and $\left(\frac{\cos \alpha}{\cos \beta \cos(\beta - \alpha)}\right)^2$ for the type-II model. In either case this ratio becomes quite large in the neighborhood of $\cos(\beta - \alpha) = 0$. For the radion, in the limit that its fully aligned with the unmixed radion, $R_{bb/ZZ}^r \propto \frac{(c_L - c_R)^2}{(1 - \frac{3m_Z^2}{\Lambda^2})^2} \approx (c_L - c_R)^2$. This is typically less than one and thus measurement of this ratio might distinguish $r$ from $H$.

4.3 Constraints From Heavy Higgs searches

The radion interactions with the scalar sector come from the following sources:

1. The quartic interactions in the 2HDM potential

$$V(\Phi_1, \Phi_2) \supset \frac{\lambda_1}{2} (\Phi_1\Phi_1^\dagger)^2 + \frac{\lambda_2}{2} (\Phi_2\Phi_2^\dagger)^2 + \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \frac{\lambda_4}{2} (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1)^2.$$  

2. The coupling of the radion with the trace of the energy momentum tensor

$$\mathcal{L} \supset -\frac{r}{\Lambda} ((\partial_\mu h)^2 - 2m_h^2 h^2 + ...).$$  

---

\textbf{Figure 2:} The parameter space of $\xi_1$ and $\xi_2$ allowed by the chi-square goodness of fit. The blue and red points correspond to $\Lambda = 3$ TeV and $\Lambda = 5$ TeV respectively.
3 The curvature-scalar mixing term $\mathcal{L} = -\xi_{ab} \mathcal{R} \Phi_a^\dagger \Phi_b$, where we expand the Ricci scalar up to second order in $\gamma$:

$$
\mathcal{R} \supset -\frac{\gamma}{v} \Box r + 2\frac{\gamma^2}{v^2} r \Box r + \frac{\gamma^2}{v^2} (\partial_\mu r)^2 + \mathcal{O}(\gamma^3).
$$

4 There is a model dependent contribution coming from the potential of the GW scalar field that one can consider however we will assume this interaction to be small as it is proven in [26] that addition of this extra term doesn’t affect the phenomenology.

5 Non-zero mixing will also induce tree-level interactions of the radion with a gauge field and a scalar, namely $rW^\pm H^\mp$ and $rZA$ coming from a direct expansion of the kinetic term in equation (2.35).

In this model the amount of kinetic mixing between the Higgs field and the radion is parametrized by the parameter $K_h$ of equation (3.5). Similarly the amount of kinetic mixing between the heavy Higgs state and the radion is encoded in the parameter $K_H$ given in equation (3.6). We use the most recent LHC direct searches for a heavy scalar decaying into a pair of SM Higgs bosons [56, 57], into $WW$ bosons [58] and into a pair of $ZZ$ bosons [59] to find bounds on the amount of mixing. The most relevant decay channels, when kinematically accessible, are $\phi_i \rightarrow hh, \phi_j \phi_j, h\phi_j, bb, tt, WW, ZZ, gg, AA, H^+H^-, ZA, W^\pm H^\mp$ with $\phi_i = r, H$. The trilinear interactions coming from the 2HDM potential have a dependence on the pseudoscalar mass $m_A$ and on the quartic coupling of the potential $\lambda_4$. 
Figure 3: Scatter plots of the amount of mixing between the Higgs and the radion, $K_h$, defined in equation (3.5), as function of the radion mass. The black region is theoretically allowed and the points colored yellow, green and red are forbidden by heavy scalar searches in the $WW$, $ZZ$ and $hh$ channels respectively. The benchmark point $\Lambda = 3(5)$TeV was used on the left (right).

