

Top and bottom partners, Higgs boson on the brane, and the $t\bar{t}h$ signalGilles Couture,^{1,*} Mariana Frank,^{2,†} Cherif Hamzaoui,^{1,‡} and Manuel Toharia^{3,2,§}¹*GTPP, Département des Sciences de la Terre et de L'Atmosphère, Université du Québec à Montréal, Case Postale 8888, Succursale Centre-Ville, Montréal, Québec H3C 3P8, Canada*²*Department of Physics, Concordia University, 7141 Sherbrooke Street West, Montreal, Québec H4B 1R6, Canada*³*Physics Department, Dawson College, 3040 Sherbrooke Street, Westmount, Québec H3Z 1A4, Canada*
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Current LHC results indicate a possible enhancement in the production of Higgs bosons in association with top quarks ($t\bar{t}h$) over the Standard Model (SM) expectations, suggesting an increase in the top Yukawa coupling. To explain these results, we study the effect of adding to the SM a small set of vectorlike partners of the top and bottom quarks with masses of order ~ 1 TeV. We consider Yukawa coupling matrices with vanishing determinant and show that Higgs production through gluon fusion is not affected by deviations in the top quark Yukawa coupling, and in fact depends only on deviations in the bottom quark Yukawa coupling. We call this scenario the “brane Higgs limit,” as it can emerge naturally in models of warped extra dimensions with all matter fields in the bulk, except the Higgs (although it could also occur in 4D scenarios with vectorlike quarks and special flavor symmetries forcing the vanishing of the Yukawa determinants). We show that the scenario is highly predictive for all Higgs production/decay modes, making it easily falsifiable, maybe even at the LHC Run 2 with higher luminosity.

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Run 1 of the LHC culminated in the discovery of the Higgs boson at 125 GeV. After the Higgs discovery, the most important question is, naturally, where is the new physics beyond the Standard Model (SM), and how will it manifest itself? Minimal additions to the SM could be revealed as an unexpected excess in dileptons, W^+W^- , ZZ , diphotons, $t\bar{t}$ or $b\bar{b}$, indicating the presence of a new boson resonance. More involved new physics could also appear as missing energy, or any other signals indicating the presence of supersymmetry or extra dimensions or other exotic particles. But also, new physics could manifest itself indirectly. In particular it could affect the production cross section and decay widths of the Higgs boson, expected to be measured with increased precision during the current Run 2 of the LHC. There are already several promising signals in the Run 1 data indicating possible deviations from the SM expectations. In particular, both ATLAS and CMS report a possible increase in the signal strength of the $t\bar{t}h$ associated production in the LHC data. Of particular interest from Run 1 at CMS and ATLAS are the same-sign dilepton ($SS2l$) and trilepton ($3l$) signals coming from leptonic Higgs decays in the associated $t\bar{t}h$ production events. The best-fit signal strengths are found to be $\mu_{SS2l} = 5.3_{-1.8}^{+2.1}$ and $\mu_{3l} = 3.1_{-2.0}^{+2.4}$ at CMS [1], and $\mu_{SS2l} = 2.8_{-1.9}^{+2.1}$

and $\mu_{3l} = 2.8_{-1.8}^{+2.2}$ at ATLAS [2]. These leptonic excesses are associated with the channels $tt(h \rightarrow WW^*)$, $tt(h \rightarrow ZZ^*)$ and $tt(h \rightarrow \tau\tau)$ where one of the tops decays leptonically. Within the preliminary results of Run 2 in those same leptonic channels, both ATLAS and CMS still report excesses with, for example, $\mu_{SS2l} = 1.9_{-0.8}^{+0.9}$ at CMS [3] and $\mu_{SS2l} = 4.0_{-1.7}^{+2.1}$ at ATLAS [4]. The most recent preliminary results reported by CMS in the $t\bar{t}h$ associated production searches make use of an integrated luminosity of 35.9 fb^{-1} and seem to show mixed results. In the leptonic channels (W^+W^- and ZZ channels) they still show an enhancement of 1.5 ± 0.5 times the SM prediction, with an observed (expected) significance of 3.3σ (2.5σ) obtained from combining these results with the 2015 data [5]. On the other hand in the $h \rightarrow \tau\tau$ decay channel search, a slight suppression of $0.72_{-0.53}^{+0.62}$ times the SM prediction is found, with an observed (expected) significance of 1.4σ (1.8σ) [6]. Note that, unlike the W^+W^- and ZZ decays, this last signal is sensitive to both the top-Higgs Yukawa and the τ -Higgs Yukawa couplings, and thus enhancements or suppressions are possible as long as there are variations in *either* the top quark or the τ lepton Yukawa couplings.

All measurements are still hindered by having few events so far, but nevertheless, should these tantalizing signals survive more precise measurements at higher luminosities, they will provide the much awaited signals for new physics. We summarize relevant production and decay channels in Table I with the overall combinations obtained by the ATLAS and CMS collaborations, for the signal strengths associated with each Higgs production and decay channel.

