Can the 750 GeV diphoton LHC excess be due to a radion-dominated state?

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Abstract

We discuss the possibility of interpreting the 125 GeV scalar boson and the 750 GeV diphoton excess recently reported by the ATLAS and CMS experiments as a Higgs-dominated and a radion-dominated states respectively in a stabilized brane-world model. It is shown that in the simplest variant of the model, where only the gravitational degrees of freedom propagate in the bulk, the production cross section of the radion-dominated state with mass 750 GeV turns out to be too small in the allowed region of the model parameter space for explaining the nature of the excess in this approach.

1 Introduction

The 750 GeV diphoton excess that has been recently presented by the LHC experiments \[1, 2\] can be the first direct evidence of physics beyond the Standard Model (SM), if it is confirmed by further searches at larger statistics. The most plausible explanation of this excess seems to be the production of a new heavy scalar state \[1\]. There have been already numerous attempts either to explain the excess in a model independent way (see, e.g., \[3\]), or to associate it with the scalar states predicted by concrete SM extensions (see, e.g., \[4, 5\]), in particular, to interpret it as a mixed Higgs-radion state \[6, 7, 8\] in the Randall-Sundrum model \[9\].

In the present paper we consider the case, where the two scalar states are interpreted as mixed Higgs-radion states in an extension of the SM based on the Randall-Sundrum model with two branes stabilized by a bulk scalar field \[10, 11, 21, 13\], which is necessary for the model to be phenomenologically acceptable. A characteristic feature of this extension is the presence of a massive scalar radion field together with its Kaluza-Klein tower. These fields have the same quantum numbers as the neutral Higgs field and can mix with the Higgs field, if they are coupled. It is worth noting that in the model under consideration all the SM fields are located on the TeV brane, which does not lead to any contradiction with the Electroweak Precision Data \[14, 15, 16\].

Usually, the Higgs-radion coupling in the Randall-Sundrum model is obtained by introducing a Higgs-curvature term on the brane \[17\]. In paper \[18\] a model was discussed, where a Higgs-radion coupling arises due to a mechanism of spontaneous symmetry breaking on the brane involving the stabilizing scalar field. This approach takes into account the influence of the KK tower of higher scalar excitations on the parameters of the Higgs-radion mixing, which turns out to be of importance. It also has the advantage that, unlike the approach based on the Higgs-curvature term, it modifies only the scalar sector of the model and leaves intact the

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\(^{1}\)A vector state cannot decay to two photons because of the Landau-Yang theorem. Spin two states can. However, if they are somehow related to the graviton excitations similar excesses should appear in di-lepton and di-jet decay modes, which are not seen in the LHC experiments.
coupling constants and the masses of the graviton KK excitations. In this case the latter are of the order of the fundamental five-dimensional energy scale. In the present paper we take this energy scale to be 2 TeV or larger. Therefore, the masses of the graviton KK excitations should also be larger than 2 TeV and the absence of any excess in the di-lepton mode in the region of 750 GeV is quite consistent with the property of the stabilized brane world models to have the KK graviton states that are much heavier than the radion.

Here we will use the effective interaction Lagrangian obtained in paper [18] to describe the couplings of the mixed Higgs-radion states to the SM fields and to analyze the interpretation of the excess as a 750 GeV mixed Higgs-radion scalar state.

2 Higgs-radion interaction and the effective Lagrangian

First let us briefly recall the main features of the model (details are given in [18]). It is a variant of the stabilized RS model with all the SM fields located on the TeV brane [19, 20, 21] and the stabilizing Goldberger-Wise scalar field propagating in the bulk along with the gravitational field. The mixing of the Higgs and gravitational scalar fields arises in the model due to an interaction Lagrangian leading to a relation between the Higgs vev and the value of the stabilizing field on the UV brane that correspond to the minimum of the potential.

