Higgs-Radion Mixing with Enhanced Di-Photon Signal

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Abstract
In the context of warped scenarios in which Standard Model (SM) fields are allowed to propagate in the bulk, we revisit the possible mixing between the IR localized Higgs field and the Radion graviscalar. The phenomenology of the resulting mostly-Higgs field does not suffer important deviations with respect to the case in which all the SM is localized in the IR brane (original Higgs-Radion mixing scenario). On the contrary, the phenomenology of the mostly-Radion field can present important differences with respect to the original scenario. At the LHC, the most striking effect is now the possibility of sizeable Radion decays into photons in a mass range well beyond the ZZ and WW thresholds, not due to dramatically enhanced couplings to photons but to suppressed couplings to massive fields.

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I. INTRODUCTION

Models of warped extra dimensions have generated a lot of attention as an interesting and novel framework in which to address simultaneously issues such as the hierarchy problem and fermion mass hierarchies. In the original setup of Randall and Sundrum [1], the hierarchy problem is dealt with by localizing all the SM particles in the IR brane. We will refer to this original proposal as RS1 when comparing it with more recent proposals in which SM fields are allowed to propagate in the bulk, which will be loosely referred to as Bulk Fields scenarios.

Precision tests from Electroweak observables put strong bounds on extensions of the Standard Model (SM), and in the case of the RS1 proposal one should expect higher dimension operators of IR fields to contribute too importantly to these observables given that the IR cutoff scale is in the TeV region. By allowing gauge fields (of a generically extended gauge group) and fermions to propagate in the bulk one might effectively suppresses the contribution of higher dimension operators containing the SM fields [2, 3, 4, 5, 6, 7, 8, 9, 10, 11]. The Higgs field should be close to the IR boundary if the scenario is to address the hierarchy problem, and now as a benefit of having bulk fermions, one can use their geographical location to explain the wide differences in their couplings with the IR localized Higgs. But precision electroweak tests as well as tight bounds from flavor physics will now force the new KK modes of the bulk fields to be heavier than a few TeV [2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16].

In all the previous scenarios there is one scalar which, like the Higgs must be located near the IR brane. It is a scalar mode of the 5D gravitational fluctuations, parametrizing a vibration mode of the inter-brane proper distance, the Radion [1, 17, 18, 19]. The original setup predicts its mass to be zero due to the fact that the actual inter-brane distance is not fixed by the spacetime background setup. Much research has been devoted to study the stabilization of this setup and therefore the mechanism to generate a Radion mass [17, 20, 21, 22, 23]. The phenomenological interest in the Radion lies in the fact that its interactions with SM matter are TeV scale and therefore could be observable in high energy collisions [21, 22, 26, 27, 28, 29]. Moreover its mass remains more or less a free parameter, and with the assumption that the stabilization mechanism does not alter importantly the gravitational background, the Radion is expected to be a much lighter field than the rest of KK excitations [22].
As already mentioned, the Radion is located near the IR Brane and its interactions are generically proportional to the mass of the fields it couples to (it couples through the trace of the energy momentum tensor). These same attributes are shared by the Higgs scalar and so it becomes important to study carefully the phenomenology of the full scalar sector to understand how to distinguish between fields. As it turns out the Higgs and Radion can also mix through a gravitational kinetic mixing term [24] opening the door to even more interesting consequences [22, 26, 27, 28]. Surprisingly most of the research related to the Radion was limited to the RS1 scenario (except for [31]), and only recently this research gap was addressed [32]. We plan to extend the study of Radion phenomenology when Fields are in the Bulk by allowing for the possibility of Higgs-Radion mixing. In Section II we will review the Radion setup and its couplings before any Higgs-Radion mixing, which will then be introduced. In section III we will compare the phenomenology of the Radion between the RS1 scenario and the Fields in the Bulk models. The most striking difference will lie in the di-photon channel and we will concentrate our attention mostly on this channel. Finally in section IV we will present our conclusions.

II. SETUP

The spacetime structure consists of one extra dimension with warping such that the metric takes the usual Randall-Sundrum form [1]:

$$ds^2 = e^{-2\sigma} \eta_{\mu\nu} dx^\mu dx^\nu - dy^2$$

where $\sigma(y) = ky$, and $k$ is the 5D curvature. This is a 4D Poincare invariant metric solution in a 5D setup with bulk cosmological constant and fine-tuned brane tensions at the two boundaries $y = 0$ and $y = y_{IR}$ of the extra dimension.

We assume that the origin of Electroweak Symmetry Breaking (EWSB) is localized in the TeV brane and is well described by a Higgs doublet. The SM fermions do have profiles along the extra dimension but their couplings with the Higgs are also localized in the IR brane, and will depend on the value of their wave-functions at that boundary.

In the gravitational sector, one needs to add full 5D tensor perturbations $h_{AB}(x,y)$ around the metric background $g_{AB}^{RS}$ of Eq. (1), i.e.

$$g_{AB} = g_{AB}^{RS} + \hat{\kappa} \ h_{AB}$$
where \( \hat{\kappa} \) is a small parameter.