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TABLE I. Higgs signal strengths used in our analysis, from ggh and $t\bar{t}h$ production modes measured at the LHC from Run 1 (combined $\sqrt{s} = 7$ and 8 TeV results) and Run 2 ($\sqrt{s} = 13$ TeV).

Production mode	Channel	Run 1 [7]	Production mode	Channel	ATLAS Run 2	CMS Run 2
ggh	$\gamma\gamma$	$1.1^{+0.23}_{-0.22}$	ggh	$\gamma\gamma$	$0.62^{+0.30}_{-0.29}$ [8]	$0.77^{+0.25}_{-0.23}$ [9]
	WW^*	$0.84^{+0.17}_{-0.17}$		WW^*
	ZZ^*	$1.13^{+0.34}_{-0.31}$		ZZ^*	$1.34^{+0.39}_{-0.33}$ [8]	$0.96^{+0.44}_{-0.33}$ [10]
$t\bar{t}h$	$\gamma\gamma$	$2.2^{+1.6}_{-1.3}$	$t\bar{t}h$	$\gamma\gamma, b\bar{b}, \text{leptons}$	$1.8^{+0.7}_{-0.7}$ [11]	...
	$b\bar{b}$	$1.15^{+0.99}_{-0.94}$		$WW^*, ZZ^*, \tau^+\tau^-$		$2.0^{+0.8}_{-0.7}$ [3]
	WW^*	$5.0^{+2.6}_{-2.2}$				

One possibility to explain the $SS2l$ excess is that it could be due to a modified Higgs coupling to the SM top quark, resulting in an enhanced $t\bar{t}$ ($h \rightarrow \text{multileptons}$) production. A simple explanation put forward to explain this latest possible signal of physics beyond the SM has been to invoke the presence of vectorlike quarks [12]. Previous studies have adopted an effective theory approach, involving generic couplings and mixings with the third generation quarks, by which they induce modifications of the Yukawa couplings of the top and bottom quarks. The scenario has been put forward to explain deviations from SM expectations in the forward-backward asymmetry in b decays A_{FB}^b and the enhancement of the $pp \rightarrow t\bar{t}h$ cross section at the LHC. Mixing with the additional states in the bottom sector allows for a sufficiently large increase of the $Zb_R\bar{b}_R$ coupling to explain the forward-backward anomaly, as well as to imply new effects in Higgs phenomenology [13,14]. The mixing could provide a strong enhancement of the $t\bar{t}h$ Yukawa coupling, which would explain an increase of the cross section at the LHC. In this scenario, rates for the loop-induced processes stay SM-like due to either small vectorlike contributions or compensating effects between fermion mixing and loop contributions. For this to be a viable scenario, vectorlike quark masses of order 1–2 TeV are required, still safe from the LHC lower limits on their masses, $m_{\text{VLQ}} \gtrsim 800$ GeV [15].¹

In Sec. II of this article, we first revisit the SM augmented by the addition of one vectorlike quark doublet and two singlets (one toplike and one bottomlike) and review the mixing with the third family of quarks. In Sec. II A, we then show that there is a specific region of parameter space where large corrections to the top Yukawa coupling (caused by contributions from the new vectorlike quarks) do not cause large corrections to the radiative coupling of the Higgs and gluons. We show that in that limit, Higgs signal strengths can be simply parametrized in terms of four variables only, related to the top and bottom quark shifts in Yukawa couplings, which we analyze in Sec. II B, in terms

¹Alternative explanations involving supersymmetric partners have also been put forward [16], as well as early studies on the implications on some coefficients of operators of the effective Lagrangian [17].

of the parameters of the model. Based on these results, we present simple predictions between signal strength in Higgs production (through gluon fusion and $t\bar{t}h$) and decays into $\gamma\gamma, WW$ and ZZ , in Sec. II C. As seen in that subsection, large regions of parameter space are excluded by both theoretical considerations and experimental constraints. Finally, in Sec. III we describe how to reproduce our scenario within the conventional Randall-Sundrum model and then summarize our findings in Sec. IV.

II. TOP AND BOTTOM MIRRORS: A DOUBLET AND TWO SINGLETS

The simple scenario that we wish to consider contains the usual SM gauge groups and matter fields, with the addition of a vectorlike quark $SU(2)$ doublet and two vectorlike quarks $SU(2)$ singlets, one with up-type gauge charge and another with down-type gauge charge. They can be regarded as top or bottom partners as we will consider that their Yukawa couplings are large. As we will show in the next section, this structure can appear naturally in models of warped extra dimensions with the Higgs localized near the TeV brane, and with fermions in the bulk. The presence of brane kinetic terms can lower significantly the mass of some of the heavy Kaluza-Klein (KK) fermions [18]. The rest of the KK fields decouple due to their heavy masses, giving rise to something similar to the simple setup considered here. We denote $q_L^0 \equiv \begin{pmatrix} t_L^0 \\ b_L^0 \end{pmatrix}$ as the SM third generation doublet, and t_R^0 as the SM right-handed top. Using similar notation we define $Q_{L,R} \equiv \begin{pmatrix} Q_{L,R}^t \\ Q_{L,R}^b \end{pmatrix}$ as the new vectorlike quark doublet, $T_{R,L}$ as the new vectorlike up-type quark singlet, and $B_{L,R}$ as the new vectorlike down-type singlet.