The model in five-dimensional space-time \( E = M_4 \times S^1 / \mathbb{Z}_2 \) with coordinates \( \{ x^M \} \equiv \{ x^\mu, y \} \), \( M = 0, 1, 2, 3, 4 \), \( \mu = 0, 1, 2, 3 \), the coordinate \( x^4 \equiv y \), \( -L \leq y \leq L \) parameterizing the fifth dimension is defined by the action

\[
S = S_g + S_{\phi + SM}, \tag{1}
\]

where \( S_g \) and \( S_{\phi + SM} \) are given by

\[
S_g = -2M^3 \int d^4x \int_{-L}^{L} dy R \sqrt{g}, \tag{2}
\]

\[
S_{\phi + SM} = \int d^4x \int_{-L}^{L} dy \left( \frac{1}{2} g^{MN} \partial_M \phi \partial_N \phi - V(\phi) \right) \sqrt{g} - \int_{y=0}^{y=L} \sqrt{-g} V_1(\phi) d^4x + \int_{y=L}^{y=-L} \sqrt{-g} (-V_2(\phi) + L_{SM-HP} + L_{int}(\phi, H)) d^4x. \tag{3}
\]

The signature of the metric \( g_{MN} \) is chosen to be \((+, -, -, -, -)\), \( M \) is the fundamental five-dimensional energy scale, \( V(\phi) \) is a bulk stabilizing scalar field potential and \( V_{1,2}(\phi) \) are brane scalar field potentials, \( \tilde{g} = \text{det} \tilde{g}_{\mu \nu} \), and \( \tilde{g}_{\mu \nu} \) denotes the metric induced on the branes. The space of extra dimension is the orbifold \( S^1 / \mathbb{Z}_2 \), which is realized as the circle of circumference \( 2L \) with the points \( y \) and \( -y \) identified. Correspondingly, the metric \( g_{MN} \) and the scalar field \( \phi \) satisfy the standard orbifold symmetry conditions and the branes are located at the fixed points of the orbifold, \( y = 0 \) and \( y = L \). The SM fields are assumed to be localized on the brane at \( y = L \), and the Lagrangian \( L_{SM-HP} \) is the SM Lagrangian without the Higgs potential. The key point of the approach under consideration is the replacement of the Higgs potential by the interaction Lagrangian

\[
L_{int}(\phi, H) = -\lambda (|H|^2 - \frac{\xi}{M^2} \phi^2)^2 \tag{4}
\]
of the Higgs and Goldberger-Wise fields, $\xi$ being a positive dimensionless parameter.

The background solutions for the metric and the scalar field, which preserve the Poincaré invariance in any four-dimensional subspace $y = \text{const}$, look like

$$ds^2 = e^{-2A(y)} \eta_{\mu\nu} dx^\mu dx^\nu - dy^2 \equiv \gamma_{MN}(y) dx^M dx^N,$$  \(5\)

$$\phi(x, y) = \phi(y),$$  \(6\)

$\eta_{\mu\nu}$ denoting the flat Minkowski metric, whereas the background (vacuum) solution for the Higgs field is standard

$$H_{\text{vac}} = \left( \frac{0}{\sqrt{2}} \right),$$  \(7\)

all the other SM fields being equal to zero.

If one substitutes this ansatz into the equations corresponding to action (1), one gets a system of differential equations for the functions $A(y)$ and $\phi(y)$, the brane scalar field potentials $V_{1,2}(\phi)$ and interaction Lagrangian (4) defining the boundary conditions for these equations on the branes. The potentials $V_{1,2}(\phi)$ are chosen so that they fix the values of the stabilizing field $\phi$ on the branes and stabilize the interbrane distance [10, 11]. Interaction Lagrangian (4) does not affect this stabilization mechanism, if the relation

$$\phi(L)^2 = \frac{M v^2}{2 \xi}$$  \(8\)

between the values of the Higgs and stabilizing scalar field on the brane at $y = L$ is valid, which defines the vacuum value of the Higgs field. This means that in such a scenario the Higgs field vacuum expectation value, being proportional to the value of the stabilizing scalar field on the TeV brane, arises dynamically as a result of the gravitational bulk stabilization.

