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Bulk Higgs with a heavy diphoton signal

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We consider scenarios of warped extra dimensions with all matter fields in the bulk and in which both the hierarchy and the flavor puzzles of the Standard Model are addressed. Inspired by the puzzling excess of diphoton events at 750 GeV reported in the early LHC Run II data (since then understood as a statistical excess), we consider here the general question as to whether the simplest extra-dimensional extension of the Standard Model Higgs sector, i.e., a five-dimensional bulk Higgs doublet, can lead to an intermediate mass resonance (between 500 GeV and 1.5 TeV) of which the first signature would be the presence of diphoton events. This surprising phenomenology can happen if the resonance is the lightest CP-odd state coming from the Higgs sector. No new matter content is required, the only new ingredient being the presence of (positive) brane localized kinetic terms associated to the five-dimensional bulk Higgs (which reduce the mass of the CP-odd states). Production and decay of this resonance can naturally give rise to observable diphoton signals, keeping dijet production under control, with very low ZZ and WW signals and with a highly reduced top pair production in an important region of parameter space. We use the 750 GeV excess as an example case scenario.

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I. INTRODUCTION

The original motivation for warped extra dimensions was to address the hierarchy problem, so that the fundamental scale of gravity is exponentially reduced along the extra dimension, from the Planck mass scale to the TeV scale. Thus, the TeV scale becomes the natural scale of the Higgs sector if this one is localized near the TeV boundary of the extra dimension, as first introduced by Randall and Sundrum (RS) [1]. If Standard Model (SM) fields are allowed to propagate in the extra dimension [2], the scenario can also address the flavor puzzle of the SM, explaining fermion masses and mixings from the geographical location of fields along the extra dimension. However, processes mediated by the heavy resonances of the five-dimensional (5D) bulk fields, Kaluza-Klein (KK) modes, generate dangerous contributions to electroweak and flavor observables (including dangerous deviations to the $Zb\bar{b}$ coupling) [3–5], pushing the KK mass scale to 5–10 TeV [6]. A popular mechanism to lower the KK scale involves using a custodial gauge $SU(2)_R$ symmetry [7], which ensures a small contribution to electroweak precision parameters, lowering the KK scale bound to about 3 TeV.

Alternatively, one can study scenarios in which the metric is slightly modified from the RS metric background (AdS_5). This can be achieved quite naturally from the backreaction on the metric caused by a 5D scalar field stabilizing the original AdS_5 warped background [8]. When the 5D Higgs field is sufficiently leaking into the bulk and when the metric background is modified near the TeV boundary, the scenario allows for KK scales as low as 1–2 TeV, with precision electroweak and flavor constraints under control [9]. An inconvenience is that these scenarios are typically hard to probe experimentally as the couplings of all particles are very suppressed [10–13]. Still, it has been shown that it can still lead to interesting deviations in Higgs phenomenology, as the Higgs couplings can receive sufficient radiative corrections from the many KK fermions of the model [14]. This is the scenario we want to pursue further.

In a nutshell, the modified warped scenario that we use here has the attractive features that precision electroweak and flavor constraints are kept in check for lower KK scales. In the same region of the parameter space, Higgs production and decay are consistent with the SM expectations. The drawback is that this scenario could be challenging to observe at colliders, due to suppressed particle couplings. The stabilization of the geometry of this model has also been established in Ref. [9], by considering a version of the Goldberger-Wise mechanism [8]. The challenges lie in finding some distinguishing signatures of the model.

In particular, we turn our attention here to the possibility of observing a heavier scalar resonance in the very clean diphoton channel at the LHC. Our work is in part motivated by the (now defunct) 750 GeV state [15,16] but also by the prospect of detecting resonances at the LHC with this

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exotic signature in the high luminosity run. We propose a simple scenario within warped extra-dimensional models that contains a heavier resonance with the unexpected signature of an enhanced diphoton rate over other more probable signatures (see also Ref. [17] for another possible candidate in these same scenarios). It would require the presence of a 5D bulk Higgs, with the Higgs as delocalized as possible from the TeV brane (but still close enough to address the hierarchy problem). The reason for localizing the Higgs in the bulk is that the masses of the Higgs KK excitations will increase as the Higgs is pushed toward the brane, getting infinitely heavy in the limit of a brane Higgs. Of these excitations some are CP-odd scalars, making them a natural candidate for a mainly diphoton signal since they do not couple at tree level to ZZ or WW. We show that if the typical mass of the KK gluon (typically the lightest and most visible KK particle) is around 1–2 TeV, it is simple to obtain a lighter CP-odd Higgs with the help of small (and positive) brane localized kinetic terms of the 5D Higgs. Due to suppressed couplings to ZZ and WW (loop level), the CP-odd scalars should have its largest couplings with top pairs. As we will show, this coupling can be naturally small in an important region of the allowed parameter space. This way, the radiative coupling to gluons, large enough for producing CP-odd scalars at the LHC, could also dominate the decays, and (also) the radiative decay into photons could then receive a sufficiently large branching fraction.

Explanations of the (now defunct) 750 GeV diphoton signal within warped scenarios have been put forward previously, with the resonance interpreted as a radion [18] (and/or dilaton [19]), as a KK graviton [20–22], a 5D field-related axion [23], or as an additional 5D singlet scalar added to the model [24]. The scenario proposed here, while preserving minimality, would predict a strong diphoton excess in a significant region of the parameter space.

We proceed as follows. In Sec. II, we describe briefly the warped scenario, followed by its Higgs and gauge sector in Sec. III and the CP-odd sector in more detail in Sec. IV. Within that section, we look at the fermion couplings in Sec. IV A, the $\gamma\gamma$ and $glu-glu$ couplings in Sec. IV B, and the Zh couplings in Sec. IV C. Our numerical estimates are presented in Sec. IV D, and we conclude in Sec. V. We leave some of the details for the Appendix.

II. BACKGROUND METRIC

The (stable) static spacetime background is

$$ds^2 = e^{-2\sigma(y)}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2, \tag{1}$$

where the extra coordinate y ranges between the two boundaries at $y=0$ and $y=y_1$ and where $\sigma(y)$ is the warp factor responsible for exponentially suppressing mass scales at different slices of the extra dimension. In the original RS scenario, $\sigma(y)=ky$, with k the curvature scale of the AdS_5 interval that we take of the same order as M_{Pl}. Nevertheless, this configuration is not stable as it contains a massless radion, a result of having the length of the interval not fixed. In more general warped scenarios with a stabilization mechanism, $\sigma(y)$ is a more general (growing) function of y.

We consider here the specific case where a 5D bulk stabilizer field backreacts on the AdS_5 metric producing the warp factor [9,10]

$$\sigma(y) = ky - \frac{1}{\nu^2}\log\left(1-\frac{y}{y_s}\right), \tag{2}$$

where $y=y_s$ is the position of a metric singularity, which stays beyond the physical interval considered here, i.e., $y_s > y_1$. In these modified metric scenarios, the Planck-TeV hierarchy is reproduced with a shorter extra-dimensional length due to a stronger warping near the TeV boundary, so that, whereas in RS we have $ky_1 \simeq 35$, in the modified scenarios, we can have $ky_1 \simeq 20-30$. The appeal of this particular modification lies on the possibility of allowing for light KK particles (∼1 TeV), thus allowing for their observation at the present LHC run II, while keeping flavor and precision electroweak bounds at bay. This happens when the Higgs profile leaks sufficiently out of the TeV brane so that all of its couplings to KK particles are suppressed compared to the usual RS scenario [9–11]. We thus fix the Higgs localization to a point where it is maximally pushed away from the IR brane, while still solving the hierarchy problem (i.e., making sure that we are not reintroducing a new fine-tuning of parameters within the Higgs potential parameters [9,12].)

III. GAUGE AND HIGGS SECTOR

The matter content of the model is that of a minimal 5D extension of the Standard Model, so that we assume the usual strong and electroweak gauge groups $SU(3)_c\times SU(2)_L\times U(1)_Y$, with all fields propagating in the bulk. The fermions of the model are also bulk fields, with different 5D bulk masses, so that their zero-mode wave functions are localized at different sides of the interval. This way the scenario also addresses the flavor puzzle of the SM, since hierarchical masses and small mixing angles for the SM fermions become a generic feature due to fermion localization and small wave function overlaps [25].

In the electroweak and Higgs sector, we consider the following action,

$$S = \int d^4x dy\sqrt{g}\left(-\frac{1}{4}F_{MN}^2 + |D^M H|^2 - V(H)\right) + \sum_{i=1}^{2}\int d^4x dy\sqrt{g}\delta(y-y_i)\left(\frac{d_i}{k}|D^M H|^2 - \lambda_i(H)\right), \tag{3}$$

where the capital index M will be used to denote the five spacetime directions, while the greek index μ will be used for the four-dimensional (4D) directions. Note that we have introduced brane-localized kinetic terms associated to the 5D Higgs, which are proportional to $\delta(y - y_i)$. This type of terms has been studied previously in the context of flat extra-dimensions [26] as well as warped extra dimensions [27,28]. In the warped case, the brane kinetic terms considered in Refs. [27,28] were associated to bulk gauge fields as well as to the gravitational metric perturbations and fermion fields [28]. We introduce them here in the particular case of a bulk Higgs, which to our knowledge has not been presented before. The effects of such terms are well established: they lead to deviations in the predicted spectrum of the associated KK modes as well as to changes in the strength of their interactions with other fields. We will study here their effects and viability in lowering the CP-odd Higgs mass and include an explicit derivation in the Appendix. The brane kinetic coefficients d_i (in units of k) are essentially free parameters encoding the size of brane localized kinetic terms associated with the bulk Higgs field. These terms will allow for a slight modification of the spectrum of the KK Higgs excitations, particularly useful in lowering the CP-odd masses to intermediate values, say between 500 GeV to 1.5 TeV. These brane kinetic terms can be thought of as exactly localized operators, or as bulk operators that happen to be dynamically localized due to couplings to some localizer vacuum expectation value (VEV).[1] Intuitively, their effect on masses can be understood by the fact that they modify the overall kinetic terms of the fields, forcing a canonical redefinition, which in turn modifies the quadratic terms in the 4D effective action (the mass terms).

The 5D Higgs doublet is expanded around a nontrivial VEV profile $v(y)$ as

$$H = \frac{1}{\sqrt{2}} e^{ig_5 \Pi} \begin{pmatrix} 0 \\ v(y) + h(x,y) \end{pmatrix}. \tag{4}$$

The covariant derivative is $D_M = \partial_M + i g_5 A_M$ with

$$A_M = \begin{pmatrix} s_W A_M^{em} + \frac{c_W^2 - s_W^2}{2c_W} Z_M & \frac{1}{\sqrt{2}} W_M^+ \\ \frac{1}{\sqrt{2}} W_M^- & -\frac{1}{2c_W} Z_M \end{pmatrix}. \tag{5}$$

The CP-odd and charged Higgs part is

$$\Pi = \begin{pmatrix} \frac{c_W^2 - s_W^2}{2c_W} \Pi_z & \frac{1}{\sqrt{2}} \Pi^+ \\ \frac{1}{\sqrt{2}} \Pi^- & -\frac{1}{2c_W} \Pi_z \end{pmatrix} \tag{6}$$

with the weak angle defined like in the SM, i.e., $s_W/c_W = g_5'/g_5$, where g_5 and g_5' are the 5D coupling constants of $SU(2)_L$ and $U(1)_Y$.