Thanks to the 5D diffeomorphism invariance we can reduce some of the linear metric perturbation degrees of freedom to obtain the simpler perturbative metric [19, 30]:

\[
\begin{align*}
\text{ds}^2 &= \left( e^{-2\sigma} \left[ \eta_{\mu\nu} + \hat{\kappa} h_{\mu\nu}^{TT}(x,y) \right] - \hat{\kappa} \eta_{\mu\nu}r(x) \right) dx^\mu dx^\nu - \left( 1 + \hat{\kappa} 2 e^{2\sigma} r(x) \right) dy^2
\end{align*}
\]  

(3)

where \( h_{\mu\nu}^{TT}(x,y) \) is transverse and traceless and \( r(x) \) is the Radion graviscalar which cannot be gauged away due to the presence of the two brane boundaries. In the absence of a stabilization mechanism the Radion is massless (and therefore a problematic long-range force mediator), but it was quickly realized that a very simple fix to this was to add an extra bulk scalar field to the setup, such that it acquires a nontrivial background vev along the extra dimension [17]. This space-time background solution will fix the inter-brane distance and generically give a positive mass squared to the Radion. The stabilized background metric solution can be in some limit very close to the Randall-Sundrum solution, so that we can still use Eq. (1) as the background metric. We will refer to this limit as the “small back-reaction limit” and will assume it for the remaining of the paper. We will consider the radion mass as a free parameter although we should expect it to be relatively light in the small back-reaction limit, at least when invoking a Goldberger-Wise type stabilization mechanism [22].

A. Interactions

From the perturbative ansatz of Eq. (3), it is simple to extract the tree-level interactions between the Radion and the matter fields since these are just gravitational interactions. One obtains the linear Radion-matter interactions as [31, 32]

\[
S_{1}^{\text{int}} = -\frac{\hat{\kappa}}{2} \int d^5x \; e^{-2\sigma} \left( -T^\mu_\mu + 2T_{55} \right) \; r(x)
\]

(4)

where the graviscalar field \( r(x) \) is not a 4D canonically normalized scalar field. The physical Radion field \( \phi_0(x) \) is obtained with the redefinition

\[
\frac{\hat{\kappa}}{2} r(x) = -\frac{1}{\Lambda_\phi} \phi_0(x)
\]

(5)

where \( \Lambda_\phi = \sqrt{6} M_{Pl} e^{-k y_{IR}} \) and the minus sign restores the convention of [27]. The 5D matter field information is included in \( T_{AB} \), the energy momentum tensor. From the previous tree-level interactions one needs to extract the interactions between the Radion and the lightest
modes of the 5D bulk matter, i.e the SM massive gauge bosons and fermions. Assuming
that the 5D electroweak group is simply $SU(2) \times U(1)$, and that the fermion structure is
a simple 5D extension of the Standard Model with a Higgs mechanism localized on the IR
brane, the couplings to the Radion are \[32\]

$$M_V^2 \left(1 - 6 k y_{IR} M_V^2 \right) \frac{\phi_0}{\Lambda_\phi} V^\alpha V_\alpha, \quad (6)$$

$$m_f (c_L - c_R) \frac{\phi_0}{\Lambda_\phi} \bar{f}_{UV} f_{UV}, \quad (7)$$

$$m_f \frac{\phi_0}{\Lambda_\phi} \bar{f}_{IR} f_{IR} \quad (8)$$

where $f_{UV}$ and $f_{IR}$ represent fermions localized near the UV and IR brane respectively, with
$c_L$ and $c_R$ being the mass parameters associated to the left and right handed bulk fermions.
The term proportional to $k y_{IR}$ in the massive boson couplings and the term proportional to
the fermion mass parameters $c_i$ appearing in the fermion couplings are new contributions
due to the 5D nature of matter fields, and were not present in the RS1 scenario.

In the case of massless gauge bosons, i.e. gluons and photons, the interactions with the
Radion appear with same strength from various sources \[32\]. First, there is a one-loop con-
tribution identical to the Higgs radiative couplings; also, because the gauge interactions are
5-dimensional there is a tree-level interaction with photons which can be directly extracted
from Eq. (4). Brane kinetic terms associated with the gauge fields will also contribute, if
present, and finally there is a term proportional to the total gauge group beta function
coefficient, coming from the trace anomaly for IR light fields and from loop corrections to
the gauge coupling due to UV and bulk fields. We write

$$\left[1 - 4\pi \alpha (\tau_{UV}^0 + \tau_{IR}^0) \right] + \frac{\alpha}{8\pi} \left( b - \sum_i \kappa_i F_i (\tau_i) \right) \frac{\phi_0}{\Lambda_\phi} F_{\mu\nu} F^{\mu\nu}. \quad (9)$$

where $\tau_{UV}^0$ and $\tau_{IR}^0$ are the brane kinetic terms, $\sum \kappa_i F_i$ are the contributions from the one-
loop diagrams and $b$ is the total beta function coefficient associated with the corresponding
gauge field. The first term in this formula is a consequence of the 5-dimensional nature of
the gauge fields and again was not present in the RS1 scenario.