Now the linearized theory is obtained by expanding the metric, the scalar and the Higgs field in the unitary gauge around the background solution as

$$g_{MN}(x, y) = \gamma_{MN}(y) + \frac{1}{\sqrt{2} M^3} h_{MN}(x, y),$$  \(9\)

$$\phi(x, y) = \phi(y) + \frac{1}{\sqrt{2} M^3} f(x, y),$$  \(10\)

$$H(x) = \left( \frac{0}{\sqrt{2}} + \sigma(x) \right).$$  \(11\)

After substituting this representation into action (1) and keeping the terms of the second order in $h_{MN}$, $f$ and $\sigma$ one gets the Lagrangian of this action which is the standard free Lagrangian of the SM (i.e. the masses of all the SM fields are expressed in the same way as usually in terms of the vacuum value of the Higgs field and the coupling constants) together with the standard second variation Lagrangian of the stabilized RS model [13] supplemented by an interaction term of the scalar fields $f$ and $\sigma$ on the brane coming from interaction Lagrangian (4) [18].

Besides the fields $f$ and $\sigma$, there two more scalar fields in the linearized theory, – the fields $h_{44}(x, y)$ and $\gamma^\mu\nu h_{\mu\nu}(x, y)$. However, the fields $f(x, y)$, $h_{44}(x, y)$ and $\gamma^\mu\nu h_{\mu\nu}(x, y)$ are not independent: they are connected by the equations of motion of the linearized theory and a
gauge condition \[13, 14\]. For this reason we can use any one of them to describe the scalar states.

The field \( h_{44}(x, y) \) can be expanded in KK modes, the lowest mode \( \phi_1(x) \) is called the radion field and the modes \( \phi_n(x) \), \( n > 1 \) belong to its KK tower. This expansion induces the corresponding expansion of the bulk scalar field \( f(x, y) \). Substituting the latter expansion into the second variation Lagrangian and integrating over the extra dimension coordinate, one gets an effective four-dimensional Lagrangian. In case the Higgs and the radion masses are much lower than the masses of the radion excitations one can pass to a low energy approximation in the four-dimensional Lagrangian by integrating out the radion excitation fields. This gives an effective Lagrangian for the interactions of the Higgs and radion fields with the SM fields. However, due to the Higgs-radion mixing terms the fields \( \sigma(x) \) and \( \phi_1(x) \) are not mass eigenstates.

The physical mass eigenstate fields \( h(x), r(x) \) are, as usually, obtained by a rotation diagonalizing the mass matrix

\[
L_{h-r} = \frac{1}{2} \partial_\mu h(x) \partial^\mu h(x) - \frac{1}{2} m_h^2 h(x)^2 + \frac{1}{2} \partial_\mu r(x) \partial^\mu r(x) - \frac{1}{2} \mu_r^2 r^2(x)
\]

\[-\left(\frac{c \cos \theta + \sin \theta}{\Lambda_r} h(x)(T^\mu_\mu + \Delta T^\mu_\mu) + \frac{\sin \theta - \cos \theta}{\Lambda_r} r(x)(T^\mu_\mu + \Delta T^\mu_\mu) - \right.
\]

\[-\sum_f \frac{m_f}{v} \bar{\psi}_f \psi_f (\cos \theta h(x) - \sin \theta r(x)) + \frac{2M^2}{v} (W^\mu_- W^{\mu+})(\cos \theta h(x) - \sin \theta r(x)) +
\]

\[+ \frac{M^2}{v} (Z^\mu_- Z^{\mu+})(\cos \theta h(x) - \sin \theta r(x))^2 +
\]

\[+ \frac{M^2}{2v^2} (Z^\mu_- Z^{\mu+})(\cos \theta h(x) - \sin \theta r(x))^2.
\]

Here \( m_h^2 \) and \( m_r^2 \) are the masses of the fields \( h(x) \) and \( r(x) \), \( \Lambda_r \) is the (inverse) coupling constant of the radion to the trace of the SM energy-momentum tensor, \( \Delta T^\mu_\mu \) is the conformal anomaly of massless vector fields explicitly given by

\[
\Delta T^\mu_\mu = \frac{\beta(g_s)}{2g_s} G^{ab}_{\rho \sigma} G^{\rho \sigma}_{ab} + \frac{\beta(e)}{2e} F^{\rho \sigma} F^{\rho \sigma}
\]

with \( \beta \) being the well-known QCD and QED \( \beta \)-functions.