The extraction of degrees of freedom in this context has been performed in Refs. [11,12,29], and we outline here the main results, while the effect of brane kinetic terms in the Higgs sector is new, and its derivation is outlined in the Appendix. The 5D equations of motion for all these fields are coupled [except for the case of the real Higgs excitation $h(x,y)$], and in order to decouple them, one can partially fix the gauge, or add a gauge fixing term to the previous 5D action. For example, in the CP-odd case, the fields $Z_\mu(x,y)$, $Z_5(x,y)$, and $\Pi_z(x,y)$ must be unmixed. The partial gauge fixing constraint[2]

$$\partial^\mu Z_\mu - M_z^2(y)\Pi_z + (e^{-2\sigma} Z_5)' = 0 \tag{7}$$

manages to decouple the fields Z_μ from Z_5 and Π_z in the bulk. We defined here $M_z(y) = \frac{g_5}{2c_W} v(y) e^{-\sigma(y)}$.

However, the presence of the Higgs brane kinetic terms, proportional to d_i in the action, forces us to extend the gauge choice on the branes, producing a lifting of the Z_5 field so that the decoupling is maintained at the boundaries.[3] The appropriate boundary condition at the IR brane is

$$Z_5(x, y_1) = -\frac{d_1}{k} M_z^2(y_1) e^{2\sigma(y_1)} \Pi_z(x, y_1), \tag{8}$$

where y_1 denotes the position of the boundary (note that if the brane kinetic term parameter d_1 tends to zero, the condition on Z_5 becomes Dirichlet, as expected). With this type of gauge choice, the 5D fields Z_μ, W_μ, and A_μ have independent 5D equations of motion. In order to extract the effective 4D degrees of freedom, we expand the gauge fields as $Z_\mu(x,y) = Z_\mu^n(x) f_z^n(y)$, $W_\mu(x,y) = W_\mu^n(x) f_w^n(y)$, and $A_\mu(x,y) = A_\mu^n(x) f_\gamma^n(y)$ (summation over n is understood) and where $Z_\mu^0(x)$, $W_\mu^0(x)$, and $A_\mu^0(x)$ are the Z, W, and γ gauge bosons of the SM. The extra-dimensional profiles $f_z^n(y)$, $f_w^n(y)$, and $f_\gamma^n(y)$ are solutions of

$$(e^{-2\sigma} f_a')' + (m_n^2 - M_a^2(y)) f_a = 0, \tag{9}$$

where $a = z,\ w,\ \gamma$, $M_z(y) = \frac{g_5}{2c_W} v(y) e^{-\sigma(y)}$, as defined before, $M_w(y) = \frac{g_5}{2} v(y) e^{-\sigma(y)}$, and $M_\gamma = 0$. The boundary conditions for these profiles are[4]

[1]In order to avoid tachyons and/or ghosts, the sign of the purely brane localized kinetic terms will be kept positive, i.e., $d_i > 0$.

[2]There is still some gauge freedom left, so that the towers of 4D Goldstone bosons that appear can be gauged away.

[3]In the absence of brane kinetic terms, Z_5 must have vanishing boundary conditions (Dirichlet) if Z_μ is to have Neumann conditions and thus develop a zero-mode KK excitation in the effective 4D theory.

[4]We ignore here possible brane localized gauge kinetic terms and keep only the effects from Higgs brane kinetic terms. We include everything in the derivation outlined in the Appendix.

$$\frac{d_i}{k} M_a^2 f_a(y_i) = -e^{-2\sigma} f'_a(y_i). \quad (10)$$

The CP-even Higgs field is expanded as $h(x,y) = h^n(x) h_y^n(y)$, and the equations for the Higgs profiles are, with $h_y \equiv h_y^n(y)$,

$$e^{4\sigma}(e^{-4\sigma} h'_y)' + (m_{h_n}^2 e^{2\sigma} - \mu_{\rm bulk}^2) h_y = 0, \quad (11)$$

where $\mu_{\rm bulk}^2 = \frac{\partial^2 V}{\partial H^2}|_{H=v}$. The boundary conditions are

$$\left(\mu_{{\rm brane}_i}^2 - \frac{d_i}{k} m_{h_n}^2 e^{2\sigma}\right) h_y = -h'_y, \quad (12)$$

with $\mu_{{\rm brane}_i}^2 = \frac{\partial^2 \lambda_i}{\partial H^2}|_{H=v}$. Note that the CP-even Higgs modes obey an equation of motion which involves both brane potentials and brane kinetic terms at the boundaries. The $n = 0$ KK mode of this equation, $h^0(x)$, must be very light as it is the SM Higgs boson. The next state (the $n = 1$ KK CP-even component of the Higgs tower) must be heavier, at the KK scale, as it has additional gradient energy.

There are still some degrees of freedom left, and their 5D equations of motion still happen to be mixed. One of the coupled systems involves Z_5 and Π_z, and the other coupled system involves $\Pi^\pm$ and $W_5^\pm$. In order to disentangle these systems, one must perform a mixed expansion, so that the decoupling of fields will happen KK level by KK level. The mixed expansions are, in the CP-odd sector,

$$Z_5(x,y) = G^n(x) \frac{f'_{G_n}(y)}{m_{G_n}^2} + \Pi_n(x) \frac{e^{2\sigma}}{m_{\pi_n}^2} X_\pi(y) \quad (13)$$

$$\Pi_z(x,y) = G^n(x) \frac{f_{G_n}(y)}{m_{G_n}^2} + \Pi_n(x) \frac{1}{m_{\pi_n}^2 M_z^2} X'_\pi(y), \quad (14)$$

and in the charged scalar sector, they are

$$W_5^\pm(x,y) = G_n^\pm(x) \frac{f_{G_n^\pm}(y)}{m_{G_n^\pm}^2} + \Pi_n^\pm(x) \frac{e^{2\sigma}}{m_{\pi_n^\pm}^2} X_\pm(y) \quad (15)$$

$$\Pi^\pm(x,y) = G_n^\pm(x) \frac{f_{G_n^\pm}(y)}{m_{G_n^\pm}^2} + \Pi_n^\pm(x) \frac{1}{m_{\pi_n^\pm}^2 M_w^2} X'_\pm(y), \quad (16)$$

where $M_z(y)$ and $M_w(y)$ were defined below Eq. (9).

The effective 4D physical fields are the tower of CP-odd neutral scalars $\Pi_n(x)$ and the tower of charged scalars $\Pi_n^\pm(x)$. Their associated extra-dimensional profiles $X_\pi(y)$ and $X_\pm(y)$ obey the equations

$$\left(\frac{1}{M_a^2(y)} X'_a\right)' + \left(\frac{m_{\pi_a}^2}{M_a^2(y)} - 1\right) e^{2\sigma} X_a = 0, \quad (17)$$

where $M_a(y) = (M_z(y), M_w(y))$ and $X_a = (X_\pi, X_\pm)$. The boundary conditions are

$$\frac{d_i}{k} X'_a = -X_a, \quad (18)$$

and note that vanishing Higgs brane kinetic terms imply Dirichlet boundary conditions for X_a. Unlike for the CP-even boson, this condition depends on the brane kinetic terms only and involves no Higgs brane potential terms. In the presence of the brane kinetic terms, this equation has mixed boundary conditions, and as such, we expect the lowest state ($n = 0$) to lie between the lowest Neumann state and the lowest Dirichlet state. In Table I, Sec. IV D, we show numerical results for the lowest-level KK scalar masses in two parameter space regions. We find that the lightest CP-odd mass is indeed in between the lightest CP-even (the SM Higgs) and the first excited CP-even state. We also checked that these bulk equations agree with Refs. [11,12,29], the only new addition being the boundary conditions imposed by the presence of Higgs brane kinetic terms.

In order for these 4D scalars to be canonically normalized, we require

$$\frac{1}{m_a^2} \int dy e^{2\sigma} \frac{X_a^2}{M_a^2} = 1, \quad (19)$$

and this condition includes the effect of Higgs brane kinetic terms.

The remaining 4D fields are $G_n(x)$ and $G_n^\pm(x)$, which are Goldstone bosons at each KK level. The profile wave functions $f_{G_a}(y)$ obey the same differential equations as the gauge profiles, Eq. (9), as well as the same boundary conditions, Eq. (10). The spectrum is thus identical to the gauge bosons spectrum level by level. These fields appear in the effective 4D action coupled to $(\partial^\mu Z_\mu^n)$ or $(\partial^\mu W_\mu^n)$, and of course there is a leftover gauge freedom allowing us to gauge them away (i.e., they are pure gauge).

We wish to identify the lightest CP-odd scalar $\Pi_0(x)$ with a possible diphoton peak at the LHC in the intermediate mass region. In order to have an idea of the effects of the Higgs brane kinetic terms on the CP-odd scalar spectrum, we consider two different parameter points. The first one is one where the background metric is essentially the RS metric. In that case, we take $\nu = 10$ and $y_s = 4 \times y_1$, where ν is the exponent appearing in the modified metric. If ν is relatively large, the location of the spurious singularity is sent away from the boundary, recovering essentially the AdS_5 metric. The other case considered is the situation where the metric modification allows for TeV size KK masses that are safe from precision electroweak constrains. The parameters chosen there are $\nu = 0.5$ and $y_s = 1.04 \times y_1$. In both parameter points, we fix the KK mass of the first gluon KK excitation to be 1500 GeV[5]

[5]Of course, the RS point is presented for comparison only, since such light KK masses should produce too large deviations in the precision electroweak observables.

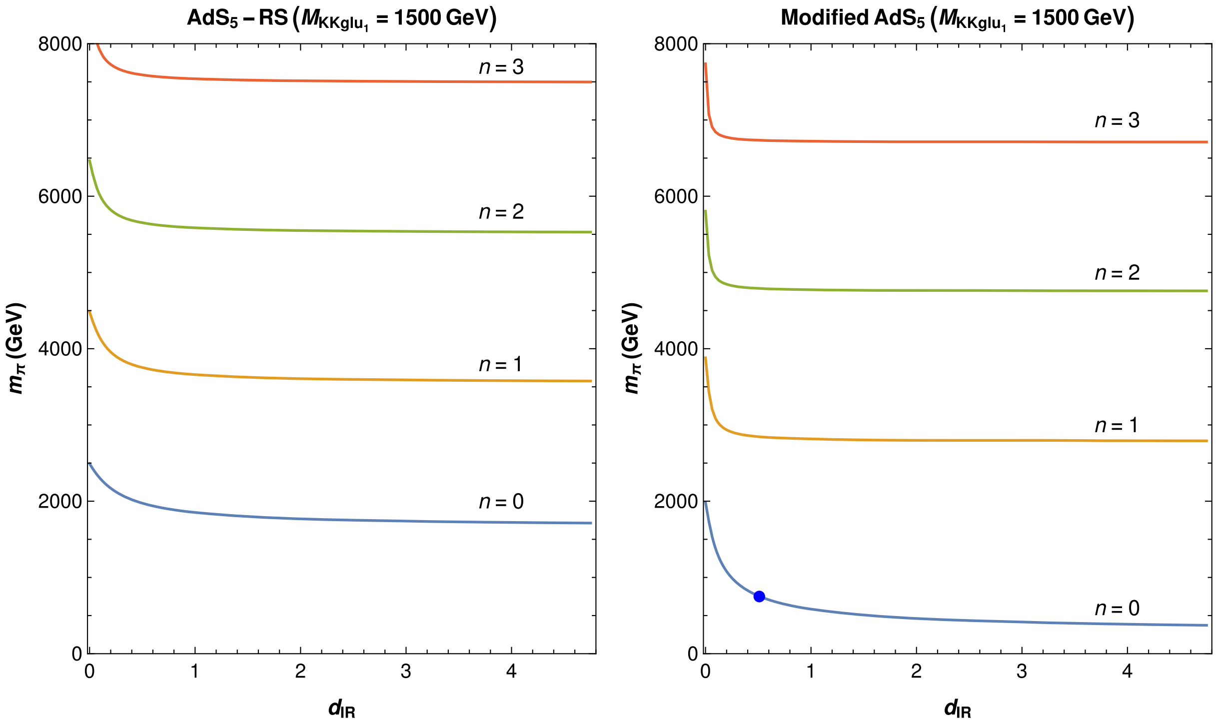

FIG. 1. Mass spectrum of the first CP-odd Higgs scalars as a function of the brane kinetic term coefficient d_1 in the RS metric limit ($\nu = 10$ and $y_s = 4y_1$, left panel) and within a noticeably IR-modified AdS_5 metric ($\nu = 0.5$ and $y_s = 1.04\ y_1$, right panel). In both cases, the first KK gluon mass is fixed at 1500 GeV. Within the modified metric, a brane kinetic coefficient $d_1 \simeq 0.51$ can produce an intermediate mass of 750 GeV, shown here as a dot, while in RS the lightest CP-odd scalar would be much heavier (closer to 2 TeV).