In principle we should also consider the mixings with the up and charm quarks, and down and strange quarks of the SM when writing down the most general Yukawa couplings in the up and down sectors. Without any additional assumption or theory input, we should write down the most general Yukawa couplings between SM quarks and the new vectorlike quarks, leading to 5×5 fermion mass matrices. In models of warped extra dimensions with bulk

fermions, the couplings between up or charm and heavy fermions are suppressed by factors of order $\sqrt{m_f/m_t}$ with respect to the couplings to the top, which are of order 1, and so we will just neglect those terms, leading to a much simpler 3×3 fermion mass matrix (and similarly in the down sector with the bottom quark).

The mass and interaction Lagrangian in the top sector, including its Yukawa couplings with the SM-like Higgs doublet \tilde{H} can be then written as

$$\begin{aligned} \mathcal{L}_{\text{mass}} = & Y_t^0 \bar{q}_L \tilde{H} t_R + Y_{qT} \bar{q}_L \tilde{H} T_R + Y_Q \bar{Q}_L \tilde{H} t_R \\ & + Y_1 \bar{Q}_L \tilde{H} T_R + Y_2 \bar{Q}_R \tilde{H} T_L \\ & + M_Q \bar{Q}_L Q_R + M_T \bar{T}_L T_R, \end{aligned} \quad (1)$$

with a similar expression for the bottom sector. After electroweak symmetry breaking, the Yukawa couplings induce off-diagonal terms into the fermion mass matrix. In the basis defined by the vectors $(\bar{q}_L, \bar{Q}_L, \bar{T}_L)$ and (t_R, Q_R, T_R) we can write the heavy quark mass matrix as

$$\mathbf{M}_t = \begin{pmatrix} vY_t^0 & 0 & vY_{qT} \\ vY_{Qt} & M_Q & vY_1^t \\ 0 & vY_2^t & M_T \end{pmatrix}, \quad (2)$$

where, in general, all entries are complex² and where v is the Higgs vacuum expectation value (VEV). We can also express the bottom sector heavy quark mass matrix as

$$\mathbf{M}_b = \begin{pmatrix} vY_b^0 & 0 & vY_{qB} \\ vY_{Qb} & M_Q & vY_1^b \\ 0 & vY_2^b & M_B \end{pmatrix}, \quad (3)$$

where again all entries can be complex and the value of M_Q is the same in both \mathbf{M}_t and \mathbf{M}_b . The associated top and bottom Yukawa coupling matrices are

$$\tilde{\mathbf{Y}}_t = \begin{pmatrix} Y_t^0 & 0 & Y_{qT} \\ Y_{Qt} & 0 & Y_1^t \\ 0 & Y_2^t & 0 \end{pmatrix} \quad \text{and} \quad \tilde{\mathbf{Y}}_b = \begin{pmatrix} Y_b^0 & 0 & Y_{qB} \\ Y_{Qb} & 0 & Y_1^b \\ 0 & Y_2^b & 0 \end{pmatrix}. \quad (4)$$

The mass matrices \mathbf{M}_t and \mathbf{M}_b are diagonalized by biunitary transformations, $\mathbf{V}_{tL}^\dagger \mathbf{M}_t \mathbf{V}_{tR} = \mathbf{M}_t^{\text{diag}}$ and $\mathbf{V}_{bL}^\dagger \mathbf{M}_b \mathbf{V}_{bR} = \mathbf{M}_b^{\text{diag}}$. At the same time, the Higgs Yukawa couplings are obtained after transforming the Yukawa matrices into the physical basis, $\mathbf{V}_{tL}^\dagger \tilde{\mathbf{Y}}_t \mathbf{V}_{tR} = \mathbf{Y}_t^{\text{phys}}$ and $\mathbf{V}_{bL}^\dagger \tilde{\mathbf{Y}}_b \mathbf{V}_{bR} = \mathbf{Y}_b^{\text{phys}}$.

²We can eliminate five phases from the mass matrix \mathbf{M}_t through phase redefinitions. We keep the notation general, since the phases regroup together easily in all the expressions.