The parameter \( c \) accommodates the contributions of the integrated out heavy scalar modes and is expressed in terms of the physical masses and the mixing angle as follows:

\[
c = \frac{(m_r^2 - m_h^2) \sin 2\theta}{m_r^2 \cos^2 \theta + m_h^2 \sin^2 \theta} \left( \sum_{n=2}^{\infty} \alpha_n^2 \right).
\]
It also depends on the sum of the coefficients \( \alpha_n^2 \), where \( \alpha_n \) is the ratio of the wave functions in the extra dimension of the modes \( \phi_n \) and \( \phi_1 \) taken at \( y = L \). These ratios are, of course, model dependent and should fall off with \( n \) in order for the sum to be convergent. In certain models several first ratios may be of the order of unity [13]. Thus, one can conservatively estimate this sum to be also of the order of unity.

The effective four-dimensional interaction Lagrangian (13) expressed in terms of the physical Higgs-dominated \( h(x) \) and radion-dominated \( r(x) \) fields involves only five extra parameters in addition to those of the SM: the masses of the Higgs-dominated and radion-dominated fields \( m_h \) and \( m_r \), the mixing angle \( \theta \), the (inverse) coupling constant of the radion to the trace of the energy-momentum tensor of the SM fields \( \Lambda_r \) and the parameter \( c \).

Let us consider several interesting subspaces of the parameter space of the effective theory. If we put the parameters \( c \) and \( \theta \) equal to zero, i.e. consider the case of the zero mixing, Lagrangian (13) becomes just the SM Higgs Lagrangian plus the usual Lagrangian of the radion interaction with the trace of the SM energy-momentum tensor. If we put \( 1/\Lambda_r = 0 \), we get the effective Lagrangian of the real Higgs singlet extension of the Standard Model [22, 23]. If we formally put the parameter \( c \) equal to zero while keeping the mixing angle \( \theta \) non-zero, we obtain an effective interaction Lagrangian which is very similar to the one of the unstabilized RS model with the Higgs-curvature term on the brane [17]. In particular, the couplings of the Higgs-dominated and the radion-dominated states to the conformal anomaly of massless vector fields, which turn to be very important for their production and decay to two photons, are the same. However, one should keep in mind that the observable parameters of the effective Lagrangian in different cases depend differently on the fundamental parameters of the models.

Thus, Lagrangian (13) is a very general effective Lagrangian of the interaction of two mixed scalar states with the fermion and vector fields of the Standard Model extended by a singlet scalar. The extra parameters of this Lagrangian are natural for considering the phenomenology of the mixed scalar states and allow one to compare easily the predictions of different models. In fact, it is unimportant, how they depend on the fundamental parameters of a particular model as long as the latter belong to the phenomenologically acceptable parameter subspace.

In the general case of a non-zero mixing, when all the parameters and \( \theta, 1/\Lambda_r, \) and \( c \) are not equal to zero, the additional terms in the Lagrangian, coming from the integrated out heavy modes and containing the parameter \( c \) that depends on the scalar state masses and the mixing angle, may lead to certain changes in the collider phenomenology of the Higgs-dominated and radion-dominated states as was demonstrated in [18], where also the Feynman rules for the model are explicitly given.

3 Phenomenology of the 750 GeV radion-dominated state

In order to understand whether or not the 750 GeV observed excess may be interpreted as a radion-dominated state in the above described stabilized brane world model let us consider the main decay and production properties of such a state as follow from effective Lagrangian (13). For the computations a special version of the CompHEP code [24, 25] was used in the same manner as was done in [18]. The version includes a special routine for \( \chi^2 \)-analysis of signal strengths and implements the Feynman rules corresponding to effective Lagrangian (13).