In Fig. 1, we show the spectrum of the first four KK levels of CP-odd Higgs bosons $\Pi_n(x)$, for $n = 0, 1, 2, 3$, as a function of the brane kinetic term d_1 in units of the curvature $k \sim M_{\mathrm{Pl}}$. The effects of the UV localized brane kinetic term are highly suppressed (warped down), and so we do not consider them here anymore. We can see that in the model with modified metric the mass of the CP-odd scalar can be as light as 500 GeV, or even lighter. In the RS limit, the CP-odd scalar mass can be lowered from about 2500 GeV (without brane kinetic terms) to about 1750 GeV for very large brane kinetic terms. All these masses are obtained with a fixed KK gluon mass of 1500 GeV, and so with the modified metric, it becomes possible to have most of the KK resonances of the scenario sitting at 1.5 TeV or more, with a much lighter CP-odd scalar with intermediate mass between 500 to 1500 GeV, generated with relatively small brane kinetic term coefficients. When the metric modification lies between the two limits considered here, there will be an intermediate behavior, with a lightest CP-odd mass having increased asymptotic values as one moves toward a RS-like background. At Run I at the LHC with $\sqrt{s} = 8$ TeV and 19.7 fb^{-1}, the lower bounds on the RS KK gluon mass are found to be about 2.5 TeV [30]. This limit is not as stringent in the modified metric scenario, since KK particle couplings are suppressed compared to the RS scenario. In Ref. [13], the authors evaluated the cross section of $pp \to KKglu_1 \to t\bar{t}$ within the modified metric model for a KK gluon mass of 2.4 TeV, in the same parameter space region considered here. The resulting cross section was about five to ten times smaller than the RS cross section for the same process (depending on the precise parameter points in RS and the modified scenario). We roughly extrapolate this result, so that a mass of $M_{KKglu_1} \sim 1$–1.5 TeV may still be safe in our scenario. In any case, the precise value of the KK gluon mass can be moderately increased without affecting our main conclusions, although larger Higgs brane kinetic terms would be required to decrease the CP-odd scalar mass.

Finally, also note that the spectrum for the charged scalars is essentially the same as the CP-odd scalars, since their differential equations and boundary conditions are identical except for the functions $M_z(y)$ and $M_w(y)$, which differ by about 10%. The deviation with respect to the CP-odd scalar spectrum is less than 5%. This means that the scenario under consideration would also contain a lightest charged Higgs scalar with a mass about the same as the CP-odd scalar.

The next question to ask is how big the effect of the Higgs brane kinetic term on the gauge bosons is, in particular on the lowest ones, i.e., the SM W and Z bosons. These terms represent an additional (brane localized) contribution to the mass of the gauge bosons. In principle, their mass is generated here from a bulk Higgs mechanism, unless the brane kinetic terms are overly important (not the limit we are working with here). We can quickly estimate its effect on the lowest lying gauge fields. These are essentially flat (like all gauge zero modes), and thus their wave function is $f_z^0 \sim 1/\sqrt{y_1}$. The contribution of a brane localized mass squared term is $\delta m_z^2 \lesssim d_1 v^2/y_1 \sim d_1 \times 700\ \mathrm{GeV}^2$. For IR brane kinetic term coefficient d_1 of $\mathcal{O}(1)$, this represents naively at most some 10% contribution to the overall mass squared of either W or Z. In the particular case of the modified metric, for our parameter choice with a brane coefficient $d_1 \simeq 0.51$ and metric parameters $\nu = 0.5$ and $y_s = 1.04 \times y_1$, the exact numerical effect on the zero-mode gauge boson masses (W and Z) is a shift of 3 GeV with respect to the no-brane-kinetic-term limit. Of course, in the presence of brane kinetic terms, one redefines the VEV normalization constant, and the value of g_5, in order to correctly account for the SM gauge boson masses and electroweak couplings. Note that we consider the dimensionless brane kinetic parameter d_1 to be a free parameter as long as it is not hierarchical. In the effective theory considered here, brane localized kinetic operators can be generated from interactions with brane localized matter, via radiative corrections or after some brane localized fields acquire vacuum expectation values. The specific size of the brane kinetic parameter depends on the unknown UV completion of the theory. As it is common in the literature dealing with such terms [26–28], we will allow the parameter d_1 to have order 1 values, with the only theoretical constraint being that it remains strictly positive in order to avoid tachyons and/or ghosts. Of course, the resulting spectrum must remain consistent with precision bounds and with the new mass limits coming from LHC searches. With these constraints in mind, we pursue the study including these brane terms since they lower significantly the pseudoscalar mass, thus yielding new and interesting phenomenology.

IV. CP-ODD HIGGS COUPLINGS

Now that an intermediate mass CP-odd Higgs scalar $\Pi_0(x)$ is allowed in the spectrum, thanks to the effect of small brane localized Higgs kinetic terms, we will study its couplings to SM particles within the modified AdS_5 metric scenario in order to see if it could possibly explain a diphoton excess at the LHC. Of course being a CP-odd scalar, its tree-level couplings to ZZ and WW are zero, making it an ideal candidate for exotic events. We thus need to focus on its tree-level couplings to fermions (and top quark pairs in particular), to Zh (where h is the 125 GeV Higgs), and to its radiative couplings to photons and gluons. We study these in the subsequent subsections.

A. Fermion couplings

The couplings of $\Pi_0(x)$ to fermions arise from two sources in the action. The first source comes from the 5D Higgs Yukawa couplings, and the second comes from the gauge-fermion couplings. This is because the physical field $\Pi_0(x)$ contains some of CP-odd Higgs scalar and some of Z_5 excitation, where Z_5 is the fifth component of the 5D vector boson Z_M. However, the 5D Yukawa coupling allows for direct coupling of $\Pi_0(x)$ to two zero-mode fermions, whereas the gauge-fermion coupling allows only couplings between fermion zero modes and higher KK fermion levels. As we will see, it is important to keep both couplings, since after electroweak symmetry breaking the physical SM fermions (top quarks in particular) are mostly zero modes but also contain a small amount of higher KK excitations and could thus inherit some of the original gauge-fermion coupling, especially if the tree-level Yukawa coupling between $\Pi_0(x)$ and zero-mode top quarks is suppressed (as it can be).

The relevant terms in the action are the 5D Higgs Yukawa couplings and the fermion gauge interaction term,

$$S_{ff\Pi_0} \subset \int d^4x dy \sqrt{g}[Y_u HQU + Y_d HQD + \mathrm{H.c.} + \overline{Q}\not{D}Q + \overline{U}\not{D}U + \overline{D}\not{D}D], \tag{20}$$

where Q, U, D represent the 5D fermion $SU(2)_L$ doublets and up-type and down-type singlets (with generation indices and isospin indices suppressed). The kinetic terms contain the 5D covariant derivative, and from them, we extract the terms containing the CP-odd component $Z_5(x,y)$, and from the Higgs Yukawa couplings, we extract the terms containing the CP-odd Higgs component $\Pi_z(x,y)$.

We follow the approach of Refs. [14,31] and compute these couplings in the modified AdS_5 metric by considering only the effects of three full KK levels, i.e., computing 21×21 fermion Yukawa coupling matrices (with three *up* and three down families, each containing zero modes and three KK levels with an $SU(2)_L$ doublet and a singlet in each level, i.e., three zero modes plus $3 \times 3 \times 2$ KK modes). Note that we are interested in the couplings of the 750 GeV CP-odd scalar $\Pi_0(x)$ to SM fermions (top quarks primarily), but we also need its couplings with the rest of KK fermions, since these interactions will be crucial to generate large enough radiative couplings to photons and gluons.

We first write the effective 4D up-type quark mass matrix as

$$\begin{pmatrix} q_L^0(x) & Q_L(x) & U_L(x) \end{pmatrix} \quad M_u \begin{pmatrix} u_R^0(x) \\ Q_R(x) \\ U_R(x) \end{pmatrix} \tag{21}$$

in a basis where $q_L^0(x)$ and $u_R^0(x)$ represent three zero-mode flavors each [doublets and singlets of $SU(2)_L$], $Q_L(x)$ and $Q_R(x)$ represent three flavors and three KK levels of the vectorlike KK up-type doublets, and $U_L(x)$ and $U_R(x)$

represent three flavors and three KK levels of vectorlike KK up-type singlets. The mass matrix is thus

$$M_u = \begin{pmatrix} (y_u^0)_{3\times3} & (0)_{3\times9} & (Y^{qU})_{3\times9} \\ (Y^{Qu})_{9\times3} & (M_Q)_{9\times9} & (Y_1)_{9\times9} \\ (0)_{9\times3} & (Y_2)_{9\times9} & (M_U)_{9\times9} \end{pmatrix} \quad (22)$$

with the down sector mass matrix M_d computed in the same way.

The submatrices are obtained by evaluating the overlap integrals

$$y_u^0 = \frac{(Y_u^{5D})_{ij}}{\sqrt{k}} \int_0^{y_1} dy e^{-4\sigma(y)} \frac{v(y)}{\sqrt{2}} q_L^{0,i}(y) u_R^{0,j}(y) \quad (23)$$

$$Y^{qU} = \frac{(Y_u^{5D})_{ij}}{\sqrt{k}} \int_0^{y_1} dy e^{-4\sigma(y)} \frac{v(y)}{\sqrt{2}} q_L^{0,i}(y) U_R^{n,j}(y) \quad (24)$$

$$Y^{Qu} = \frac{(Y_u^{5D})_{ij}}{\sqrt{k}} \int_0^{y_1} dy e^{-4\sigma(y)} \frac{v(y)}{\sqrt{2}} Q_L^{m,i}(y) u_R^{0,j}(y) \quad (25)$$

$$Y_1 = \frac{(Y_u^{5D})_{ij}}{\sqrt{k}} \int_0^{y_1} dy e^{-4\sigma(y)} \frac{v(y)}{\sqrt{2}} Q_L^{m,i}(y) U_R^{n,j}(y) \quad (26)$$

$$Y_2 = \frac{(Y_u^{5D^*})_{ij}}{\sqrt{k}} \int_0^{y_1} dy e^{-4\sigma(y)} \frac{v(y)}{\sqrt{2}} Q_R^{m,i}(y) U_L^{n,j}(y), \quad (27)$$

where the indices m and n track the KK level and i, $j = 1$, 2, 3 are 5D flavor indices. The diagonal matrices $(M_Q)_{9\times9}$ and $(M_U)_{9\times9}$ are constructed with the masses of all the KK quarks involved. The masses and the profiles of the KK fermions appearing in these overlap integrals [$Q_L(y)$, $Q_R(y)$, $U_L(y)$, and $U_R(y)$] are obtained by solving differential equations for the fermion profiles

$$\partial_y(e^{(2c-1)\sigma(y)}\partial_y(e^{-(c+2)\sigma(y)}))f(y) + e^{(c-1)\sigma(y)} m_n^2 f(y) = 0, \quad (28)$$

where $f(y)$ is the KK profile. The mass eigenvalues m_n are found by imposing Dirichlet boundary conditions on the wrong chirality modes.