A. Higgs production in the brane Higgs limit

In the physical basis, the top quark mass and the top Yukawa coupling (the first entries in the physical mass matrix and the physical Yukawa matrix) are not related anymore by the SM relationship $m_t^{\text{phys}} = v y_t^{\text{SM}}$ [19] (with v normalized to $v = 174$ GeV for simplicity). The same goes for the bottom quark, and we thus define the shifts, δy_t and δy_b , between the SM and the physical Yukawa couplings, due to the diagonalization, as

$$y_t^{\text{phys}} = y_t^{\text{SM}} - \delta y_t, \quad (5)$$

and

$$y_b^{\text{phys}} = y_b^{\text{SM}} - \delta y_b. \quad (6)$$

Later we will give an approximate expression of these shifts in terms of the model parameters of Eqs. (2) and (3). But before that, we consider the radiative coupling of the Higgs to gluons. This coupling depends on the physical Yukawa couplings y_{nn} of all the fermions running in the loop and on their physical masses m_n . The real and imaginary parts of the couplings (the scalar and pseudoscalar parts) contribute to the cross section through different loop functions, $A_{1/2}^S$ and $A_{1/2}^P$, as they generate the two operators $hG_{\mu\nu}G^{\mu\nu}$ and $hG_{\mu\nu}\tilde{G}^{\mu\nu}$.

The cross section, depicted in Fig. 1, is

$$\sigma_{gg \rightarrow h} = \frac{\alpha_s^2 m_h^2}{576\pi} [|c_{ggh}^S|^2 + |c_{ggh}^P|^2] \delta(s - m_h^2) \quad (7)$$

where

$$\begin{aligned} c_{ggh}^S &= \sum_{n=1}^3 \text{Re} \left(\frac{y_{nn}}{m_n} \right) A_{1/2}^S(\tau_f) \quad \text{and} \\ c_{ggh}^P &= \sum_{n=1}^3 \text{Im} \left(\frac{y_{nn}}{m_n} \right) A_{1/2}^P(\tau_f) \end{aligned} \quad (8)$$

with $\tau = m_h^2/4m_n^2$ and with the loop functions $A_{1/2}^S(\tau)$ and $A_{1/2}^P(\tau)$ as defined in [20]. Note that we use a normalization of the loop functions such that for very heavy quarks with masses m_n much greater than the Higgs mass m_h (i.e. when τ is very small) they behave asymptotically as $\lim_{\tau \rightarrow 0} A_{1/2}^S = 1$ and $\lim_{\tau \rightarrow 0} A_{1/2}^P = 3/2$. On the other hand, for light quarks (all the SM quarks except top and to some extent, bottom), the loop functions essentially vanish since $\lim_{\tau \rightarrow \infty} A_{1/2}^S = \lim_{\tau \rightarrow \infty} A_{1/2}^P = 0$, and we thus neglect contributions from the light SM quarks and consider only the effect of the top, the bottom and the four remaining physical heavy quarks.

The amplitudes c_{ggh}^S and c_{ggh}^P can then be written in terms of traces involving the fermion mass and Yukawa matrices

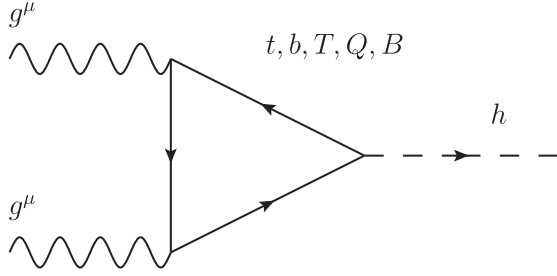


FIG. 1. Feynman diagram for the production cross section $gg \rightarrow h$ in a setup with new vectorlike fermions Q, T and B .

involving top and vectorlike up quarks, and bottom and any vectorlike down quarks, \mathbf{M}_i and \mathbf{Y}_i with $i = t, b$, so that we obtain

$$c_{ggh}^S = \sum_n \text{Re} \left(\frac{y_{nn}^u}{m_n^u} \right) + \sum_n \text{Re} \left(\frac{y_{nn}^d}{m_n^d} \right) - \text{Re} \left(\frac{y_b}{m_b} \right) + \text{Re} \left(\frac{y_b}{m_b} \right) A_{1/2}^S(\tau_b) \quad (9)$$

where we have added and subtracted the bottom quark loop contribution in order to keep the dependence in $A_{1/2}(\tau_b)$, and with a similar expression holding for c_{ggh}^P . We evaluate exactly the sums in the top and bottom sectors and find

$$\sum_n \left(\frac{y_{nn}^u}{m_n^u} \right) = \frac{1}{v} \frac{1 + 3\varepsilon_{Q_t} \varepsilon_T \frac{|Y_2^t|}{|Y_1^t|} e^{i\theta_2^t} \left(1 - e^{i\theta_1^t} \frac{|Y_1^t| |Y_1^0|}{|Y_{Q_t}| |Y_{qT}|} \right)}{1 + \varepsilon_{Q_t} \varepsilon_T \frac{|Y_2^t|}{|Y_1^t|} e^{i\theta_2^t} \left(1 - e^{i\theta_1^t} \frac{|Y_1^t| |Y_1^0|}{|Y_{Q_t}| |Y_{qT}|} \right)}, \quad (10)$$