In Fig. 1 we show the branchings of the Radion in the absence of Higgs-Radion mixing
for the two scenarios we wish to compare, RS1 and the Fields in the Bulk. In this last case
we have not included any brane kinetic term associated with the bulk (gauge) fields and we
FIG. 1: Branching fractions for the Radion as a function of its mass $m_\phi$ in the RS1 scenario (dotted curves) and the Fields in the Bulk scenario (thick curves). The individual curves are very similar for both scenarios except for the $\gamma\gamma$ channel (blue curve) which drops quickly with the Radion mass in the RS1 case but becomes flat in the Fields in the Bulk case.

will not include them either in the rest of this study$^1$. The figure agrees reasonably well with $[32]$, the main difference being that we have included here decays into one intermediate off-shell $Z$ or $W$ boson, a process which can still be quite significant below the $ZZ$ and $WW$ physical thresholds. One observes that the branchings do not vary much between the two scenarios and we mainly point out for the case of the Fields in the Bulk, a slight increase in $gluon-gluon$ branchings and an interesting plateau for the $\gamma\gamma$ branchings in the high Radion mass range.

\[ S_i^{BKT} = \pm M_i^2 \int d^4x \, \sqrt{-g_i} \, R_i \]  

$^1$ Brane kinetic terms for photons and/or gluons will actually have interesting effects in both suppressing or enhancing the $gluon-gluon$ and the $\gamma\gamma$ decays as shown in $[32]$. 

B. Higgs-Radion mixing

One can still modify this setup since it is always possible to write down localized gravity kinetic terms in the boundaries,
where $g_i$ is the determinant of the induced metric on the $i$-th brane, and $M_i$ is some dimensionful parameter that should be naturally related to the fundamental scales $M_5$ and $k$. These terms will contribute to the kinetic term of the Radion and their presence will therefore require a new canonical normalization. In the same lines, in the IR brane we have at our disposal the Higgs bilinear $H^+ H$ which can be coupled to the induced Ricci scalar $R$ to obtain an effective dimension-4 operator \[ S_\xi = \xi \int d^4x \sqrt{-g_{IR}} \ R H^+ H \] where $\xi$ is a dimensionless parameter.

In the small back-reaction limit, and after EWSB, the two lightest (mixed) states of the scalar sector will be the Radion graviscalar and the Higgs; the effective 4d action for these fields up to quadratic order will be:

\[
\mathcal{L} = -\frac{1}{2} \left\{ 1 \pm \frac{6M_{IR}^2}{\Lambda_\phi^2} + 6\gamma^2 \xi \right\} \phi_0 \Box \phi_0 - \frac{1}{2} \phi_0 m_{\phi_0}^2 \phi_0 - \frac{1}{2} h_0 (\Box + m_h^2) h_0 - 6\gamma \xi \phi_0 \Box h_0 ,
\]

where $m_{h_0}$ and $m_{\phi_0}$ are the Higgs and Radion masses before mixing and $\gamma = v/\Lambda_\phi$ is a dimensionless parameter reflecting the suppression of the Radion-to-matter couplings with respect to the Higgs-to-matter couplings. For simplicity and to make close contact with previous work, we will set $M_{IR} = 0^2$.

The states $h(x)$ and $\phi(x)$ which diagonalize (12) can be introduced as

\[
\begin{pmatrix}
 h_0 \\
 \phi_0
\end{pmatrix} = \begin{pmatrix} d & c \\
 b & a
\end{pmatrix} \begin{pmatrix} h \\
 \phi
\end{pmatrix}
\]

It is interesting to remark that this transformation is not orthogonal due to the nature of the mixing (kinetic mixing and extra contribution to the Radion kinetic term). We can decompose the previous transformation in terms of two operations. The first redefinition diagonalizes and normalizes the kinetic terms, and the second one, an orthogonal transformation this time, diagonalizes the mass matrix. To maintain positive definite kinetic energy terms for the Radion $\phi$, we must have $Z^2 > 0$, where

\[
Z^2 \equiv 1 + 6\xi \gamma^2 (1 - 6\xi) .
\]

\[\text{Non-zero } M_{IR} \text{ will only affect the Radion kinetic term and can be recast as an overall modification of the “bare” Radion couplings to matter (i.e. a redefinition of } \Lambda_\phi) . \text{ The effects of such terms on the KK graviton spectrum and couplings have been studied in } [33] .\]
This allows us to obtain theoretical limits on the $\xi$ parameter in terms of the scale $\Lambda_\phi$ (i.e. the parameter $\gamma = v_0/\Lambda_\phi$):

$$ \frac{1}{12} \left( 1 - \sqrt{1 + \frac{4}{\gamma^2}} \right) \leq \xi \leq \frac{1}{12} \left( 1 + \sqrt{1 + \frac{4}{\gamma^2}} \right). $$

(15)

The corresponding mass-squared eigenvalues of the new physical states are

$$ m^2_\pm = \frac{1}{2Z^2} \left( m^2_{\phi_o} + \beta m^2_{h_o} \pm \sqrt{[m^2_{\phi_o} + \beta m^2_{h_o}]^2 - 4Z^2 m^2_{\phi_o} m^2_{h_o}} \right), $$

(16)

where we have defined $\beta = 1 + 6\xi \gamma^2$. These masses must satisfy the inequality

$$ \frac{m^2_+}{m^2_-} > 1 + \frac{2\beta}{Z^2} \left( 1 - \frac{Z^2}{\beta} \right) + \frac{2\beta}{Z^2} \left[ 1 - \frac{Z^2 \gamma}{\beta} \right]^{1/2}, $$

(17)

in order for the ‘bare’ masses $m^2_{\phi_o}$ and $m^2_{h_o}$ to be real [27]. This constraint on the physical masses is quite interesting since the larger the value of $|\xi|$, the larger the mass-squared difference of the two scalar fields must be.