As mentioned before, we have included three full KK levels so that the mass matrices in the gauge basis are 21×21 dimensional matrices, which are not diagonal. One needs to diagonalize them, and by doing so, to move to the quark physical basis where all the fermion couplings can then be extracted.

In the CP-odd scalar sector, we can write the effective 4D Yukawa-type couplings to fermions in the same gauge basis as before,

$$(q_L^0(x) \quad Q_L(x) \quad U_L(x)) \mathbf{Y}_\pi \begin{pmatrix} u_R^0(x) \\ Q_R(x) \\ U_R(x) \end{pmatrix} \Pi_0(x), \quad (29)$$

where now the 21×21 coupling matrix $\mathbf{Y}_\pi$ is given by

$$\mathbf{Y}_\pi = \begin{pmatrix} (y_{\pi qu}^0)_{3\times3} & (a_{\pi qQ})_{3\times9} & (Y_{\pi qU})_{3\times9} \\ (Y_{\pi Qu})_{9\times3} & (a_{\pi QQ})_{9\times9} & (Y_1^\pi)_{9\times9} \\ (a_{\pi uU})_{9\times3} & (Y_2^\pi)_{9\times9} & (a_{\pi UU})_{9\times9} \end{pmatrix}. \quad (30)$$

The submatrices are obtained by the overlap integrals

$$y_{\pi qu}^0 = i\frac{(Y_u^{5D})_{ij}}{\sqrt{2k}} \int_0^{y_1} dy e^{-3\sigma(y)} q_L^{0,i}(y) u_R^{0,j}(y) \frac{X_\pi'(y)}{m_{\pi_0}^2 M_z(y)} \quad (31)$$

$$Y_{\pi qU} = i\frac{(Y_u^{5D})_{ij}}{\sqrt{2k}} \int_0^{y_1} dy e^{-3\sigma(y)} q_L^{0,i}(y) U_R^{n,j}(y) \frac{X_\pi'(y)}{m_{\pi_0}^2 M_z(y)} \quad (32)$$

$$Y_{\pi Qu} = i\frac{(Y_u^{5D})_{ij}}{\sqrt{2k}} \int_0^{y_1} dy e^{-3\sigma(y)} Q_L^{m,i}(y) u_R^{0,j}(y) \frac{X_\pi'(y)}{m_{\pi_0}^2 M_z(y)} \quad (33)$$

$$Y_1^\pi = i\frac{(Y_u^{5D})_{ij}}{\sqrt{2k}} \int_0^{y_1} dy e^{-3\sigma(y)} Q_L^{m,i}(y) U_R^{n,j}(y) \frac{X_\pi'(y)}{m_{\pi_0}^2 M_z(y)} \quad (34)$$

$$Y_2^\pi = i\frac{(Y_u^{5D^*})_{ij}}{\sqrt{2k}} \int_0^{y_1} dy e^{-3\sigma(y)} Q_R^{m,i}(y) U_L^{n,j}(y) \frac{X_\pi'(y)}{m_{\pi_0}^2 M_z(y)}, \quad (35)$$

and

$$a_{\pi qQ} = \frac{g_L^{5D}}{\sqrt{k}} \int_0^{y_1} dy e^{-2\sigma(y)} q_L^{0,i}(y) Q_R^{n,j}(y) \frac{X_\pi(y)}{m_{\pi_0}^2} \quad (36)$$

$$a_{\pi uU} = \frac{g_R^{5D}}{\sqrt{k}} \int_0^{y_1} dy e^{-2\sigma(y)} u_R^{0,i}(y) U_L^{n,j}(y) \frac{X_\pi(y)}{m_{\pi_0}^2} \quad (37)$$

$$a_{\pi QQ} = \frac{g_L^{5D}}{\sqrt{k}} \int_0^{y_1} dy e^{-2\sigma(y)} Q_L^{m,i}(y) Q_R^{n,j}(y) \frac{X_\pi(y)}{m_{\pi_0}^2} \quad (38)$$

$$a_{\pi UU} = \frac{g_R^{5D}}{\sqrt{k}} \int_0^{y_1} dy e^{-2\sigma(y)} U_R^{m,i}(y) U_L^{n,j}(y) \frac{X_\pi(y)}{m_{\pi_0}^2}, \quad (39)$$

where the $g_{L,R}^{5D}$ couplings are given by

$$g_L^{5D} = \frac{g^{5D}}{\cos\theta_W}(T_3 - Q_q \sin^2\theta_W) \tag{40}$$

$$g_R^{5D} = \frac{g^{5D}}{\cos\theta_W} Q_q \sin^2\theta_W, \tag{41}$$

with Q_q the charge of the quark (here $\frac{2}{3}$), θ_W the weak angle, and $T_3 = \frac{1}{2}$. Note that when the interaction originates in the 5D Yukawa couplings the profile to use is the one coming from the CP-odd Higgs component, i.e., proportional to $X'_\pi(y)$. When the interaction originates in the gauge-fermion coupling and thus comes from the Z_5 component, the profile to use is proportional to $X_\pi(y)$, with X_π being the solution of Eq. (17), using the decompositions of Eqs. (13) and (14).

When the fermion matrix in (22) is diagonalized, the coupling matrix of fermions with the CP-odd field $\Pi_0(x)$ in (30) is rotated, and we can then extract all the physical Yukawa couplings. All these couplings are needed later in order to compute the radiative couplings of $\Pi_0(x)$ with gluons and photons.

Let us first analyze the very important Yukawa coupling between $\Pi_0(x)$ and top quarks, as this coupling should dominate the decays of the CP-odd scalar. The coupling comes essentially from the entry $(y^0_{\pi qu})_{33}$ (before rotation to the physical basis) although it receives small corrections after going to the physical basis. We focus on $(y^0_{\pi qu})_{33}$ which comes from the overlap integral

$$(y^0_{\pi qu})_{33} = i\frac{(Y_u^{5D})_{33}}{\sqrt{2k}} \int_0^{y_1} dy e^{-3\sigma(y)} q_t^0(y) u_t^0(y) \frac{X'_\pi(y)}{m^2_{\pi_0} M_z(y)}. \tag{42}$$

In the RS limit, the warp factor is $\sigma(y) = ky$, and the top profiles are $q_t^0(y) = f(c_q)e^{(2-c_q)ky}$ and $u_t^0(y) = f(-c_u)e^{(2+c_u)ky}$, where $f(x)$ is a normalization factor. We also have $M_z(y) = \frac{g_5}{2c_W} v_0 e^{(a-1)ky}$, with v_0 a constant factor, so that the previous overlap integral in this limit reads

$$(y^0_{\pi qu})_{33} = i\frac{(Y_u^{5D})_{33}}{\sqrt{2k}} \frac{2c_W}{g_5} \frac{f(c_q)f(-c_u)}{v_0 m^2_{\pi_0}} \times \int_0^{y_1} dy e^{(2-a-c_q+c_u)ky} X'_\pi(y). \tag{43}$$

We integrate this by parts to find

$$(y^0_{\pi qu})_{33} = -i\frac{(Y_u^{5D})_{33}}{\sqrt{2k}} \frac{2c_W}{g_5} \frac{f(c_q)f(-c_u)}{v_0 m^2_{\pi_0}} \times \int_0^{y_1} dy (2-a-c_q+c_u) e^{(2-a-c_q+c_u)ky} X_\pi(y) + BT, \tag{44}$$

where $BT = i\frac{(Y_u^{5D})_{33}}{\sqrt{2k}} \frac{2c_W}{g_5} \frac{f(c_q)f(-c_u)}{v_0 m^2_{\pi_0}} e^{(2-a-c_q+c_u)ky} X_\pi(y)|_0^{y_1}$ is a boundary term. Note that the profile $X_\pi(y)$ has vanishing boundary conditions in the absence of Higgs localized brane kinetic terms. In that limit, we can see that the coupling of the CP-odd scalar can actually vanish, when $(2-a-c_q+c_u) = 0$ [29]. Note also that the Higgs localizer parameter a is, in this RS limit, $a \gtrsim 2$ and the bulk parameters c_q and c_u are defined such that, for example, charm or bottom quarks are assigned values more or less $c_q \in (0.45, 0.55)$ and $c_u \in (-0.5, -0.6)$, whereas for the top, we have $c_{q_3} \sim 0.45$ and $c_{u_3} > -0.45$. This means that in the RS metric, one should expect the term $(2-a-c_q+c_u)$ to vanish, in the limit of $a \sim 2$, when $c_q - c_u \sim 0$, so that the suppression in this case seems only possible for the top quark, where both c_q and c_u could be small.

Of course, when the metric background is modified away from AdS_5 (the case we consider here) and when the boundary conditions include brane kinetic terms, there will be deviations from the RS expectations outlined in the previous paragraph. Nevertheless, it is clear that the Yukawa coupling of the CP-odd scalar field to top quarks can have highly suppressed values. Another way to see this is to consider the overlap integral in Eq. (42). Because the profile $X_\pi(y)$ vanishes at the boundaries (or *almost* vanishes, for small brane kinetic terms), then its derivative $X'_\pi(y)$ will have a node in the bulk and therefore will change sign. That means that there can be some parameter choice for which it is possible for the overlap integral to vanish, since the fermion zero-mode profiles have no nodes in the bulk.

This feature is clearly seen in Fig. 2, where we plot the absolute value of the Yukawa couplings between zero-mode fermions and both the Higgs and the CP-odd scalar $\Pi_0(x)$.[6] The couplings shown are relative to the 5D bulk Higgs Yukawa coupling Y_5 and are plotted as functions of the fermion bulk mass parameter c_q and c_u (for the case where we take $c_q = -c_u$, for simplicity), for different overall KK scales. We observe that the CP-odd Yukawa couplings are fairly similar to the Higgs Yukawa couplings (i.e., exponentially sensitive to UV localization and then toplike when the zero mode is IR localized) except that there is a range of parameters where the coupling vanishes. Interestingly enough, this suppression happens for preferred values of the top-quark bulk mass parameters. This means that the existence of suppressed couplings to top quarks of

[6]We are actually plotting the values defined in Eqs. (31) and (23), i.e., the zero-mode Yukawa couplings before going to the fermion mass basis. In that basis, the couplings will inherit a small correction due to mixing with heavy KK fermions [32], so that the exact cancellation of the coupling will be replaced by a strong suppression.