and

$$\sum_n \left(\frac{y_{nn}^d}{m_n^d} \right) = \frac{1}{v} \frac{1 + 3\varepsilon_{Q_b} \varepsilon_B \frac{|Y_2^b|}{|Y_1^b|} e^{i\theta_2^b} \left(1 - e^{i\theta_1^b} \frac{|Y_1^b| |Y_1^0|}{|Y_{Q_b}| |Y_{qB}|} \right)}{1 + \varepsilon_{Q_b} \varepsilon_B \frac{|Y_2^b|}{|Y_1^b|} e^{i\theta_2^b} \left(1 - e^{i\theta_1^b} \frac{|Y_1^b| |Y_1^0|}{|Y_{Q_b}| |Y_{qB}|} \right)}, \quad (11)$$

where we have defined the small parameters $\varepsilon_T = \frac{v|Y_{qT}|}{|M_T|}$, $\varepsilon_B = \frac{v|Y_{qB}|}{|M_B|}$, $\varepsilon_{Q_t} = \frac{v|Y_{Q_t}|}{|M_{Q_t}|}$ and $\varepsilon_{Q_b} = \frac{v|Y_{Q_b}|}{|M_{Q_b}|}$, and with the relative phases θ_1^i and θ_2^i defined as $\theta_1^t = \text{Arg} \left(\frac{Y_{Q_t}^0 Y_1^t}{Y_{qT} Y_1^0} \right)$ and $\theta_2^t = \text{Arg} \left(\frac{Y_{qT} Y_{Q_t} Y_2^t}{M_T M_{Q_t} Y_1^0} \right)$, with similar definitions for θ_1^b and θ_2^b . In the SM limit, the expression in Eq. (9) should tend to $\sim \frac{1}{v} (1 + A_{1/2}^S(\tau_b))$, and so if one wished to limit the contribution coming from the top partners to Higgs production (in gluon fusion), we must reduce/eliminate these corrections. We note the following observations.

- (i) Of course for heavier and heavier vectorlike fermions, the parameters ε_i become more and more suppressed, and thus we can smoothly recover the SM limit, but the new physics effect will decouple

from everywhere else (and in particular the top Yukawa quark will also tend to its SM value).

- (ii) It might also be possible to reduce the couplings $|Y_2^t|$ or $|Y_2^b|$, but then this will also affect the physical top or bottom Yukawa coupling shifts, and in particular no enhancement in the top quark Yukawa coupling will be possible (although suppression might still be possible), as we will show later.
- (iii) Another interesting possibility would be to set the overall phase of the correction term in Eqs. (10) or (11) to be $\pi/2$ so that the real part vanishes (in general, we expect that the real part would dominate the overall corrections, at least for $\varepsilon_i < 1$). This possibility might limit the amount of enhancement in the top Yukawa coupling, since that correction also depends on the phase θ_2 . Again when we compute the approximate expression of the Yukawa couplings shift, we will see that the phase θ_2 should be close to 0 to yield an enhancement in the top Yukawa coupling.
- (iv) In our considerations, we will impose a seemingly contrived constraint on the model parameters, which we call the ‘‘brane Higgs limit,’’ such that

$$\det \tilde{\mathbf{Y}}_t = \det \tilde{\mathbf{Y}}_b = 0, \quad (12)$$

with the matrices $\tilde{\mathbf{Y}}_t$ and $\tilde{\mathbf{Y}}_b$ defined in Eq. (4). This constraint implies that

$$Y_2^t \left(1 - e^{i\theta_1^t} \frac{|Y_1^t| |Y_1^0|}{|Y_{Q_t}| |Y_{qT}|} \right) = 0, \quad (13)$$

and thus ensures that the top sector contribution to Higgs production, given in Eq. (10), gives the same result as the SM top quark contribution to the same process. The vanishing determinant condition could come from a specific flavor structure in the Yukawa matrix, emerging for example from democratic textures, etc. We will show in the next section that the flavor structure required can also be obtained in models of extra dimensions, so that the cancellation in Eq. (13) is satisfied exactly if the scenario arises out of the usual Randall-Sundrum warped extra-dimensional scenario with matter fields in bulk. It is necessary, though, that the Higgs be sufficiently localized towards the brane and that the KK modes of the top quark (and bottom quark) be much lighter than the KK partners of the up and charm quarks (and the down and strange quarks). We will then refer to the vectorlike partners of the top and bottom quarks throughout as KK partners, and we return to this scenario in Sec. III.