We scanned over all the parameters and chose as benchmark values $\Lambda = 3, 5$ TeV, $m_A = 200, 500, 700$ GeV and fixed $\lambda_4 = 0.1$. Changing the value of the quartic coupling does not affect significantly the results. The results are presented as scattered plots in figures 3 and 4 where we show the allowed region in $m_r$-$K_h$ and $m_H$-$K_H$ parameter space. In those figures the background black points correspond to the points that are both theoretically allowed
and that survived the chi-square analysis of the previous subsection while the points colored yellow, green and red correspond to regions that are forbidden by LHC searches of a heavy scalar decaying in the \(WW, ZZ\) and \(HH\) channels respectively. No bounds were found from Higgs resonant production searches in [57]. One can immediately notice that direct searches in the \(WW\) and \(ZZ\) channel forbid mainly the low mass region \(m_r = 200 - 400\) GeV with the bounds from the \(WW\) being weaker than those from the \(ZZ\) channel and no bounds at all from the \(WW\) channel were found for the heavy Higgs. The di-Higgs search channels put constraints mostly in the intermediate mass region \(m_{r/H} = 300 - 800\) GeV.

From the figure we can notice that as the pseudoscalar mass increases the bounds coming from the di-Higgs boson and \(ZZ\) channels become more stringent. This is reasonable since an increase in the pseudoscalar mass corresponds, via the 2HDM potential, to an increase in the trilinear coupling of the radion to a pair of SM Higgs fields and the branching fraction becomes bigger.

The LHC has also searched for a CP-odd Higgs scalar in the processes \(pp \rightarrow H/A \rightarrow ZA/H\) [60–62] where the final state \(Z\) boson decays into two oppositely charged electrons or muons and the scalar, either \(H\) or \(A\), is assumed to decay into a pair of \(b\) quarks. These final states were motivated by the large branching fractions predicted in a 2HDM with type-II Yukawa structure and the benchmark values \(\tan \beta = 0.5-1.5\) and \(\cos(\beta - \alpha) = 0.01\) are used in those references. In those papers, the charged Higgs boson masses were kept equal to the highest mass involved in the benchmark signal, namely \(m_{H\pm}^2 \approx m_H^2\) for \(H \rightarrow ZA\) or \(m_{H\pm}^2 \approx m_A^2\) for \(A \rightarrow ZH\).

Due to the custodial symmetry imposed in the 2HDM potential we can only account for the latter triplet mass degeneracy but we can consider both decay topologies. To the best of our knowledge there has been no search for the signal \(H \rightarrow ZA\) with \(m_{H\pm} \approx m_A\). If such a search appears in the literature we would expect more stringent bounds since the branching fraction \(BR(H \rightarrow ZA)\) would be reduced by the opening of the channels \(H^+H^-\) and \(W^\pm H^\mp\).
Figure 4: Scatter plots of the amount of mixing between the heavy Higgs and the radion, $K_H$, defined in equation (3.6), as function of the heavy Higgs mass. The black region is theoretically allowed and the points colored yellow, green and red are forbidden by heavy scalar searches in the $WW$, $ZZ$ and $hh$ channels respectively. The benchmark point $\Lambda = 3(5)$TeV was used on the left (right).

In figure 5 we show the production cross section, via gluon fusion, for $A$ times the branching fractions $BR(A \rightarrow ZX) BR(Z \rightarrow l^+l^-) BR(X \rightarrow bb)$ in the type-I (top) and type-II model (bottom) as a function of the mass $m_X$ where $X = H$ (red), $r$ (blue). The values $m_A = 700$ GeV and $\lambda_4 = 0.1$ were fixed.
Figure 5: The observable $\sigma(gg \to A \to ZX)\,BR(Z \to l^+l^-)\,BR(X \to b\bar{b})$ as a function of the resonance mass with $X = H$ (red), $r$ (blue) for type-I (top) and type-II (bottom) models. We fixed $\Lambda = 3$ TeV, $m_A = 700$ GeV and $\lambda_4 = 0.1$.

By visual inspection of the 95% CL upper limits from figure 5 (b) of Ref. [62] for $m_A = 700$ GeV one can infer an upper bound of about 0.3-3 fb at $m_H = 200$ GeV to 3-6 fb at $m_H = 600$ GeV after multiplying by $BR(Z \to l^+l^-) \approx 0.0336$ [63]. Using this rough estimates we can see that the predictions for this observable in our model lie within those limits for both radion and heavy Higgs in both models, but more encroachment from experiments is expected on the predicted range.