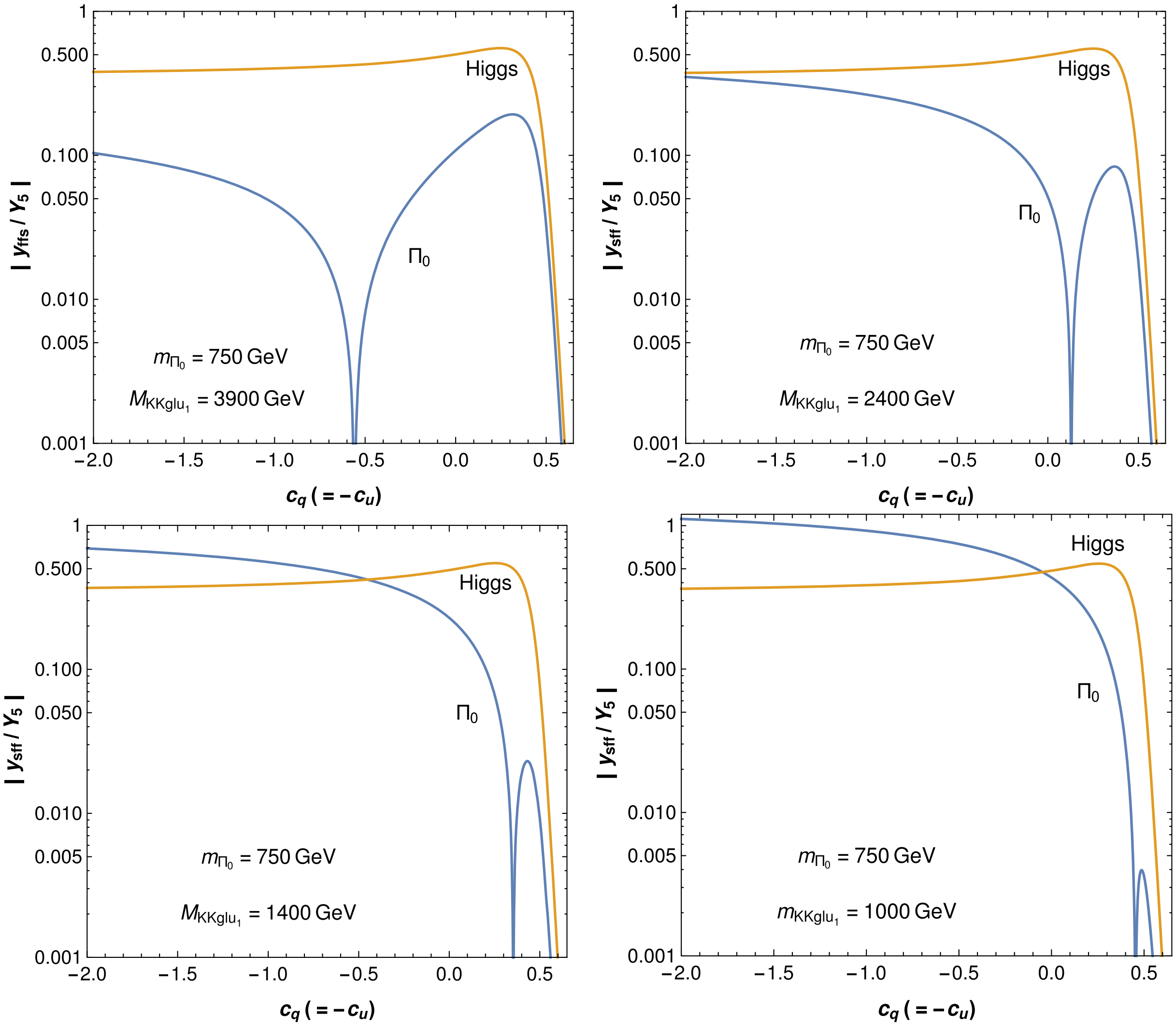

FIG. 2. Yukawa couplings between zero-mode fermions and the two lightest neutral scalars of the scenario, the 125 GeV Higgs and a 750 GeV CP-odd Π_0. The couplings are evaluated relative to the 5D bulk Higgs Yukawa coupling Y_5 and are shown as a function of the fermion bulk mass parameter c_q (in the case $c_q = -c_u$ for simplicity) for different overall KK scales, $M_{KKglu_1} = 3900$ GeV (upper left panel), $M_{KKglu_1} = 2400$ GeV (upper right panel), $M_{KKglu_1} = 1400$ GeV (lower left panel), and $M_{KKglu_1} = 1000$ GeV (lower right panel). The CP-odd scalar mass is set to 750 GeV, for illustrative purposes, and for certain values of c_q, its Yukawa coupling to top quarks can be highly suppressed for typical top-quark values of the c_i's.

the CP-odd Π_0 is a natural possibility in this scenario, thus reducing the rate of top pair production in the CP-odd decays.

B. Radiative couplings to photons and gluons

Just like in the Higgs boson case, the radiative couplings of $\Pi_0(x)$ to gluons and photons will depend on the physical Yukawa couplings y_{nn} between Π_0 and the fermions (zero modes and KK modes) running in the loop, as well as on the fermion masses m_n [the eigenvalues of the mass matrix in Eq. (22)]. The real and imaginary parts of the couplings are associated with different loop functions, $A^S_{1/2}$ and $A^P_{1/2}$, as they generate the two operators $\Pi_0 G_{\mu\nu}G^{\mu\nu}$ and $\Pi_0 G_{\mu\nu}\tilde{G}^{\mu\nu}$.[7]

The production cross section through gluon fusion is

$$\sigma_{gg\to\Pi_0} = \frac{\alpha_s^2 m_{\Pi_0}^2}{576\pi}\left[\left|\sum_{\text{quarks}} c_n^S\right|^2 + \left|\sum_{\text{quarks}} c_n^P\right|^2\right], \qquad (45)$$

[7]The Yukawa couplings of Π_0 are mostly imaginary, and thus the dominant contribution will come, as expected, from the operator $\Pi_0 G_{\mu\nu}\tilde{G}^{\mu\nu}$. Still, small real Yukawa coupling components are generated when going to the fermion mass basis, and so we keep the general formalism in our formulas.

and the decay widths to gluons and photons are

$$\Gamma_{\Pi_0\to gg} = \frac{\alpha_s^2 m_{\Pi_0}^3}{54\pi^2}\frac{1}{v^2}\left[\left|\sum_{\text{quarks}} c_n^S\right|^2 + \left|\sum_{\text{quarks}} c_n^P\right|^2\right] \quad (46)$$

$$\Gamma_{\Pi_0\to\gamma\gamma} = \frac{\alpha^2 m_{\Pi_0}^3}{192\pi^3}\frac{1}{v^2}\left[\left|\sum_{\substack{\text{quarks}\\\text{leptons}}} N_c Q_n^2 c_n^S\right|^2 + \left|\sum_{\substack{\text{quarks}\\\text{leptons}}} N_c Q_n^2 c_n^P\right|^2\right], \quad (47)$$

where α_s and α are the strong and weak coupling constants, N_c is the number of colors, and Q_n is the charge of the fermion and where

$$c_n^S = Re\left(\frac{y_{nn}}{m_n}\right)A_{1/2}^S(\tau_n) \quad \text{and} \quad c_n^P = \text{Im}\left(\frac{y_{nn}}{m_n}\right)A_{1/2}^P(\tau_n) \quad (48)$$

with $\tau_n = m_{\Pi_0}^2/4m_n^2$ and with the loop functions defined as [33]

$$A_{1/2}^S(\tau) = \frac{3}{2}[\tau + (\tau-1)f(\tau)]\tau^{-2}, \quad (49)$$

$$A_{1/2}^P(\tau) = -\frac{3}{2}f(\tau)/\tau \quad (50)$$

and with

$$f(\tau) = \begin{cases} [\arcsin\sqrt{\tau}]^2 & (\tau \le 1) \\ -\frac{1}{4}\left[\ln\left(\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}}\right) - i\pi\right]^2 & (\tau > 1). \end{cases} \quad (51)$$

For heavy KK quarks with masses m_n much greater than the CP-odd mass m_{Π_0} (i.e., when τ is very small), the loop functions are essentially constant, as they behave asymptotically as $\lim_{\tau\to 0} A_{1/2}^S = 1$ and $\lim_{\tau\to 0} A_{1/2}^P = 3/2$. On the other hand, for light quarks (all the SM quarks except the top and bottom), the loop functions essentially vanish asymptotically as $\lim_{\tau\to\infty} A_{1/2}^S = \lim_{\tau\to\infty} A_{1/2}^P = 0$.

Moreover, we investigate a parameter region where the couplings of Π_0 to top quarks are highly suppressed. This means that the production mechanism must rely exclusively on the heavy KK fermions running in the loop, and as we have seen, this coupling depends on the ratio $\frac{y_{nn}}{m_n}$ between the physical Yukawa coupling and the mass of the fermion running in the loop. To have an idea of the relative contribution of each of these KK fermions in the loop, in Fig. 3, we plot the mass normalized Yukawa couplings of Standard Model Higgs with top quarks, of Higgs with the first KK fermion, and of Π_0 to the first KK fermion, for different values of the KK scale. As expected, we see that the c_q dependence is mild (i.e., all KK fermions of any flavor will couple with similar strength), and also, as expected, we observe that the mass normalized couplings are quite suppressed with respect to the SM top-quark case. Still the multiplicity of KK fermions is high, since there are six families of quarks and three families of charged leptons (the latter run in the diphoton loop), and for each family, there are a few KK levels that give important contributions to the rate.

A numerical scan of the couplings, including all families and three full KK levels is computationally too intensive, so in order to produce the couplings plotted in Fig. 3, we performed an approximation, sufficient for the purposes of the graph.

The KK fermion Yukawa couplings plotted neglect mixings between different KK levels and different quark flavours, and with the zero-mode fermions. They are obtained as follows. Consider the 2×2 KK mass matrix

$$(\,Q_L(x) \quad U_L(x)\,)M_u\begin{pmatrix} Q_R(x) \\ U_R(x)\end{pmatrix}. \quad (52)$$

Here $Q_L(x)$ and $Q_R(x)$ represent a single flavour and a single KK level of the up-type doublets, and $U_L(x)$ and $U_R(x)$, represent a single flavour and a single KK level of the up-type singlets. The mass matrix is thus

$$M_u = \begin{pmatrix} m_Q & Y_1 \\ Y_2 & m_U \end{pmatrix}, \quad (53)$$

where the diagonal entries are the KK masses (large), whereas the off-diagonal entries are coming from Yukawa couplings and are therefore smaller. In order to give a simple estimate, we take for simplicity the fermion bulk mass parameters as $c_q = -c_u$ and the bulk Higgs Yukawa Y^{5D} to be real, which leads to $Y_1 = Y_2$ and $m_Q = m_U = m_{KK}$, with the masses and profiles obtained by solving Eq. (28). With the KK fermion profiles, one obtains the off-diagonal entries

$$Y_1 = \frac{(Y_u^{5D})}{\sqrt{k}}\int_0^{y_1} dy e^{-4\sigma(y)}\frac{v(y)}{\sqrt{2}}Q_L(y)U_R(y). \quad (54)$$

The matrix that diagonalizes (53) in this simple limit ($c_q = -c_u$) is $\begin{pmatrix} -\frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} \\ \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}}\end{pmatrix}$ and the eigenvalues are $m_1 = m_{KK} - Y_1$ and $m_2 = m_{KK} + Y_1$.