The parameters needed to fix the scalar sector are the Higgs and Radion ‘bare’ masses $m^2_{\phi_o}$ and $m^2_{h_o}$ before mixing, the Radion interaction scale $\Lambda_\phi$ (related to the KK masses and the solution of the hierarchy problem) and the mixing parameter $\xi$. We will trade the ‘bare’ masses of the original fields for the phenomenologically more interesting physical masses of the two mixed scalar states. We will refer to these as the Higgs and the Radion, even though they are an admixture of both, and will fix the convention by defining the Higgs scalar as the field which becomes the SM Higgs in the limit of $\xi \to 0$ and similarly for the Radion. The important parameters are thus

$$ m_h, \ m_\phi, \ \Lambda_\phi \ and \ \xi, $$

with the understanding that they are not completely independent since the masses are bound by the non-degeneracy constraint from Eq. (17), which defines a theoretically forbidden region for these parameters.

### III. PHENOMENOLOGY

The couplings of the physical scalar fields with the SM matter fields will be obtained using the redefinitions of Eq. (13). Calling $g^0_{h_{ii}}$ and $g^0_{\phi_{ii}}$ the coefficients of the ‘bare’ Higgs

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3 Note that the quantity inside the square root is positive definite so long as $m^2_{h_o} m^2_{\phi_o} > 0$. 

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and Radion couplings to the fields “i”, one obtains

\[ g_{hii} = d \, g_{h_{0}ii} + b \, g_{\phi_{0}ii} \]  
(18)

\[ g_{\phi_{0}ii} = c \, g_{h_{0}ii} + a \, g_{\phi_{0}ii} \]  
(19)

for couplings of the physical fields.

Now, the couplings \( g_{\phi_{0}ii} \) of the ‘bare’ Radion are basically the same as the couplings of the ‘bare’ Higgs \( g_{h_{0}ii} \), but suppressed by a factor of \( \sim v/\Lambda_{\phi} \), where \( v \) is the Higgs vev.

When the mixing parameter \( \xi \) is small, the redefinition coefficients \( b \) and \( c \) must be small too (basically they play the role of \( \sin \theta \) in the case of a typical orthogonal mass mixing), whereas \( d \) and \( a \) must lie close to 1 (like \( \cos \theta \) for small \( \theta \)). A quick glance at the previous couplings shows that for small mixing, the couplings will look like

\[ g_{h_{0}ii} \sim g_{h_{0}ii}(1 + b \frac{v}{\Lambda_{\phi}}) \]  
(20)

\[ g_{\phi_{0}ii} \sim g_{h_{0}ii}(c + \frac{v}{\Lambda_{\phi}}) \]  
(21)

where both \( b \) and \( c \) are small numbers. One sees that the Higgs couplings do not receive much corrections, since \( b \) is small and is multiplied by the also small \( v/\Lambda_{\phi} \). Of course for larger values of the mixing parameter \( \xi \), the couplings of the mostly-Higgs scalar will start deviating significantly from the SM Higgs values. This effect was extensively studied in [27] in the context of RS1 and we will not pursue it further here, but instead concentrate on the mostly-Radion sector were the couplings are more sensitive in the small mixing region. Moreover, when the mixing is large one should also consider carefully the bounds coming from precision electroweak constraints, as now both the Radion and the Higgs are expected to contribute importantly to electroweak observables such as S and T [22, 28].

On the other hand, the Radion couplings can quickly change for small mixing since now, even if \( c \) is small, it is to be compared with \( v/\Lambda_{\phi} \), which is small too (see Eq. (21)). In particular this means that for a small mixing in the appropriate direction, one can actually suppress completely the physical coupling of the Radion to the fields “i”. Of course, in principle, the point where the Radion couplings to the “i” particles vanish does not mean that the couplings to some different “j” particles also vanish.

Nevertheless, in the RS1 scenario in which all fields lie on the IR brane, it turns out that the point where the physical Radion is phobic to W bosons, is the same where it is phobic to Z bosons and to all massive fermions. Moreover, the suppressed couplings to \( \gamma \gamma \) happen
FIG. 2: Branching fractions for the Radion as a function of its mass $m_\phi$ in the RS1 scenario (left panel) and the Fields in the Bulk scenario (right panel). We have taken a small and negative mixing parameter value, $\xi = -0.1$, and a physical Higgs mass of $m_h = 150$ GeV. In both panels, the vertical gray band centered at the physical Higgs mass represents the theoretically excluded region in which the values of the ‘bare’ scalar masses are complex. The individual curves are very similar in both panels except for the $\gamma\gamma$ channel (blue curve) which shows different features in each scenario reaching higher maximum values in the Fields in the Bulk scenario.

in a nearby point. These features are a consequence of the exact Higgs-like structure of the Radion couplings to massive particles in RS1 (vector bosons and fermions) where by simply replacing the parameter $\Lambda_\phi$ by the Higgs vev $v$, one obtains the Higgs couplings. Moreover the trace anomaly contribution to the Radion coupling to $\gamma\gamma$ is numerically small, and therefore this coupling still resembles the Higgs one. In the case of the Radion coupling to gluons the contribution from the trace anomaly is quite important and therefore the zero for that coupling is very separated in parameter space from the points where the other couplings vanish.