The decay branching ratios for the radion-dominated state with mass 750 GeV are shown
in Fig. 1 as functions of the mixing angle parameter $\sin \theta$. One can see that practically for all the values of the mixing parameter $\sin \theta$ the branching ratios are distributed between the modes close to those for the SM Higgs boson, if it would have had mass 750 GeV. The NLO corrections are included following the HDECAY code [26]. The dominating decay modes are the decays to heavy SM particles $W^+ W^-, ZZ$ boson pairs and the top quark pair. However, for some rather small values of the parameter $\sin \theta$ close to approximately $v/\Lambda_r$, all the branching ratios are significantly decreased, and the dominating decay mode becomes the mode to two gluons. Also in this region of the parameter space the branching to two photons is significantly increased. Such a property could be easily understood from the structure of the interaction vertices of the radion-dominated state and the SM fermions and gauge bosons (see the Feynman rules in [18]). Indeed all the vertices for the fermions and massive gauge bosons contain the factor $\frac{\cos \theta - c \sin \theta}{\Lambda_r} - \frac{\sin \theta}{v}$, which becomes very small for $\sin \theta$ close to $v/\Lambda_r$. This occurs due to the cancellation of the contributions to the vertices coming from the SM type part of the interactions and the part coming from the trace of the energy-momentum tensor. In contrast, the interaction vertices of the radion-dominated state and the massless gluons and photons have anomaly enhanced contributions and the mentioned cancellation does not take place for small values of the parameter $\sin \theta$ close to $v/\Lambda_r$. The corresponding cancellation occurs for the gluon-gluon and photon-photon vertices at much larger values of the parameter $\sin \theta$, where the branching ratios have minima as can be seen in Fig. 1. One should mention that the position of the maximum value for the gluon and photon decay modes goes to smaller and smaller values of $\sin \theta$ with the increase of the scale parameter $\Lambda_r$. The position of the maximum as well as the form of the curves close to the maximum practically does not depend on the value of the parameter $c$, which accumulates the contributions of the higher KK scalar modes. The dependence on the parameter $c$ for two different values of this parameter is demonstrated in Fig. 2 showing the gluon-gluon and photon-photon branchings as functions of $\sin \theta$. 

Figure 1: The decay branching ratios for the radion-dominated state with mass 750 GeV as functions of the mixing angle parameter $\sin \theta$.

Figure 2: The gluon-gluon and photon-photon branchings as functions of $\sin \theta$ for $c = c_{\text{max}}$ and $c = 0.25 c_{\text{max}}$. 
Fig. 3 illustrates the behavior of the total width of the radion-dominated state. The deep minimum corresponds exactly to the discussed region, where all the leading decay modes drastically go down. For a particular choice of the parameters $\Lambda_r = 2$ TeV and $c = c_{\text{max}}$ the width in the minimum becomes as small as $6 \cdot 10^{-2}$ GeV. However, for the values of the parameter $\sin \theta$ away from the minimum region the width can be significantly larger reaching a few tens or even a hundred GeV. In the two-dimensional plot (see Fig. 4) the lines of equal width values are shown depending on the parameters $\sin \theta$ and $1/\Lambda_r$. The plot demonstrates that the radion-dominated state can be rather wide. The width of the 750 GeV excess observed at the LHC has a rather large value of the order of 45 GeV [1, 2].

However, a crucial point for the possible interpretation is the production cross section, which should be in the range from a few to 10 fb. The production cross section for the radion-dominated state is shown in Fig. 5. We have included the contributions of all the production channels for the radion-dominated state (ggF, VBF, rV, rtt) with the decay to two photons. We have included the NNLO K-factors taken from the Higgs cross section working group web page [27, 28]. One can see that, as expected, the gluon-gluon fusion dominates the production cross section in the most interesting region of the parameter space, where the cross section has a maximum. The maximum occurs for the same values of parameters, for which the corresponding gg and $\gamma \gamma$ decay branching ratios have the maximum. As for the branching ratios in the range close to the maximum, the cross section depends very weakly on the the parameter $c$. Such a dependence is significant for rather large mixing angles, where the cross section becomes much smaller (see Fig. 6).

One should note that the parameter region allowed by the signal strength measurements for 125 GeV boson at 8 TeV LHC energy and presented in Fig. 7 includes the above mentioned region, where the two-photon cross section has a maximum. The inclusion of the new measurements at 13 TeV for the 125 GeV boson and the 750 GeV possible resonance into the
Figure 5: The production cross section of the radion-dominated state with mass 750 GeV as a function of the mixing angle parameter $\sin \theta$ including the contributions of all the production modes (thick curve) and only the leading contribution of the gluon-gluon fusion mode (thin curve).

Figure 6: The production cross section of the radion-dominated state in the main gluon-gluon fusion mode with mass 750 GeV as a function of the mixing angle parameter $\sin \theta$ for the parameter $c = c_{\text{max}}$ (thick curve) and $c = 0.25 c_{\text{max}}$ (thin curve).