In figure 6 we show the production cross section via gluon fusion of a heavy Higgs boson (red) and a radion (blue) times the branching fractions $BR(X \to ZA)\,BR(Z \to l^+l^-)\,BR(A \to b\bar{b})$ as a function of the mass $m_X$ and with $X = H$, $r$ for the type-I (top) and type-II (bottom) models. For type-I model we fixed $m_A = 200$GeV and in the type-II, due to lower bounds on the charged Higgs [64], we fixed $m_A = 500$GeV.

For this decay topology we can use as rough estimates the upper left triangle of figure 1 of Ref. [60], in which the upper bounds for $m_A = 200$GeV are about 40fb at $m_X = 300$GeV descending with mass to about 2-3fb at $m_H = 1000$GeV. Similarly for $m_A = 500$GeV the upper limit is 30fb at $m_X = 600$GeV to 10fb at $m_X = 1000$GeV. Rough estimates from [61] can also be drawn but are not as stringent as those from [60]. We can see from figure 6 that our predictions are within these bounds, but further improvement of the experimental
Figure 6: The observable $\sigma(gg \rightarrow X \rightarrow ZA)BR(Z \rightarrow l^+l^-)BR(A \rightarrow b\bar{b})$ as a function of the resonance mass with $X = H$ (red), $r$ (blue) in the type-I (top) and type-II (bottom) models. We fixed $\Lambda = 3\text{TeV}$, $m_A = 200\text{GeV}$ ($m_A = 500\text{GeV}$) on top (bottom) and $\lambda_4 = 0.1$.

results would restrict the predicted region.

5 Conclusions

In this work we considered two Higgs doublets coupling to the Ricci scalar in the TeV-brane of an RS model. Assuming CP-conservation, the inclusion of this term causes kinetic mixing between the CP-even scalars of the 2HDM and the radion field of the RS model.

The most up to date LHC measurements of the signal strengths of the SM Higgs boson were used to fit the model and the allowed $\cos(\beta-\alpha)$-$\tan\beta$ parameter space for type-I and type-II 2HDM were presented.

We have discussed two possible ways to differentiate this model from other scenarios with similar scalar states. One possibility is to look at the sum of squared couplings of the scalars to gauge bosons. This model predicts a small deviation of about 0.5% from the SM value which could be measured at a future ILC. The other possibility is to look at the ratio of decay widths to a pair of $b$ quarks and $Z$ bosons for both scalars. Future experiments might distinguish the scalars by determining the value of the mixing angles $\alpha$ and $\beta$. 
Throughout this work we have taken the mass of the extra scalars to be in the range of 200-1000 GeV and we study the constraints that LHC searches of heavy resonances impose on the amount of mixing. The most stringent bounds arise if we take \( \Lambda = 3 \) TeV and \( m_A = 700 \) GeV where a radion is disfavored in the mass range \( m_r < 780 \) GeV while a heavy Higgs is disfavored in the mass range \( 300 \) GeV \( < m_H < 750 \) GeV and \( m_H < 250 \) GeV and kinetic mixing for both, radion and Higgs, is constrained to \( -4 < K_h, K_H < 4 \). These constraints relax significantly by reducing \( m_A \) and increasing the value \( \Lambda \).

Finally we showed how improvements of the experimental analysis for the decay topologies \( X \rightarrow ZA \) and \( A \rightarrow ZX \) where \( X = r \) or \( H \) could further constrain the parameter space of the model.