In the left panels of both Fig. 2 and Fig. 3, we plot the Radion decay branching fractions with respect to its mass for the RS1 scenario, for two different small mixing parameters, $\xi = -0.1$ and $\xi = 0.1$. We have taken $\Lambda_\phi = 2000$ GeV and the mass of the mostly-Higgs scalar is $m_h = 150$ GeV. In both plots one observes that there is a ‘universal’ point where all the couplings to massive particles vanish. The $\gamma\gamma$ signal is suppressed in a nearby region, so one cannot take advantage of the suppression of other couplings, and the gluon branchings are suppressed only in the region close to the theoretically forbidden boundary, represented
FIG. 3: Same as Fig. 2 but with $\xi = +0.1$. Again the $\gamma\gamma$ channel (blue curve) shows very different features in each scenario, and in particular, in the right panel, one observes strikingly high values in this channel for a range of Radion masses beyond the ZZ and WW threshold. This effect is mainly due to a suppression in the couplings of the Radion to massive gauge bosons.

in the plots as a vertical gray band, centered at the physical Higgs mass, here $m_h = 150$ GeV (The Radion and the Higgs masses cannot lie in the same region, as explained below Eq. (17)).

When the matter fields are placed in the 5D bulk, the couplings of the Radion receive some corrections relative to the RS1 case, as explained below Eqs. (6-8). The consequence of this is that now the zeroes of the different Radion couplings happen in separated points of parameter space. In the case of the $WW$ coupling and the $ZZ$ coupling this separation is not very substantial, although it does happen. Couplings to different families and types of fermions will have zeros in different regions of parameter space, depending on the values of the mass parameters $c^i_{L,R}$ for each fermion. But the most important change comes from the new contribution to the Radion coupling to photons (see Eq. (9)), where the sign of the overall coupling actually flips, even if retaining a similar absolute value. Therefore, when we turn on the $\xi$-mixing, the coupling of the Radion to photons will vanish in a very different region where it used to in the RS1 scenario. This feature is apparent in the right panels of both Fig. 2 and Fig. 3, where one observes clearly how the suppression to $\gamma\gamma$ happens far away from the other zeroes (namely on the opposite side of the forbidden vertical band). Then, the branching fraction to photons increases substantially in the region where the Radion is phobic to massive particles. Specially in the right panel of Fig. 3 we
FIG. 4: **Left Panel:** Ratio of discovery significances $R_{SZ}^{ZZ} = S(gg \rightarrow \phi \rightarrow ZZ)/S(gg \rightarrow h_{SM} \rightarrow ZZ)$ in the $ZZ$ channel between the Radion and a SM Higgs of same mass as a function of $\xi$. **Right Panel:** Cross section (in fb) for the process $pp \rightarrow \phi \rightarrow \gamma\gamma$ as a function of $\xi$. In both panels we take $\Lambda_\phi = 2$ TeV and the radion mass $m_\phi = 250$ GeV, and the dotted (solid) curves are for the RS1 (Bulk Fields) scenario. We take two limits for the mass of the mostly-Higgs scalar, $m_h = 120$ GeV (Purple) and $m_h = 800$ GeV (Green). The dependance on $\xi$ in the $ZZ$ channel is similar for both RS1 and Bulk Fields scenarios, whereas the $\gamma\gamma$ signal fluctuates quite importantly with $\xi$ in the Bulk Fields scenario.

also observe how the zero to $b\bar{b}$ is moved away from the $WW$ and $ZZ$ zeroes. For this, we used $c_L - c_R = 1.1$ for the bottom quark 5D mass parameters. One could study in detail the variations of the couplings to other fermions such as tau’s and charm quarks, since the branchings could change importantly if one happens to live near a zero coupling for one of these heavier fermions. At the LHC, generically the Radion is mainly produced in gluon fusion, with other production mechanisms extremely suppressed (unlike the SM Higgs case). In that case most fermionic decays will be hard to study due to the enormous QCD background. When the Higgs-Radion mixing is large, production via vector boson fusion for example can be enhanced importantly [27], therefore opening the door to study decays into tau’s with associated forward and backward jets, just like in the SM Higgs case. We will not pursue further this line of investigation although it might be an interesting one for the future.
Instead we plan to focus on the Radion decays into photons since the branchings can be substantially enhanced for both light masses, $110 \text{ GeV} < m_\phi < 150 \text{ GeV}$, but also more massive ones, $m_\phi > 150 \text{ GeV}$, as shown in the right panel of Fig. [3]. But before concentrating on the $\gamma\gamma$ channel we want to also take a look at the Radion decays into ZZ, a very important channel for the large mass region. In [32] it was observed that, in the absence of brane gauge kinetic terms, the ZZ signal is relatively enhanced with respect to the RS1 scenario. In the presence of Higgs-Radion mixing one would expect this enhancement to remain similar. This is confirmed in the left panel of Figure 4. The ratio of discovery significances $R_{S}^{ZZ} = S(gg \rightarrow \phi \rightarrow ZZ)/S(gg \rightarrow h_{SM} \rightarrow ZZ)$, as defined in [24, 32], in the ZZ channel between the Radion and a SM Higgs of same mass is plotted as a function of $\xi$. The dotted curves are for the RS1 scenario, while the solid ones are for the Fields in the Bulk scenario. We have taken $\Lambda_\phi = 2000 \text{ GeV}$ and a Radion mass of $m_\phi = 250 \text{ GeV}$, and chose a light Higgs scenario ($m_h = 120 \text{ GeV}$) and a heavy Higgs one ($m_h = 800 \text{ GeV}$). In the two cases the RS1 curves and the Bulk Field curves follow each other closely, as expected.