Now, we apply this rotation to the CP-odd Yukawa coupling matrix

$$Y_\Pi = \begin{pmatrix} \mathcal{O}(g) & Y_\pi \\ Y_\pi & \mathcal{O}(g)\end{pmatrix}, \quad (55)$$

where

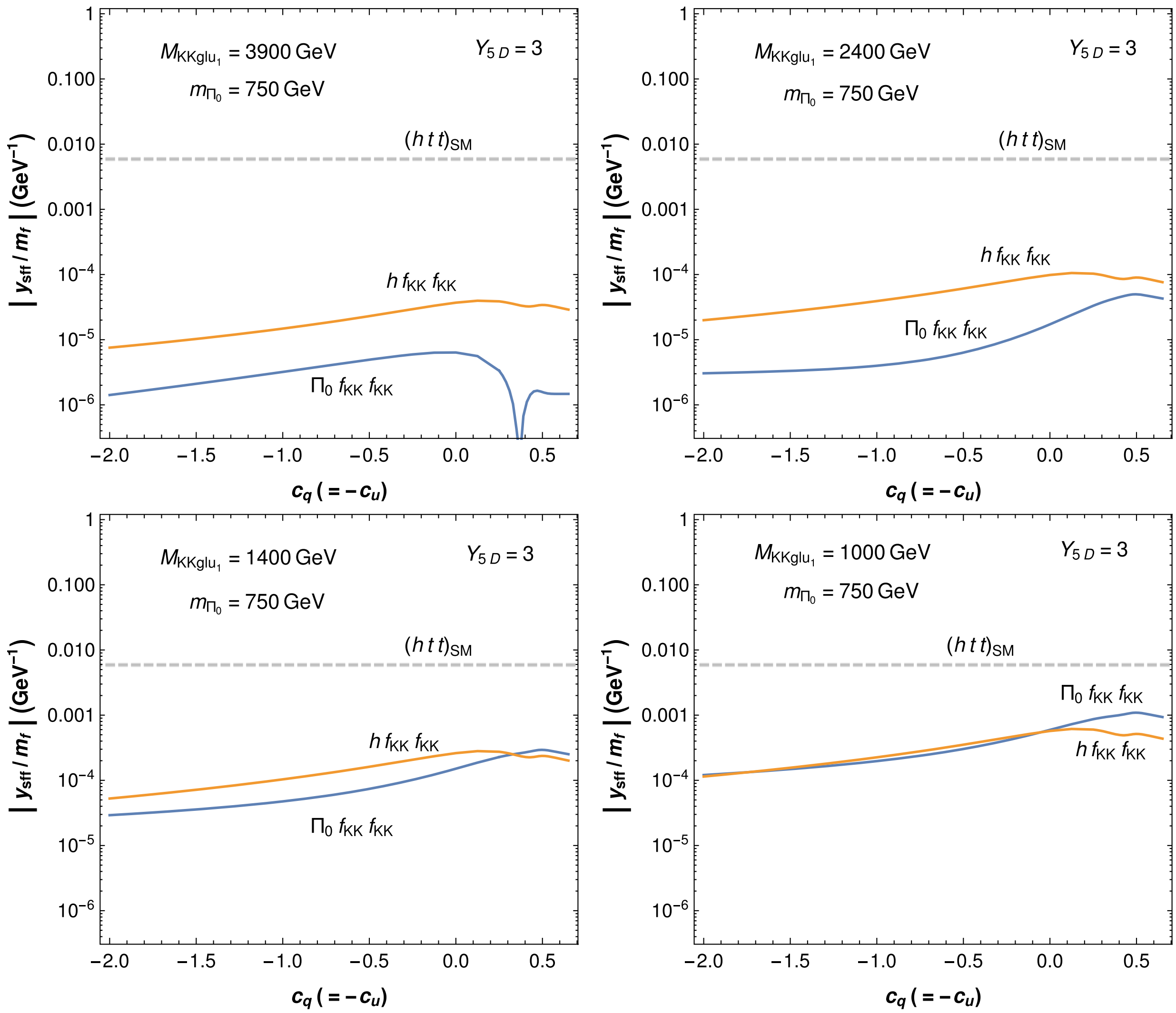

FIG. 3. Yukawa couplings between the lightest KK fermion and the 125 GeV Higgs (middle curves) and of the lightest KK fermions and the *CP*-odd Π_0 (chosen to have a mass of 750 GeV) (lower curves), divided by the KK fermion mass, for different values of the lightest KK gluon mass M_{KKglu_1}, as indicated on the panels. This mass normalized Yukawa coupling gives an estimate of the relative contribution of the respective KK fermion to the radiative coupling of the scalar to gluons and photons, to be compared with the mass normalized SM coupling of Higgs to top quarks (shown as a dashed line).

$$Y_\pi = i\frac{(Y_u^{5D})_{ij}}{\sqrt{2k}}\int_0^{y_1} dy e^{-3\sigma(y)} Q_L(y) U_R(y) \frac{X'_\pi(y)}{m_{\pi_0}^2 M_z(y)} \quad (56)$$

and where for simplicity we have neglected gauge couplings terms compared to IR Yukawa terms (a safe assumption when Y^{5D} is large).

After diagonalization, we obtain the two physical couplings between Π_0 and the KK fermions. When we normalize the couplings by the two eigenmasses and add the two contributions,[8] we obtain

$$\sum_{i=1}^{2}\frac{y_i}{m_i} = \frac{y_1^\pi}{m_1} + \frac{y_2^\pi}{m_2} = -2\frac{Y_1 Y_\pi}{m_{KK}^2 - Y_1^2} \simeq -2\frac{Y_1 Y_\pi}{m_{KK}^2}. \quad (57)$$

The last expression corresponds to the *mass normalized* Yukawa couplings of Π_0 plotted in Fig. 3, and this describes very closely the behavior of the couplings obtained in the full flavor calculation. The parametric dependence of these couplings is $Y_{5D}^2 v/m_{KK}^2$, so that if $Y^{5D} \sim 3$ we expect mass normalized couplings of order $(10^{-3}$–$10^{-4})$ GeV^{-1}, if the overlap integral is of $\mathcal{O}(1)$. Since all the profiles of the integral are IR localized, one expects that integral to be $\mathcal{O}(1)$, although the precise numerical result varies between 0.5 and 0.05, depending on the values of the c_q parameter, as shown in the plots.

[8]One needs to add the two contributions since there is a cancellation happening level by level.

All in all, it seems likely that, after taking into consideration all the fermion flavors, and for a KK scale of order 1–2 TeV, the overall KK fermion contribution to the radiative couplings of $\Pi_0(x)$ to photons and gluons can be close to the top-quark contribution to the gluon and photon couplings of the Higgs in the SM model.

C. $\Pi_0 Zh$ coupling

The coupling between the CP-odd scalar, the Z boson, and the Higgs will be extracted from the kinetic operator of the 5D Higgs,

$$\int d^4x dy e^{-2\sigma} D_\mu H^\dagger D^\mu H \left(1 + \delta(y - y_i)\frac{d_i}{k}\right). \quad (58)$$

Expanding the SM-like Higgs mode using Eq. (4) as well as the SM-like Z_μ and the 750 GeV Π_0 using Eqs. (5) and (6), we can obtain the coefficient $g_{\Pi hZ}$ of the operator $Z^\mu(x)(h(x)\partial_\mu \Pi_0(x) + \Pi_0(x)\partial_\mu h(x))$,

$$g_{\Pi hZ} = \frac{g_5^2}{4c_W^2}\int dy e^{-2\sigma} v(y) h(y) f_z(y) \times \frac{X'(y)}{M_z^2(y) m_\pi^2}\left(1 + \delta(y - y_i)\frac{d_i}{k}\right). \quad (59)$$

Now, since $M_z(y) = \frac{g_5}{2c_W} e^{-\sigma} v(y)$, $h(y) \sim v(y)/v_4$, and $f_z \simeq 1/\sqrt{y_1}$, we can write

$$g_{\Pi hZ} \simeq \frac{1}{\sqrt{y_1} v_4 m_\pi^2}\left(X(y_1) + X'(y_1)\frac{d_1}{k}\right) = 0, \quad (60)$$

where we have used the boundary conditions for the profile $X(y)$ [see Eq. (18)] and assumed no UV brane kinetic term ($d_0 = 0$).

The coupling should thus vanish in the limit of the flat Z boson profile $f_z(y)$ and when the nontrivial Higgs VEV $v(y)$ is proportional to the Higgs scalar profile $h(y)$. Corrections to these limits scale as v_4^2/m_{KK}^2 and m_h^2/m_{KK}^2 in the RS case, and so we expect the overall coupling to be highly suppressed.

The partial width for the decay $\Pi_0 \to hZ$ is [34]

$$\Gamma(\Pi_0 \to hZ) = \frac{g_{\Pi hZ}^2}{16\pi}\frac{m_Z^2}{m_\pi}\sqrt{\lambda(m_h^2, m_Z^2; m_\pi^2)}\lambda(m_h^2, m_\pi^2; m_Z^2), \quad (61)$$

where m_Z, m_π, and m_h are the masses of the particles involved and where $\lambda(x, y; z) = (1 - x/z - y/z)^2 - 4xy/z^2$. With the masses $m_Z = 91$ GeV, $m_\pi = 750$ GeV, and $m_h = 125$ GeV, the width becomes $\Gamma(\Pi_0 \to hZ) \sim (900 g_{\Pi hZ}^2)$ GeV.

For example, choosing $m_{\Pi_0} = 750$ GeV, we compute numerically $g_{\Pi hZ}$ for three different values of M_{KKglu_1} and find

M_{KKglu_1}	1000 GeV	1400 GeV	2400 GeV
$g_{\Pi hZ}$	3.8×10^{-4}	3.1×10^{-3}	1.1×10^{-2}
$\Gamma(\Pi_0 \to hZ)$	1.3×10^{-4} GeV	8.8×10^{-3} GeV	0.11 GeV

Note that the couplings and widths are small, but we observe that the partial width becomes larger as the KK mass scale is increased.

D. Estimates and numerical results

With all the previous ingredients, one can estimate the viability of this scenario in terms of the possible diphoton excess. When the bulk mass parameters of the top quark are around $|c_{u_3}| \sim 0.35$, we know that the top quark will have highly suppressed couplings to Π_0, as shown in the third panel of Fig. 2. At the same time, the couplings of the KK tops and all other KK quarks will have relatively strong Yukawa couplings to Π_0 (third panel of Fig. 3), so that the contribution of each of them to the radiative coupling of Π_0 to gluons is about an order of magnitude smaller than the top contribution to the $h - glu - glu$ coupling of the SM. Thus, it is possible that the overall contribution of all flavors and KK excitations can make up for the suppressed top couplings, so that the production cross section of Π_0 is similar to that of a heavy SM-like Higgs.

For a 750 GeV Higgs, the production cross section through gluon fusion, at the LHC running at 13 TeV is $497 fb$ [35], so, roughly, here, let us assume this to be the production cross section for the Π_0 of the same mass.

TABLE I. Boson masses obtained for $M_{KK} \simeq 1300$ GeV and $M_{KK} \simeq 1000$ GeV, with the Higgs VEV profile as delocalized as possible and where the Higgs brane kinetic terms are $d_1 \simeq 0.32$ and $d_1 \simeq 0.15$, respectively. We list, in order, the masses of the $n = 1$ gluon, the SM Higgs boson (which is the $n = 0$ CP-even Higgs), the $n = 0$ CP-odd Higgs boson, the $n = 1$ CP-even Higgs boson, the $n = 0$ charged Higgs boson, the $n = 1$ Z boson, and the $n = 1$ W boson.