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Appendix A Scalar Couplings After Mixing

The interactions of the physical scalars to SM fields can be obtained by substituting the transformation of equation (3.13) into the unmixed couplings. A summary is given by

\[
g_{\phi V V} = U_{2\phi} \sin(\beta - \alpha) + U_{3\phi} \cos(\beta - \alpha) + U_{1\phi} \gamma \left( 1 - 3 \frac{m_{h}^{2} y_{t}}{\Lambda^{2}} \right) \quad \phi = r, h, H, \tag{A.1}
\]

\[
g_{\phi f f} = U_{2\phi} \xi_{h}^{f} + U_{3\phi} \xi_{H}^{f} + U_{1\phi} \gamma (c_{L}^{f} - c_{R}^{f}), \quad \phi = r, h, H, \tag{A.2}
\]

\[
g_{\phi g g} = \left( \frac{2\pi}{\alpha_{s} k_{y_{t}}} + \gamma \right) U_{1\phi} \gamma + \sum_{q} F_{q} (\xi_{q}^{h} U_{2\phi} + \xi_{q}^{H} U_{3\phi} + \gamma U_{1\phi}) \quad \phi = r, h, H. \tag{A.3}
\]

The trilinear interactions between scalar eigenstates \( r, h, \) and \( H \) are given by

\[
L \supseteq y_{1} r \partial^{\mu} h \partial_{\mu} H + y_{2} r \partial^{\mu} H \partial_{\mu} r + y_{3} r h \partial^{\mu} H + y_{r r h} r h H, \tag{A.4}
\]

where

\[
y_{1} = \frac{2}{\nu} \gamma \left\{ -6 \gamma [\xi_{1} \sin(\beta - \alpha) + \xi_{2} \cos(\alpha + \beta)] (U_{11} U_{12} U_{23} + U_{11} U_{13} U_{22} + U_{12} U_{13} U_{21}) \\
- 6 \gamma [\xi_{1} \cos(\beta - \alpha) + \xi_{2} \sin(\alpha + \beta)] (U_{11} U_{12} U_{33} + U_{11} U_{13} U_{32} + U_{12} U_{13} U_{31}) \\
+ 6 \xi_{2} U_{11} [\sin(2\alpha)(U_{23} U_{23} - U_{22} U_{23}) + \cos(2\alpha)(U_{22} U_{33} + U_{32} U_{33})] + 6 \xi_{1} U_{11} (U_{22} U_{23} \\
+ U_{32} U_{33}) - U_{11} U_{22} U_{23} - U_{11} U_{32} U_{33} + U_{12} U_{21} U_{23} + U_{12} U_{31} U_{33} + U_{13} U_{21} U_{22} \\
+ U_{13} U_{31} U_{32}) \right\}, \tag{A.5}
\]
\begin{equation}
y_2 = \frac{2}{v} \gamma \left\{ 3U_{11}(U_{22}U_{23} + U_{32}U_{33})\xi_1 + U_{13}(U_{21}U_{22} + U_{31}U_{32})(1 + 3\xi_1) \\
+ 3(U_{13}U_{22}U_{31} + U_{13}U_{21}U_{32} + U_{11}U_{23}U_{32} + U_{11}U_{22}U_{33})\xi_2 \cos(2\alpha) \\
- 6U_{13}(U_{12}U_{31} + 2U_{11}U_{32})\gamma_1 \cos(\alpha - \beta) - 6U_{13}(U_{12}U_{21} + 2U_{11}U_{22})\gamma_2 \cos(\alpha + \beta) \\
+ 3(-U_{13}U_{21}U_{22} - U_{11}U_{22}U_{33} + U_{13}U_{31}U_{32} + U_{11}U_{32}U_{33})\xi_2 \sin(2\alpha) \\
+ 6U_{13}(U_{12}U_{21} + 2U_{11}U_{22})\gamma_1 \sin(\alpha - \beta) \\
- 6U_{13}(U_{12}U_{31} + 2U_{11}U_{32})\gamma_2 \sin(\alpha + \beta) \right\}, \tag{A.6}
\end{equation}