Therefore we work in the brane Higgs limit of the general parameter space. In the down sector, we also have

$$Y_2^b \left(1 - e^{i\theta_b^1} \frac{|Y_1^b| |Y_b^0|}{|Y_{Qb}| |Y_{qB}|} \right) = 0, \quad (14)$$

so that we have

$$\sum_n \left(\frac{y_{nn}^u}{m_n^u} \right) = \sum_n \left(\frac{y_{nn}^d}{m_n^d} \right) = \frac{1}{v}. \quad (15)$$

This means that now we can write the ggh couplings as

$$c_{ggh}^S = \frac{1}{v} (1 + A_{1/2}^S(\tau_b)) + \text{Re} \left(\frac{\delta y_b}{m_b} \right) (1 - A_{1/2}^S(\tau_b)), \quad (16)$$

and

$$c_{ggh}^P = \text{Im} \left(\frac{\delta y_b}{m_b} \right) \left(\frac{3}{2} - A_{1/2}^P(\tau_b) \right), \quad (17)$$

where we have used the definitions of the Yukawa coupling shifts in Eqs. (5) and (6). Evaluating the values for the bottom quark loop functions $A_{1/2}^{S,P}(\tau_b)$ we obtain

$$\frac{\sigma_{gg \rightarrow h}}{\sigma_{gg \rightarrow h}^{\text{SM}}} = \frac{\Gamma_{h \rightarrow gg}}{\Gamma_{h \rightarrow gg}^{\text{SM}}} = (1 + \Delta_{gg}), \quad (18)$$

where the correction term is

$$\begin{aligned} \Delta_{gg} &= 2.13v \left(\text{Re} \frac{\delta y_b}{m_b} \right) + 1.13v^2 \left(\text{Re} \frac{\delta y_b}{m_b} \right)^2 \\ &+ 2.51v^2 \left(\text{Im} \frac{\delta y_b}{m_b} \right)^2. \end{aligned} \quad (19)$$

This result links in a simple and nontrivial way Higgs production through gluon fusion to the bottom quark Yukawa coupling (or more precisely to its relative shift $v\delta y_b/m_b$). In a similar fashion we can also obtain the correction to the Higgs decay into $\gamma\gamma$ in the brane Higgs limit, since the fermion loop is the same as the gluon fusion loop (although there is an additional W loop contribution in this case). We obtain

$$\frac{\Gamma_{h \rightarrow \gamma\gamma}}{\Gamma_{h \rightarrow \gamma\gamma}^{\text{SM}}} = (1 + \Delta_{\gamma\gamma}), \quad (20)$$

with

$$\begin{aligned} \Delta_{\gamma\gamma} &= -0.14v \left(\text{Re} \frac{\delta y_b}{m_b} \right) + 0.005v^2 \left(\text{Re} \frac{\delta y_b}{m_b} \right)^2 \\ &+ 0.01v^2 \left(\text{Im} \frac{\delta y_b}{m_b} \right)^2, \end{aligned} \quad (21)$$

and where we took the SM loop contributions to be $|c_{\gamma\gamma}| = |-7A_1(\tau_W) + 16/9A_{1/2}^S(\tau_t) + 4/9A_{1/2}^S(\tau_b)| \approx 6.53$,

with the W -loop function $A_1(\tau_W)$ [20] and the fermion loop function $A_{1/2}^S(\tau_f)$ normalized so that $\lim_{\tau \rightarrow 0} A_i(\tau) = 1$. Finally, from Eqs. (5) and (6) we can now write

$$\begin{aligned} \frac{\sigma_{pp \rightarrow tth}}{\sigma_{pp \rightarrow tth}^{\text{SM}}} &= \left| \frac{y_t}{y_{\text{SM}}} \right|^2 = (1 + \Delta_{tt}) \quad \text{and} \\ \frac{\Gamma_{h \rightarrow bb}}{\Gamma_{h \rightarrow bb}^{\text{SM}}} &= \left| \frac{y_b}{y_{\text{SM}}} \right|^2 = (1 + \Delta_{bb}), \end{aligned} \quad (22)$$

where

$$\begin{aligned} \Delta_{tt} &= -2v \text{Re} \left(\frac{\delta y_t}{m_t} \right) + v^2 \left| \frac{\delta y_t}{m_t} \right|^2 \quad \text{and} \\ \Delta_{bb} &= -2v \text{Re} \left(\frac{\delta y_b}{m_b} \right) + v^2 \left| \frac{\delta y_b}{m_b} \right|^2. \end{aligned} \quad (23)$$