M_{g_1}	M_h	M_{Π_0}	M_{H_0}	$M_{\Pi_0^\pm}$	M_{Z_1}	M_{W_1}
1303 GeV	125 GeV	750 GeV	1065 GeV	665 GeV	1285 GeV	1289 GeV
1005 GeV	125 GeV	750 GeV	1115 GeV	661 GeV	1003 GeV	1004 GeV

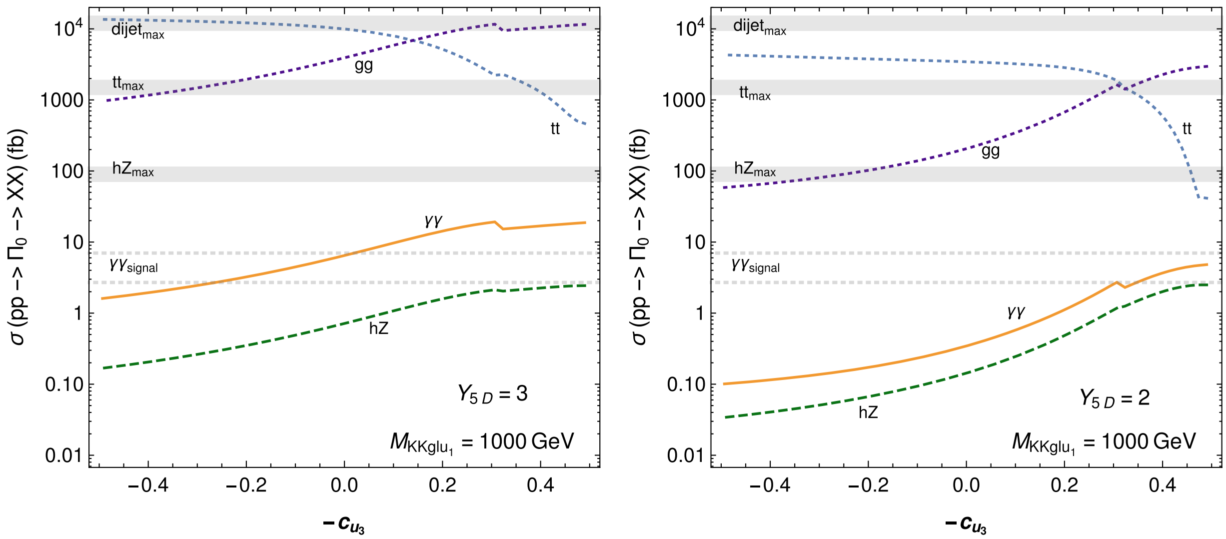

FIG. 4. Production cross sections at the LHC-13 TeV in the channels $\gamma\gamma$, $t\bar{t}$, and hZ for a 750 GeV CP-odd scalar. The KK scale is $M_{KKglu_1} = 1000$ GeV, and the 5D Yukawa couplings are $Y_{5D} \simeq 3$ (left panel) and $Y_{5D} \simeq 2$ (right panel). The shaded bands are LHC bounds as of winter of 2015, and the dotted horizontal band represents the excess reported in December 2015 (now defunct). We see that a visible signal can always be produced at this KK scale (closely below the $\gamma\gamma$ dotted band), but for larger Yukawa couplings, all rates seem too large, and thus lower Yukawa couplings seem to be preferred. Top pairs should be observed roughly at the same time as diphotons.

Since the decays of Π_0 into top quarks and hZ are suppressed in this parameter space point, and its decays to WW and ZZ can only be radiative via the CP-odd gauge boson kinetic operator, the main decay channel is into gluons, so that the branching of the diphoton channel should be very roughly

$$\mathrm{Br}(\Pi_0 \to \gamma\gamma) \sim \frac{\alpha_{em}^2}{8\alpha_s^2}\frac{N_\gamma}{N_{glu}}, \tag{62}$$

where N_γ and N_{glu} are the multiplicities of states running in the $(\Pi_0\gamma\gamma)$ loop and in the $(\Pi_0 gluglu)$ loop, respectively. In the diphoton loop, there are three extra families of charged lepton KK excitations making the multiplicity of states greater. If their multiplicity and their Yukawa couplings can partially make up for the color factor of 8, then the diphoton cross section might become of $\mathcal{O}(fb)$, and thus be easily observed.

To complete the analysis, we perform a full numerical computation of production and branching ratios in a setup where we consider an effective 4D scenario including three full KK levels for all fields; i.e., we consider 21×21 fermion mass matrices, which we diagonalize in order to obtain the physical Yukawa couplings. We choose a set of c-parameters and 5D Yukawa entries such that the SM masses and mixings are reproduced; the specific flavor choice for these parameters should not affect much the overall results since these depend on overlap integrals between IR localized fields, with very loose c-dependence. We choose the background metric parameters so that precision electroweak bounds are kept at bay, i.e., $\nu = 0.5$ and $y_s = 1.04 y_1$. Two average 5D Yukawa scales are considered, $Y_{5D} \simeq 3$ and $Y_{5D} \simeq 2$, to show the dependence on this parameter, and we also consider two different KK mass scales, $M_{KKglu_1} = 1000$ GeV and $M_{KKglu_1} = 1300$ GeV, which turn out to lead to successful signal generation. For completeness, in Table I, we give the spectrum of the lightest massive bosons corresponding to these two KK mass scales.[9]

In order to see how tuned the choice of the top c-parameter is, we plot the production cross section of the CP-odd resonance, followed by decays into $\gamma\gamma$, $t\bar{t}$, and Zh, as functions of c_{u_3} (the bulk mass parameter of the 5D singlet top quark), with the doublet bulk mass parameter fixed at $c_{q_3} = 0.4$. (This value ensures typically suppressed bounds from $Zb_L\bar{b}_L$ bounds [9].) The results are shown in Figs. 4 and 5, and in both cases, we show results for $Y_{5D} \simeq 3$ and $Y_{5D} \simeq 2$ to illustrate the sensitivity on this bulk parameter, crucial for enhancing the radiative couplings of Π_0. When the KK scale is smaller and 5D Yukawa couplings are larger, the production of top pairs and gluon pairs can be quite large. By reducing the 5D Yukawa couplings, enough visible diphoton signals can be generated with dijets and top pairs under control, as well as the Zh decay. For slightly smaller KK scales, one expects a similar behavior, but such that 5D Yukawa couplings

[9] Note that, as mentioned in the Introduction, TeV scale KK gauge bosons are difficult to probe at the LHC in these scenarios because their couplings to light fermions are very suppressed. For larger integrated luminosities of $\sim 100\ fb^{-1}$, they may yet be accessible [13].

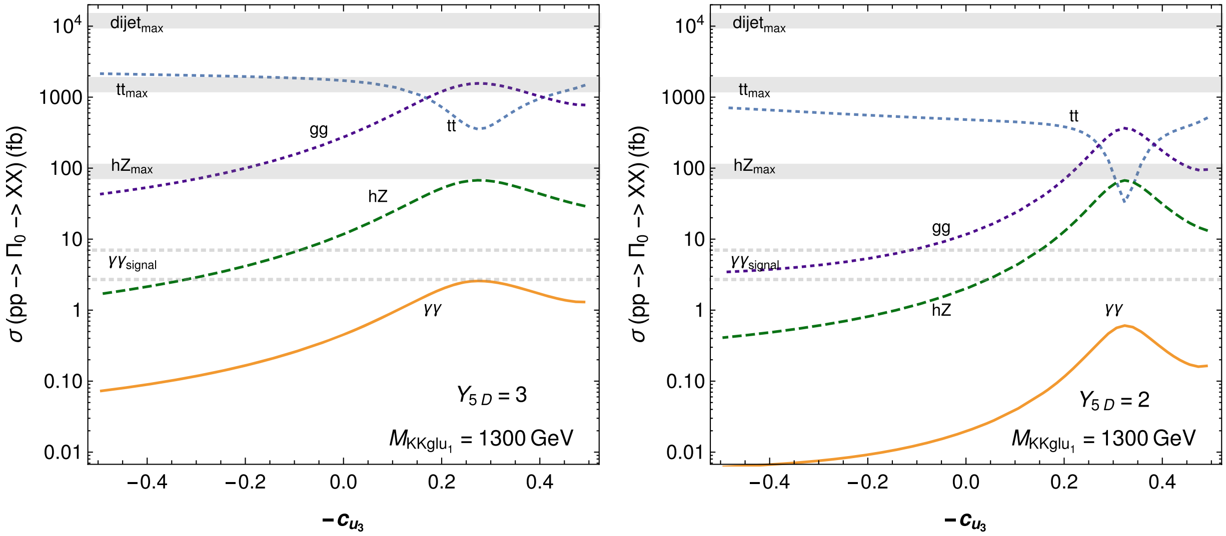

FIG. 5. Same as the previous figure but for a KK scale of 1300 GeV. In this case, a signal can still be achieved, but within a more localized region near $c_{u_3} \sim 0.25$. The bound from hZ seems to become the main constraint but also becomes a complementary signal for such a state.

should be even smaller in order to suppress overproduction of Π_0 particles.

This leads to the question of how large the KK scale can be and still manage to produce a signal close to the defunct excess reported in December 2015. We observe that at $M_{KKglu_1} = 1300$ GeV, with larger 5D Yukawa couplings ($Y_{5d} \geq 3$), one can get close enough to that signal. The branching fraction into hZ increases, threatening competition with the $\gamma\gamma$ signal. But more interestingly, now the signal production requires the value of $|c_{u_3}|$ to be located around the point where the top Yukawa couplings of Π_0 are suppressed, in this case around $|c_{u_3}| \sim 0.3$. It is interesting that for that value of c_{u3} the top Yukawa coupling Y_{33} is *not* required to be much larger than the rest of 5D Yukawa couplings in order to reproduce the top-quark mass.[10]

We conclude therefore that in order to obtain a visible diphoton excess at a low *intermediate* mass scale (500–800 GeV) the overall KK scale should be around $1000 \text{ GeV} \leq M_{KKglu_1} \leq 1300$ GeV, with 5D Yukawa couplings $Y_{5D} \lesssim 2$–3 (i.e., pretty constrained). For those values, the signal seems quite generic (i.e., small, but typical, c_{u_3} values are required, but these happen to be the values required to reproduce a heavy enough top-quark mass).

V. DISCUSSION

We performed an analysis of the scalar sector of warped space models to investigate whether the *minimal* model can accommodate a *moderately light* resonance in the diphoton channel at CMS and ATLAS. We showed that in the simplest extra-dimensional extension of the SM, that is with a 5D Higgs doublet living in the bulk, the lowest pseudoscalar KK excitation can be responsible for such a signal. We emphasize that, unlike other explanations relying on scalar fields in warped models, ours does not introduce any new fields or representations but relies exclusively on Higgs brane kinetic terms to lower the KK mass of the lightest CP-odd Higgs resonance. This makes the model extremely constrained, with the only new parameter being the IR brane kinetic coefficient d_1. The lightest CP-odd excitation, a mixture of the 5D Higgs field and Z_5, does not decay at tree level into WW or ZZ and, over a range of the parameter space, can have suppressed couplings to the top quark, and thus a small decay width into $t\bar{t}$. The production through gluon fusion can be loop enhanced through the effects of the usual KK fermion modes, and so can the diphoton decay. The coupling to Zh is also suppressed, although it starts increasing dangerously for KK masses above 1500 GeV.

We also showed that in AdS_5 spaces (RS-type models) (with a fixed KK scale of $M_{KKglu_1} = 1500$ GeV) the presence of Higgs brane kinetic terms can lower the mass of the lightest CP-odd scalar from 2500 to about 1750 GeV. On the other hand, when the metric is modified slightly away from AdS_5, the Higgs brane kinetic terms can produce CP-odd scalars as light as 500 GeV (with the same fixed KK scale of $M_{KKglu_1} = 1500$ GeV).

Within these modified metric scenarios, and for KK mass scales at around 1 TeV (consistent with precision electroweak bounds), this CP-odd resonance obeys the (current) experimental constraints. We analyzed its production and decay for several values of the lowest KK gluon mass (M_{KKglu_1}).