\begin{equation}
y_3 = \frac{2}{v} \gamma \left\{ U_{12}(U_{21}U_{23} + U_{31}U_{33}) + 3(U_{13}U_{21}U_{22} + U_{11}U_{22}U_{32} + U_{13}U_{31}U_{32}) \xi_1 \\
+ 3(U_{13}U_{22}U_{31} + U_{13}U_{21}U_{32} + U_{11}U_{23}U_{32}) \xi_2 \cos(2\alpha) - 6(U_{12}U_{13}U_{31} + U_{11}U_{13}U_{32}) \xi_1 \cos(\beta - \alpha) - 6(U_{12}U_{13}U_{21} + U_{11}U_{13}U_{22}) \xi_2 \cos(\alpha + \beta) + 3(-U_{13}U_{21}U_{22} - U_{11}U_{22}U_{33} + U_{13}U_{31}U_{32}) \xi_2 \sin(2\alpha) + 6(U_{12}U_{13}U_{21} + U_{11}U_{13}U_{22}) \xi_1 \sin(\alpha - \beta) - 6(U_{12}U_{13}U_{31} + U_{11}U_{13}U_{32}) \xi_2 \sin(\alpha + \beta) \right\}. \tag{A.7}
\end{equation}

The tree-level coupling has two contributions, one from the trace of the energy-momentum tensor and another one from the 2HDM potential, i.e. $g_{rhH} = g_{rhH}^{\text{trace}} + g_{rhH}^{2\text{HDM}}$.
where
\[
g_{rhH}^{2HDM} = \frac{1}{2v} (\cos \beta (U_{33} \cos \alpha - U_{23} \sin \alpha)(U_{21} \cos \alpha + U_{31} \sin \alpha)(U_{22} \cos \alpha \\
+ U_{32} \sin \alpha)(m_A^2 - v^2 \lambda_4 - (m_h^2 - m_H^2) \cos \alpha \csc \beta \sec \beta \sin \alpha) \\
+ \cos \beta (U_{32} \cos \alpha - U_{22} \sin \alpha)(U_{21} \cos \alpha + U_{31} \sin \alpha)(U_{23} \cos \alpha \\
+ U_{33} \sin \alpha)(m_A^2 - v^2 \lambda_4 - (m_h^2 - m_H^2) \cos \alpha \csc \beta \sec \beta \sin \alpha) \\
+ \cos \beta (U_{31} \cos \alpha - U_{21} \sin \alpha)(U_{22} \cos \alpha + U_{32} \sin \alpha)(U_{23} \cos \alpha \\
+ U_{33} \sin \alpha)(m_A^2 - v^2 \lambda_4 - (m_h^2 - m_H^2) \cos \alpha \csc \beta \sec \beta \sin \alpha) \\
+ (U_{32} \cos \alpha - U_{22} \sin \alpha)(U_{33} \cos \alpha - U_{23} \sin \alpha)(U_{21} \cos \alpha \\
+ U_{31} \sin \alpha)(m_A^2 - v^2 \lambda_4 - (m_h^2 - m_H^2) \cos \alpha \csc \beta \sec \beta \sin \alpha) \sin \beta \\
+ (U_{31} \cos \alpha - U_{21} \sin \alpha)(U_{33} \cos \alpha - U_{23} \sin \alpha)(U_{22} \cos \alpha + U_{32} \sin \alpha)(m_A^2 \\
- v^2 \lambda_4 - (m_h^2 - m_H^2) \cos \alpha \csc \beta \sec \beta \sin \alpha) \sin \beta + (U_{31} \cos \alpha \\
- U_{21} \sin \alpha)(U_{32} \cos \alpha - U_{22} \sin \alpha)(U_{23} \cos \alpha + U_{33} \sin \alpha)(m_A^2 - v^2 \lambda_4 \\
- (m_h^2 - m_H^2) \cos \alpha \csc \beta \sec \beta \sin \alpha) \sin \beta - 6(U_{21} \cos \alpha \\
+ U_{31} \sin \alpha)(U_{22} \cos \alpha + U_{32} \sin \alpha)(U_{23} \cos \alpha + U_{33} \sin \alpha)((m_A^2 \\
+ v^2 \lambda_4) \cos \beta - \csc^2 \beta ((m_h^2 \cos^2 \alpha + m_H^2 \sin^2 \alpha)) \sin \beta + 6 \sec \beta (U_{31} \cos \alpha \\
- U_{21} \sin \alpha)(U_{32} \cos \alpha - U_{22} \sin \alpha)(U_{33} \cos \alpha - U_{23} \sin \alpha)(m_A^2 \cos \alpha^2 + m_h^2 \sin \alpha^2 \\
- (m_h^2 + v^2 \lambda_4) \sin \beta - 2v^2 \lambda_4((U_{23}U_{32} + U_{22}U_{33}) \cos(2\alpha) + (-U_{22}U_{23} \\
+ U_{32}U_{33}) \sin(2\alpha))(U_{21} \cos(\alpha + \beta) + U_{31} \sin(\alpha + \beta)) + 2v^2 \lambda_4((U_{23}U_{31} \\
+ U_{21}U_{33}) \cos(2\alpha) + (-U_{21}U_{23} + U_{31}U_{33}) \sin(2\alpha))(U_{22} \cos(\alpha + \beta) \\
+ U_{32} \sin(\alpha + \beta)) + 2v^2 \lambda_4((U_{22}U_{31} + U_{21}U_{32}) \cos(2\alpha)(-U_{21}U_{22} \\
+ U_{31}U_{32}) \sin(2\alpha)(U_{23} \cos(\alpha + \beta) + U_{33} \sin(\alpha + \beta))), \tag{A.8}
\]