Figure 7: Exclusion contours for the global $\chi^2$ fit in the $(\sin \theta, 1/\Lambda_r)$ plane for the LHC at $\sqrt{s} = 7$ and 8 TeV and $m_h = 125$ GeV, $m_r = 750$ GeV, $c = c_{\text{max}}$. The dark, medium and light shaded areas correspond to CL of the fit 65%, 90% and 99% respectively.

Figure 8: Exclusion contours for the global $\chi^2$ fit in the $(\sin \theta, 1/\Lambda_r)$ plane for the LHC at $\sqrt{s} = 7$, 8 and 13 TeV and $m_h = 125$ GeV, $m_r = 750$ GeV, $c = c_{\text{max}}$. The dark, medium and light shaded areas correspond to CL of the fit 65%, 90% and 99% respectively.

overall fit leads to a minor modification of the allowed region of the model parameter space as demonstrated in Fig. 8.
However, the interpretation of the observed excess as the radion-dominated state is very problematic or even impossible in the simplest variant of the discussed brane-world models, where only the gravitational degrees of freedom are allowed to propagate in the bulk. Indeed, as one can see in Fig. 6, the cross section has a maximum of about 0.14 fb, which is by a factor of $50\div100$ smaller than what is needed to achieve the observed level of the cross section for the 750 GeV excess. In other areas of the parameter space the production cross section gets even smaller as one can see in Fig. 9, where equal value contours for the cross section are shown.

Thus, one has to conclude that in the simplest variant of the model under consideration, where only the gravitational and stabilizing scalar fields propagate in the bulk, the observed excess cannot be understood as a radion-dominated resonance. This is in accord with the results of paper [29], where the production and two-photon decay of a heavy radion-dominated state was considered in the unstabilized RS model with the Higgs-curvature term and the SM fields on the brane. As we have mentioned in Section 2, our effective Lagrangian (13) for small values of the parameter $c$ gives the same coupling of the radion-dominated state to the conformal anomaly of massless vector fields as in paper [29], which turns out to be dominating for the small values of the mixing angle, and the corresponding contact diagrams give the main contribution to the amplitude.

A similar conclusion about the impossibility to explain the 750 GeV excess by a mixed Higgs-radion state was reached in papers [5, 8]. In the present paper, in addition to these previous studies, we analyzed the influence of the tower of the higher KK scalar states having performed the statistical analysis of the allowed parameter space and having shown that the inclusion of these KK states does not change the main negative conclusion. As was explained above, the obtained effective Lagrangian of the model is rather general and therefore the conclusion remains valid for several other extensions of the SM by a singlet real scalar field.
In order to increase the two-photon signal rate some other heavy particles should propagate in the gluon-gluon-scalar and photon-photon-scalar loop vertices, the scalar in our case being the Higgs-dominated or the radion-dominated states. Such heavy particles in brane-world models could be the excited KK states of those SM fields, which are also allowed to propagate in the multidimensional bulk. The contribution of the bulk field excitations to the production cross section of a heavy radion-dominated state with mass up to 350 GeV was discussed in paper [29]. It was shown there that it can increase the cross section up to 10 fb, although it is not quite clear, whether this result can be extrapolated to the 750 GeV radion-dominated state. In fact, in paper [7] it is claimed that the bulk field excitations in the Randall-Sundrum background can really give the necessary enhancement of the production cross section of this state. However, the calculation of the excited state contributions in stabilized brane-world models is not straightforward, in particular because of the excited state impact on the energy momentum tensor, and needs a special thorough investigation.

4 Conclusions

In the present paper we have studied the possibility of interpreting the 750 GeV diphoton excess that has been recently presented by the LHC as the production of a radion-dominated state in a stabilized brane-world model. The Higgs-dominated state is fixed to be the 125 GeV boson. For our investigation we used a very general effective Lagrangian that describes the low-energy interactions of the scalar mixed states in the Standard model extended by a scalar singlet including the remaining contributions from the integrated out possible tower of heavier scalar states. Our calculations show that although one can rather well fit the signal strengths in a certain region of the model parameter space, the resulting two-photon production cross section is $50 \div 100$ smaller than what is needed to achieve the observed level of the cross section. This means that the stabilized brane-world models with only gravity and the Goldberger-Wise field propagating in the bulk, as well as the other models described by special cases of the considered effective Lagrangian, are unable to explain the excess.

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