We are interested in studying the dependence on the Yukawa shifts δy_t and δy_b of the signal strengths

$$\begin{aligned} \mu_{ggh}^{ii} &= \frac{\sigma(gg \rightarrow h) \text{Br}(h \rightarrow ii)}{\sigma_{\text{SM}}(gg \rightarrow h) \text{Br}^{\text{SM}}(h \rightarrow ii)} \\ &= \frac{\sigma(gg \rightarrow h)}{\sigma_{\text{SM}}(gg \rightarrow h)} \frac{\Gamma(h \rightarrow ii)}{\Gamma_{\text{SM}}(h \rightarrow ii)} \frac{\Gamma_{\text{SM}}^{\text{tot}}(h)}{\Gamma^{\text{tot}}(h)}, \end{aligned} \quad (24)$$

with a similar expression for μ_{tth}^{ii} (and using the small width approximation). In these expressions, the ratio of total Higgs widths can be written as

$$\frac{\Gamma_{\text{SM}}^{\text{tot}}(h)}{\Gamma^{\text{tot}}(h)} = \frac{1}{(1 + \text{Br}_{h \rightarrow bb}^{\text{SM}} \Delta_{bb} + \text{Br}_{h \rightarrow gg}^{\text{SM}} \Delta_{gg} + \text{Br}_{h \rightarrow \gamma\gamma}^{\text{SM}} \Delta_{\gamma\gamma})}, \quad (25)$$

where, taking into account numerical values for the SM Higgs branching ratios, gives simply

$$\frac{\Gamma_{\text{SM}}^{\text{tot}}(h)}{\Gamma^{\text{tot}}(h)} \approx \frac{1}{1 + 0.58\Delta_{bb} + 0.086\Delta_{gg}}, \quad (26)$$

and where we have dropped the dependence in $\Delta_{\gamma\gamma}$ as it is highly suppressed.

With all these ingredients, we find the $t\bar{t}h$ production- and decay strengths

$$\mu_{t\bar{t}h}^{VV} = \frac{(1 + \Delta_{tt})}{(1 + 0.58\Delta_{bb} + 0.086\Delta_{gg})}, \quad (27)$$

$$\mu_{t\bar{t}h}^{bb} = \mu_{t\bar{t}h}^{VV} (1 + \Delta_{bb}), \quad (28)$$

$$\mu_{t\bar{t}h}^{\gamma\gamma} = \mu_{t\bar{t}h}^{VV} (1 + \Delta_{\gamma\gamma}), \quad (29)$$

as well as the ggh strengths

$$\mu_{ggh}^{VV} = \frac{(1 + \Delta_{gg})}{(1 + 0.58\Delta_{bb} + 0.086\Delta_{gg})}, \quad (30)$$

$$\mu_{ggh}^{\gamma\gamma} = \mu_{ggh}^{VV}(1 + \Delta_{\gamma\gamma}), \quad (31)$$

with the corrections Δ_{ii} depending only on top or bottom quark Yukawa coupling shifts

$$\Delta_{tt} = -2v\text{Re}\left(\frac{\delta y_t}{m_t}\right) + v^2\left|\frac{\delta y_t}{m_t}\right|^2 \quad (32)$$

$$\Delta_{bb} = -2v\text{Re}\left(\frac{\delta y_b}{m_b}\right) + v^2\left|\frac{\delta y_b}{m_b}\right|^2 \quad (33)$$

$$\begin{aligned} \Delta_{gg} = & 2.13v\left(\text{Re}\frac{\delta y_b}{m_b}\right) + 1.13v^2\left(\text{Re}\frac{\delta y_b}{m_b}\right)^2 \\ & + 2.51v^2\left(\text{Im}\frac{\delta y_b}{m_b}\right)^2 \end{aligned} \quad (34)$$

$$\begin{aligned} \Delta_{\gamma\gamma} = & -0.14v\left(\text{Re}\frac{\delta y_b}{m_b}\right) + 0.005v^2\left(\text{Re}\frac{\delta y_b}{m_b}\right)^2 \\ & + 0.01v^2\left(\text{Im}\frac{\delta y_b}{m_b}\right)^2. \end{aligned} \quad (35)$$

B. Yukawa coupling shifts

As indicated before, the mass matrices \mathbf{M}_t and \mathbf{M}_b from Eqs. (2) and (3) are diagonalized by biunitary transformations, $\mathbf{V}_{tL}\dagger\mathbf{M}_t\mathbf{V}_{tR} = \mathbf{M}_t^{\text{diag}}$ and $\mathbf{V}_{bL}\dagger\mathbf{M}_b\mathbf{V}_{bR} = \mathbf{M}_b^{\text{diag}}$. In order to obtain simple analytical expressions for the Yukawa couplings emerging after the diagonalization, we expand the unitary matrices $\mathbf{V}_{t,bL}$ and $\mathbf{V}_{t,bR}$ in powers of $\varepsilon \sim v/M$, where v is the Higgs VEV and M represents the vectorlike masses M_Q , M_T or M_B .

In this approximation we can obtain the lightest mass eigenvalues (the top quark and the bottom quark masses) as well as the physical $t\bar{t}h$ Yukawa coupling and the $b\bar{b}h$ Yukawa coupling.