\[
g_{rhH}^{\text{trace}} = \frac{1}{2v} (m_h^2(U_{13}U_{21}U_{22} + U_{12}U_{21}U_{23} + U_{11}U_{22}U_{23}) \\
+ m_H^2(U_{13}U_{31}U_{32} + U_{12}U_{31}U_{33} + U_{11}U_{32}U_{33})). \tag{A.9}
\]

The other interactions like \(rhh, rHH\), etc. can be similarly obtained and are not illustrated here.
## Appendix B  LHC Data

### Table 2: Measured Higgs Signal Strengths

| Decay | Production | Measured Signal Strength $R_m$ |
|-------|------------|--------------------------------|
| $\gamma\gamma$ | ggF+tth | $1.19^{+0.20}_{-0.18}$ [CMS] [65] |
| | VBF +Vh | $1.01^{+0.57}_{-0.51}$ [CMS] [65] |
| | ggF | $0.8^{+0.19}_{-0.18}$ [ATLAS] [66] |
| | VBF | $2.1^{+0.6}_{-0.6}$ [ATLAS] [66] |
| | Vh | $0.7^{+0.9}_{-0.8}$ [ATLAS] [66] |
| WW* | ggF | $1.02^{+0.29}_{-0.26}$ [ATLAS] [67] |
| | VBF | $1.27^{+0.53}_{-0.45}$ [ATLAS] [67] |
| | ggF | $0.76 \pm 0.21$ [CMS] [68] |
| | VBF | $1.7^{+1.1}_{-0.9}$ [ATLAS] [69] |
| | Wh | $3.2^{+4.4}_{-4.2}$ [ATLAS] [69] |
| ZZ* | ggF | $1.7^{+0.5}_{-0.4}$ [ATLAS] [70] |
| | VBF + Vh | $0.3^{+1.6}_{-0.9}$ [ATLAS] [70] |
| | ggF | $1.20^{+0.35}_{-0.31}$ [CMS] [71] |
| | VBF | $0.00^{+1.37}_{-0.30}$ [CMS] [71] |
| bb | VBF | $-3.7^{+2.4}_{-2.5}$ [CMS] [72] |
| | Vh | $1.26^{+0.42}_{-0.36}$ [ATLAS] [73] |
| | Vh | $1.2 \pm 0.4$ [CMS] [74] |
| $\tau\tau$ | VBF | $1.2 \pm 0.4$ [ATLAS] [75] |
| | ggF | $2.0^{+1.5}_{-1.2}$ [ATLAS] [76] |
| | VBF + Vh | $1.24^{+0.50}_{-0.54}$ [ATLAS] [76] |
| | WH | $2.3 \pm 1.6$ [ATLAS] [77] |
| | tth | $1.5^{+1.2}_{-1.0}$ [ATLAS] [78] |
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