In the right panel we plot the cross section (in fb) for the process $pp \rightarrow \phi \rightarrow \gamma\gamma$ as a function of $\xi$. We have computed it for $\sqrt{s} = 14 \text{ TeV}$ without QCD corrections and used CTEQ5L pdf’s. The difference between the RS1 scenario and the Bulk Fields is quite striking, and the cross sections can reach up to $20 – 30 \text{ fb}$ for the parameters chosen. This is roughly the level of cross sections that one expects for a SM Higgs with a mass of 120-130 GeV, even though the Radion mass taken here is quite large, $m_\phi = 250 \text{ GeV}$.

This prompts us to do a parameter space scan to see how large is the region where the di-photon signal is important (and enhanced with respect to the RS1 scenario). In Figure 5 we choose the mass of the mostly-Higgs field as $m_h = 150 \text{ GeV}$ and the Radion coupling scale at $\Lambda_\phi = 2 \text{ TeV}$. We then plot, in the $(\xi, m_\phi)$ plane, contours of the ratio of discovery significances $R_{S}^{\gamma\gamma} = S(gg \rightarrow \phi \rightarrow \gamma\gamma)/S(gg \rightarrow h_{SM} \rightarrow \gamma\gamma)$, as defined in [24, 32], between the Radion and a SM Higgs with equal mass. In RS1 (left panel), for positive values of $\xi$ one can reach roughly up to 2 times better than a SM Higgs. For the Bulk Fields scenario (right panel), the enhancement happens for negative values of $\xi$, and reaches quite larger values. This is because the zeroes of the couplings to $W$’s and $Z$’s happen to be in the negative $\xi$ region while the zeroes of the $\gamma\gamma$ coupling are now in positive $\xi$ region. In the mass region considered for this figure (below WW and ZZ thresholds), and at the points of larger di-photon signal, the collider phenomenology would be quite similar to exotic Higgs
FIG. 5: Contours in the \((\xi, m_\phi)\) plane of the ratio of discovery significances \(R_{S\gamma} = S(gg \to \phi \to \gamma\gamma)/S(gg \to h_{SM} \to \gamma\gamma)\), with the SM Higgs mass equal to that of the Radion. The mass of the mostly-Higgs scalar of the setup is held at \(m_h = 150\) GeV and the Radion coupling scale at \(\Lambda_\phi = 2\) TeV. The left panel shows results for the model in which all SM fields are confined in the TeV Brane (here referred to as RS1) and the right panel shows results for the case in which all fields propagate in the bulk except the Higgs, which remains confined in the TeV brane.

scenarios with enhanced di-photon branchings (see for e.g. \[34\]).

The gray region is the theoretically excluded region while the yellow regions are parameter points excluded by LEP data \[35\], given that the Radion couplings to \(ZZ\) are quite enhanced in those corners thereby reaching the bounds set by LEP, as previously remarked in \[27\].

Surprisingly one can still obtain interesting rates in the di-photon channel for a heavier Radion. Of course a comparison with a SM Higgs of same mass becomes useless since this channel is extremely suppressed above \(\sim 160\) GeV. Therefore for the large Radion mass range, we decide to plot in both panels of Figure 6 contours in the \((\xi, m_\phi)\) plane of the di-photon cross section \(\sigma(pp \to \phi)Br(\phi \to \gamma\gamma)\) (in fb) at the LHC with \(\sqrt{s} = 14\) TeV (using CTEQ5L gluon PDF). We have not included QCD corrections, which tend to enhance the production cross section, but these are the same for the SM Higgs process and one can easily estimate their effect. In any case our main interest lies in comparing different scenarios so we can confidently use the leading order results. As in Figure 5 we take \(m_h = 150\) GeV and \(\Lambda_\phi = 2\) TeV, and again we observe a striking difference between the RS1 scenario (left panel) and the Bulk Fields scenario (right panel). In the RS1 case the cross section above 200 GeV is at most 1 fb making this signal quite difficult at least in the first years of running.
FIG. 6: Contours in the \((\xi, m_s)\) plane of the di-photon cross section \(\sigma(pp \to \phi) \times Br(\phi \to \gamma\gamma)\) (in fb) at the LHC with \(\sqrt{s} = 14\) TeV (using CTEQ5L gluon PDF). The mass of the mostly–Higgs scalar of the setup is held at \(m_h = 150\) GeV and the Radion coupling scale at \(\Lambda_\phi = 2\) TeV. The left panel shows results for the model in which all SM fields are confined in the TeV Brane (here referred to as RS1) and the right panel shows results for the case in which all fields propagate in the bulk except the Higgs, which remains confined in the TeV brane.