[10]In scenarios with a modified AdS_5 metric, generic values of $Y_{33} \sim 3 - 4$ are required in order to reproduce the top-quark mass. When c_{u3} and c_{q3} are both around $\sim 0.2 - 0.4$, the value of Y_{33} can be at the same level as all other 5D Yukawa couplings.

Our analysis is quite general, even though we show an analysis here for a CP scalar mass of 750 GeV. The general conclusion to be taken from our analysis here is that warped space models, without any new particles, can explain a (relatively) light diphoton resonance at the LHC. Should a diphoton excess be found at higher mass values, even the RS model might accommodate such a state without the need modify the metric.

Among the general features of the light CP-odd scalar resonance, resulting from the 5D Higgs doublet is the fact that its coupling to top pairs can be suppressed for appropriate top bulk mass parameters. Also, its coupling to Zh is generically suppressed due to the boundary conditions of the CP-odd state. In addition, the model predicts that the spectrum for the CP-odd and the charged scalars is essentially the same since their differential equations and boundary conditions are almost identical. This means that the lightest charged Higgs boson is expected to have a mass very close to the pseudoscalar mass, so about 750 GeV, in the scenario in which the latter is the diphoton resonance. If the charged Higgs happens to be connected to the light pseudoscalar through the mechanism envisioned here, its production through the Yukawa coupling to $t\bar{b}$ would also be suppressed, in the same manner in which the pseudoscalar couplings to $t\bar{t}$ are suppressed. Thus, another prediction is that the charged Higgs would be difficult to observe *if* the CP-odd Higgs leads to many diphotons.

Decays into ZZ and WW could be seen later, since their couplings to the light pseudoscalar are loop induced and thus one expects them to be similar to the photon couplings. However, the massive gauge bosons must decay further into leptons, suppressing the strength of the signal with respect to $\gamma\gamma$. Top pair production, dijet production, and the Zh signal should be around the corner, with rates similar in size to the diphoton channel rates.

Finally, this is a warped space scenario allowing for light KK partners in general, making it quite appealing and distinguishable. The whole scalar sector in particular might also be quite light. A study of the general features of the general scalar sector in these scenarios, without explicit focus on an exotic diphoton signal, is currently underway.

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APPENDIX: THE EQUATIONS OF MOTION WITH BRANE-LOCALIZED KINETIC TERMS

In this section, we consider the effect of brane localized kinetic terms associated with the 5D Higgs doublet and also with the gauge bosons. For simplicity, let us consider a 5D toy model with a Higgs scalar $H(x,y)$ charged under a local $U(1)$, defined by the following action,

$$S = \int d^4x dy \sqrt{g}\left(-\frac{1}{4}F_{MN}^2 + |D^M H|^2 - V(H)\right) \tag{A1}$$

$$+ \sum_i \int d^4x dy \sqrt{g}\delta(y - y_i) \times \left(\frac{1}{4}r_i F_{MN}^2 + d_i|D^M H|^2 - \lambda_i(H)\right), \tag{A2}$$

where $F_{MN} = \partial_M A_N - \partial_N A_M$ and for simplicity we set the gauge coupling constant to unity in the appropriate mass dimensions. The background spacetime metric is assumed to take the form

$$ds^2 = e^{-2\sigma(y)}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2, \tag{A3}$$

where $\sigma(y)$ is the warp factor.

We are interested in studying the effective 4D perturbative spectrum of the 5D Higgs field and the 5D gauge boson, around a nontrivial Higgs vacuum profile solution $\langle H\rangle = v(y)$,

$$H(x,y) = \frac{1}{\sqrt{2}}(v(y) + h(x,y))e^{i\pi(x,y)}. \tag{A4}$$

In particular, we are interested in the CP-odd Higgs perturbations $\pi(x,y)$, the equations of motion of which are coupled with the gauge boson perturbations. The equations read

$$\begin{aligned}&(1 + r_i\delta_i)\partial_\mu\partial^\mu A_\mu - ((1 + r_i\delta_i)e^{-2\sigma}A_\mu')' + (1 + d_i\delta_i)M_A^2 A_\mu \\ &\quad + \partial_\mu((1 + d_i\delta_i)M_A^2\pi - (1 + r_i\delta_i)\partial^\nu A_\nu \\ &\quad - ((1 + r_i\delta_i)e^{-2\sigma}A_5)') = 0\end{aligned} \tag{A5}$$

$$\begin{aligned}&(1 + r_i\delta_i)\partial_\mu\partial^\mu A_5 - (1 + r_i\delta_i)\partial^\nu A_\nu' + (1 + d_i\delta_i) \\ &\quad \times M_A^2(\pi' - A_5) = 0\end{aligned} \tag{A6}$$

$$\begin{aligned}&(1 + d_i\delta_i)\partial_\mu\partial^\mu\pi - (1 + d_i\delta_i)\partial^\nu A_\nu + M_A^{-2}((1 + d_i\delta_i) \\ &\quad \times M_A^2 e^{-2\sigma}(\pi' - A_5))' = 0,\end{aligned} \tag{A7}$$

where $M_A = v(y)e^{-\sigma}$ and where $d_i\delta_i \equiv \sum_i d_i\delta(y - y_i)$ and $r_i\delta_i \equiv \sum_i r_i\delta(y - y_i)$.

We fix partially the 5D gauge by imposing

$$(1 + d_i\delta_i)M_A^2\pi - (1 + r_i\delta_i)\partial^\nu A_\nu - ((1 + r_i\delta_i)e^{-2\sigma}A_5)' = 0. \tag{A8}$$

The previous gauge fixing equation reads in the bulk

$$M_A^2\pi - \partial^\nu A_\nu - (e^{-2\sigma}A_5)' = 0. \tag{A9}$$

Note that if we evaluate the bulk constrain Eq. (A9) at $y = y_1 - \epsilon$ (i.e., right before the IR brane) we obtain

$$\partial^\nu A_\nu|_{y_1-\epsilon} = M_A^2\pi - (e^{-2\sigma}A_5)'|_{y_1-\epsilon}. \tag{A10}$$

On the other hand, the effect of the delta functions in Eq. (A8) is to produce a discontinuity in the 5D field A_5 at the brane location as

$$d_1 M_A^2 \pi - r_1 \partial^\nu A_\nu|_{y_1-\epsilon} = [e^{-2\sigma} A_5]_{y_1-\epsilon}^{y_1}, \quad \text{(A11)}$$

and similarly for the UV brane. We can thus multiply (A10) by r_1 and use it in the previous equation and find the necessary boundary condition between π and A_5, which ensures that A_μ is completely decoupled, even on the brane. We find

$$(d_1 - r_1) M_A^2 \pi + r_1 (e^{2\sigma} A_5)'|_{y_1-\epsilon} = -e^{-2\sigma} A_5|_{y_1-\epsilon}, \quad \text{(A12)}$$

where we have taken A_5 to vanish exactly on the brane, but it jumps right before the boundary.

Inserting the gauge choice in the coupled equations of motion, one manages to decouple the gauge modes A_μ (in both the bulk and the branes) with a bulk equation

$$\partial_\mu \partial^\mu A_\mu - (e^{-2\sigma} A'_\mu)' + M_A^2 A_\mu = 0 \quad \text{(A13)}$$

and jump condition on A'_μ

$$r_1 \partial_\mu \partial^\mu A_\mu + d_1 M_A^2 A_\mu|_{y_1-\epsilon} = -e^{-2\sigma} A'_\mu|_{y_1-\epsilon}, \quad \text{(A14)}$$

where A'_μ again vanishes exactly on the brane but has a jump right before it. We separate variables

$$A_\mu(x, y) = V_\mu^{4d}(x) V_y(y) \quad \text{(A15)}$$

and find the separated equations for the gauge boson tower become

$$\partial_\mu \partial^\mu V_\mu^{4d}(x) + m_A^2 V_\mu^{4d}(x) = 0 \quad \text{(A16)}$$

$$(e^{-2\sigma} V'_y)' + (m_A^2 - M_A^2) V_y = 0 \quad \text{(A17)}$$

with jump conditions on V'_y,

$$(d_i M_A^2 - r_i m_A^2) V_y|_{y_1-\epsilon} = -e^{-2\sigma} V'_y|_{y_1-\epsilon}, \quad \text{(A18)}$$

where the 4D effective mass m_A^2 is the constant of separation of variables.

The remaining equations are, in the bulk,

$$\partial_\mu \partial^\mu A_5 + M_A^2 (\pi' - A_5) - (M_A^2 \pi)' + ((e^{-2\sigma} A_5)')' = 0 \quad \text{(A19)}$$

$$\partial_\mu \partial^\mu \pi + M_A^{-2} (M_A^2 e^{-2\sigma} (\pi' - A_5))' - M_A^2 \pi + (e^{-2\sigma} A_5)' = 0, \quad \text{(A20)}$$

and the fields must verify the boundary conditions of Eq. (A12).

We now perform a mixed separation of variables,

$$A_5(x, y) = G(x) g(y) + \pi_x(x) \eta(y) \quad \text{(A21)}$$

$$\pi(x, y) = G(x) h(y) + \pi_x(x) \xi(y), \quad \text{(A22)}$$

which is to say that both $A_5(x, y)$ and $\pi(x, y)$ each contain some Goldstone and CP-odd degrees of freedom. The profiles $g(y)$, $\eta(y)$, $h(y)$, and $\xi(y)$ quantify how much of each they contain. Of course, the functions g and h are inter-connected, and η and ξ are also inter-connected. The relationships are such that $G(x)$ and $\pi_x(x)$ decouple. With the choice

$$h(y) = \frac{K(y)}{m_G^2} \quad \text{(A23)}$$

$$g(y) = \frac{K'(y)}{m_G^2} \quad \text{(A24)}$$

$$\eta(y) = \frac{e^{2\sigma}}{m_\pi^2} X(y) \quad \text{(A25)}$$

$$\xi(y) = \frac{1}{m_\pi^2 M_A^2} X'(y) \quad \text{(A26)}$$

and using the mixed separation of variables in (A21) and (A22), the mixed equations of motion in (A20) decouple, and we obtain

$$\frac{e^{2\sigma}}{M_A^2} X(y) \partial_\mu \partial^\mu \pi_x(x) + \pi_x(x) e^{2\sigma} X(y) - \pi_x(x) [M_A^{-2} X'(y)]' = 0 \quad \text{(A27)}$$

$$K(y) \partial_\mu \partial^\mu G(x) + G(x)[(K' e^{-2\sigma})' + M_A^2 K(y)] = 0. \quad \text{(A28)}$$

Once separated, we obtain, for the CP-odd physical scalars

$$\partial_\mu \partial^\mu \pi_x(x) + m_\pi^2 \pi_x(x) = 0 \quad \text{(A29)}$$

$$(M_A^{-2} X')' + e^{2\sigma} \left(\frac{m_\pi^2}{M_A^2} - 1 \right) X = 0, \quad \text{(A30)}$$

with boundary conditions

$$d_i X' = -X, \quad \text{(A31)}$$

and for the Goldstone modes

$$\partial_\mu \partial^\mu G(x) + m_G^2 G(x) = 0 \quad \text{(A32)}$$

$$(K' e^{-2\sigma})' + (m_G^2 - M_A^2) K = 0, \quad \text{(A33)}$$

with boundary conditions

$$(d_i M_A^2 - r_i m_G^2) K = -e^{-2\sigma} K'. \quad \text{(A34)}$$

Note that both the equations and boundary conditions for the Goldstone bosons are identical to the ones for the gauge boson tower, as they should be, so that they can then be gauged away level by level with the remaining gauge fixing freedom.

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