On the other hand, when fields are in the Bulk, one sees that there is a thin tower region in the allowed parameter space where large cross sections are possible. When we calculate the cross section \(\sigma(pp \to h_{SM}) Br(h_{sM} \to \gamma\gamma)\) for a SM Higgs with mass \(m_{h_{SM}} = 130\) GeV we find it to lie at around 40 fb which is therefore comparable to the red contour (the shortest vertical band) which reaches Radion masses of about 220 GeV. A cross section of 10 fb is possible in a much larger region and can actually be reached for Radion masses well above 350 GeV. We should note here that the dependence of the cross sections with the scale \(\Lambda_\phi\) goes as \(\sim 1/\Lambda_\phi^2\), and so for example the numbers in the contours should be roughly divided by 4 if we were to take \(\Lambda_\phi = 4\) TeV and scan for the new allowed regions in the \((m_\phi - \xi)\) plane. In any case the larger the mass, the better the backgrounds are, so one should definitely consider this surprising channel as a new possibility arising from the scalar sector of Randall-Sundrum scenarios. The gray region in both panels is again the theoretically excluded region. The yellow region marked “Tevatron Excluded” makes use of the latest Tevatron Data from the Higgs search [36]. The most important channel for Higgs searches in the mass range 150 – 200 GeV is the process \(p\bar{p} \to h_{SM} \to W^+W^-\) with the two \(W\)’s decaying leptonically. One can then easily convert bounds on this process into bounds
in the Radion parameter space [37], and as is seen in the figure for $\Lambda_\phi = 2$ TeV these bounds do reach our parameter space.

IV. CONCLUSIONS

In this paper we have extended the study of Radion phenomenology with matter fields in the Bulk to the case of Higgs and Radion mixing. From the study of Higgs-Radion mixing in RS1 [22, 24, 27, 31] one might think that very similar results would hold modulo the relative enhancements/suppressions observed when Fields are in the Bulk [32]. This was actually the case for the $pp \rightarrow \phi \rightarrow ZZ$ channel but surprising effects happened in the $\gamma\gamma$ channel. Of course in this channel, both the production and the decay happen via loops and perhaps it is not surprising that presumably benign changes in the underlying model can produce significant phenomenological differences. In particular we pointed out that a non negligible region of parameter space allows for important di-photon signals even in a mass range well above the $WW$ and $ZZ$ thresholds. In this case the main effect is caused not because of a huge enhancement of the Radion coupling to photons (although it is enhanced for large Radion masses) but because of a suppression of the couplings to the SM massive particles. In the context of Higgs-Radion mixing, this signal could then play an important discriminating role between different models of warped extra dimensions.

Other interesting features of the Radion with Bulk Fields pointed out in [32] such as its coupling to fermions depending directly on the 5D bulk mass parameters $c_{L,R}$ might also be enhanced by the effects of the Higgs-Radion mixing and should also be looked at carefully although we leave this for the future.

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[1] L. Randall and R. Sundrum, Phys. Rev. Lett. 83, 3370 (1999) arXiv:hep-ph/9905221.
[2] H. Davoudiasl, J. L. Hewett and T. G. Rizzo, Phys. Lett. B 473, 43 (2000) arXiv:hep-ph/9911262; A. Pomarol, Phys. Lett. B 486, 153 (2000) arXiv:hep-ph/9911294.
[3] S. Chang, J. Hisano, H. Nakano, N. Okada and M. Yamaguchi, Phys. Rev. D 62, 084025 (2000) arXiv:hep-ph/9912498.
[4] C. Csáki, J. Erlich and J. Terning, Phys. Rev. D 66, 064021 (2002) arXiv:hep-ph/0203034; J. L. Hewett, F. J. Petriello and T. G. Rizzo, JHEP 0209, 030 (2002) arXiv:hep-ph/0203091.
[5] K. Agashe, A. Delgado, M. J. May and R. Sundrum, JHEP 0308, 050 (2003) arXiv:hep-ph/0308036.
[6] C. Csáki, C. Grojean, L. Pilo and J. Terning, Phys. Rev. Lett. 92, 101802 (2004) arXiv:hep-ph/0308038.
[7] R. Contino, Y. Nomura and A. Pomarol, Nucl. Phys. B 671, 148 (2003) arXiv:hep-ph/0306250; K. Agashe, R. Contino and A. Pomarol, Nucl. Phys. B 719, 165 (2005) arXiv:hep-ph/0412089.
[8] C. Csáki, C. Grojean, J. Hubisz, Y. Shirman and J. Terning, Phys. Rev. D 70, 015012 (2004) arXiv:hep-ph/0310355.
[9] Y. Grossman and M. Neubert, Phys. Lett. B 474, 361 (2000) arXiv:hep-ph/9912408.
[10] T. Gherghetta and A. Pomarol, Nucl. Phys. B 586, 141 (2000) arXiv:hep-ph/0003129.
[11] S. J. Huber and Q. Shafi, Phys. Lett. B 498, 256 (2001) arXiv:hep-ph/0010195.
[12] G. Burdman, Phys. Rev. D 66, 076003 (2002) arXiv:hep-ph/0205329; G. Burdman, Phys. Lett. B 590, 86 (2004) arXiv:hep-ph/0310144.
[13] S. J. Huber, Nucl. Phys. B 666, 269 (2003) arXiv:hep-ph/0303183.
[14] K. Agashe, G. Perez and A. Soni, Phys. Rev. D 71, 016002 (2005) arXiv:hep-ph/0408134; K. Agashe, M. Papucci, G. Perez and D. Pirjol, arXiv:hep-ph/0509117.
[15] G. Cacciapaglia, C. Csaki, J. Galloway, G. Marandella, J. Terning and A. Weiler, JHEP 0804, 006 (2008) arXiv:0709.1714 [hep-ph]; C. Csaki, A. Falkowski and A. Weiler, JHEP 0809, 008 (2008) arXiv:0804.1954 [hep-ph].
[16] S. Davidson, G. Isidori and S. Uhlig, Phys. Lett. B 663, 73 (2008) arXiv:0711.3376 [hep-ph].

17
[17] W. D. Goldberger and M. B. Wise, Phys. Rev. Lett. 83, 4922 (1999) arXiv:hep-ph/9907447; Phys. Lett. B 475, 275 (2000) arXiv:hep-ph/9911457.

[18] J. Garriga and T. Tanaka, Phys. Rev. Lett. 84, 2778 (2000) arXiv:hep-th/9911055.

[19] C. Charmousis, R. Gregory and V. A. Rubakov, Phys. Rev. D 62, 067505 (2000) arXiv:hep-th/9912160.

[20] C. Csáki, M. Graesser, L. Randall and J. Terning, Phys. Rev. D 62, 045015 (2000) arXiv:hep-ph/9911406.

[21] T. Tanaka and X. Montes, Nucl. Phys. B 582, 259 (2000) arXiv:hep-th/0001092.

[22] C. Csáki, M. L. Graesser and G. D. Kribs, Phys. Rev. D 63, 065002 (2001) arXiv:hep-th/0008151.

[23] R. Hofmann, P. Kanti and M. Pospelov, Phys. Rev. D 63, 124020 (2001) arXiv:hep-ph/0012213; P. Kanti, K. A. Olive and M. Pospelov, Phys. Lett. B 538, 146 (2002) arXiv:hep-ph/0204202.

[24] G. F. Giudice, R. Rattazzi and J. D. Wells, Nucl. Phys. B 595, 250 (2001) arXiv:hep-ph/0002178.

[25] K. m. Cheung, Phys. Rev. D 63, 056007 (2001) arXiv:hep-ph/0009232.

[26] J. L. Hewett and T. G. Rizzo, JHEP 0308, 028 (2003) arXiv:hep-ph/0202155.

[27] D. Dominici, B. Grzadkowski, J. F. Gunion and M. Toharia, Nucl. Phys. B 671, 243 (2003) arXiv:hep-ph/0206192; D. Dominici, B. Grzadkowski, J. F. Gunion and M. Toharia, Acta Phys. Polon. B 33, 2507 (2002) arXiv:hep-ph/0206197.

[28] J. F. Gunion, M. Toharia and J. D. Wells, Phys. Lett. B 585, 295 (2004) arXiv:hep-ph/0311219.

[29] S. Bae, P. Ko, H. S. Lee and J. Lee, arXiv:hep-ph/0103187.

[30] M. Toharia, Mod. Phys. Lett. A 19, 37 (2004) arXiv:hep-th/0212036.

[31] T. G. Rizzo, JHEP 0206, 056 (2002) arXiv:hep-ph/0205242.

[32] C. Csaki, J. Hubisz and S. J. Lee, Phys. Rev. D 76, 125015 (2007) arXiv:0705.3844 [hep-ph].

[33] H. Davoudiasl, J. L. Hewett and T. G. Rizzo, JHEP 0308, 034 (2003) arXiv:hep-ph/0305086.

[34] S. Mrenna and J. D. Wells, Phys. Rev. D 63, 015006 (2001) arXiv:hep-ph/0001226.

[35] R. Barate et al. [LEP Working Group for Higgs boson searches and ALEPH, DELPHI, L3 and OPAL Collaborations], Phys. Lett. B 565, 61 (2003) arXiv:hep-ex/0306033.

[36] [CDF and D0 Collaborations], “Combined CDF and D0 upper limits on standard model physics at Tevatron energy,” arXiv:hep-ex/0306033.
Higgs boson production at high mass (155-200 GeV) with 3.0 fb-1 of data 08/03/08, 3.0 fb-1”,
http://www-d0.fnal.gov/Run2Physics/WWW/results/higgs.htm [D0 Collaboration], “Search
for the Higgs boson in H→WW(*)→ll’nuν (l,l’=e,μ) decays with 3.0 fb-1 at D0 in Run II
07/31/08, 3.0 fb-1”, http://www-d0.fnal.gov/Run2Physics/WWW/results/higgs.htm

[37] N. Okada and M. Toharia, work in progress.