This yields the relative deviation between the physical Yukawa couplings y_t^{phys} and y_b^{phys} , and the SM Yukawa couplings, defined as $y_t^{\text{SM}} = m_t^{\text{phys}}/v$ and $y_b^{\text{SM}} = m_b^{\text{phys}}/v$. In terms of the mass matrix parameters from Eqs. (2) and (3), we obtain

$$\frac{\delta y_t}{y_t^{\text{SM}}} = \varepsilon_T^2 + \varepsilon_{Q_t}^2 - 2\varepsilon_T\varepsilon_{Q_t}\frac{|Y_2|}{|Y_t^0|}e^{i\theta_2} + \mathcal{O}(\varepsilon^4), \quad (36)$$

and similarly for the bottom quark

$$\frac{\delta y_b}{y_b^{\text{SM}}} = \varepsilon_B^2 + \varepsilon_{Q_b}^2 - 2\varepsilon_B\varepsilon_{Q_b}\frac{|Y_2^b|}{|Y_b^0|}e^{i\theta_2^b} + \mathcal{O}(\varepsilon^4). \quad (37)$$

As previously, $\varepsilon_T = \frac{v|Y_{qT}|}{|M_T|}$, $\varepsilon_B = \frac{v|Y_{qB}|}{|M_B|}$, $\varepsilon_{Q_t} = \frac{v|Y_{Q_t}|}{|M_Q|}$ and $\varepsilon_{Q_b} = \frac{v|Y_{Q_b}|}{|M_Q|}$ and the relative phases θ_2^t and θ_2^b are given

as $\theta_2^t = \text{Arg}\left(\frac{Y_{qT}Y_{Q_t}Y_t^0}{M_T M_Q Y_t^0}\right)$ and $\theta_2^b = \text{Arg}\left(\frac{Y_{qB}Y_{Q_b}Y_b^0}{M_B M_Q Y_b^0}\right)$. Note that these perturbative expressions are only valid for $\varepsilon_i < 1$. Nevertheless they are very useful in identifying limits and parameter behavior, and moreover the limit $\varepsilon_i < 1$ is the natural one as the top and bottom KK partners are expected to be heavy enough to make the expansions converge. The first two terms of both expressions always yield a suppression in the physical Yukawa coupling strength, irrespective of the phases within the original fermion mass matrices. However, the third term, proportional to Y_2 , could induce an overall enhancement of the top Yukawa coupling or of the bottom Yukawa coupling, when the phases $\theta_2^{t,b}$ are such that $-\pi/2 < \theta_2^{t,b} < \pi/2$, and for sufficiently high values of Y_2 . An enhancement effect would be maximal when $\theta_2^{t,b} = 0$.

Note that the top/bottom mirror sector, even though essentially decoupled from the light quarks, should still have some impact in the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix. In the same perturbative limit used to obtain Yukawa shifts, we can also obtain an approximation to the corrections on V_{tb} , due to the presence of the top/bottom vectorlike mirrors. The value is shifted as

$$|V_{tb}| \simeq 1 - \frac{1}{2}|V_{cb}|^2 - \frac{1}{2}|V_{ub}|^2 - \frac{1}{2}(\varepsilon_T - \varepsilon_B)^2, \quad (38)$$

where the first two terms represent the usual SM CKM unitarity constraint, and the last term is the new contribution (where we have eliminated the relative phases between Y_{qT} and M_T , and between Y_{qB} and M_B through a phase redefinition). The current Tevatron and LHC average on V_{tb} , coming from single top production is $V_{tb} = 1.009 \pm 0.031$ [21], which gives a lowest bound of about $V_{tb} \sim 0.97$. That means that the corrections from our scenario should be limited to about

$$\frac{1}{2}(\varepsilon_T - \varepsilon_B)^2 \lesssim (0.15)^2, \quad (39)$$

requiring that $Yv/M \lesssim 0.15$ or $M \gtrsim 1$ TeV, unless a strong cancellation between the top and bottom terms happens. In a similar way, the rest of the third row and third column CKM mixing angles V_{ub} , V_{cb} , V_{td} and V_{ts} will receive corrections producing deviations on the usual SM unitarity relations. For example we have

$$(1 - \varepsilon_B^2)|V_{cb}|^2 \simeq (1 - |V_{cd}|^2 - |V_{cs}|^2) \quad (40)$$

so that imposing the experimental uncertainties in $|V_{cd}|$, $|V_{cs}|$ and $|V_{cb}|$ [21], we find that

$$\varepsilon_B^2 \lesssim (0.44)^2. \quad (41)$$

This is a slightly less constraining bound on the vectorlike sector, compared to the one in Eq. (39). A thorough full fit analysis on CKM unitarity is beyond the scope of this paper, although should the $t\bar{t}h$ signal survive the higher luminosity

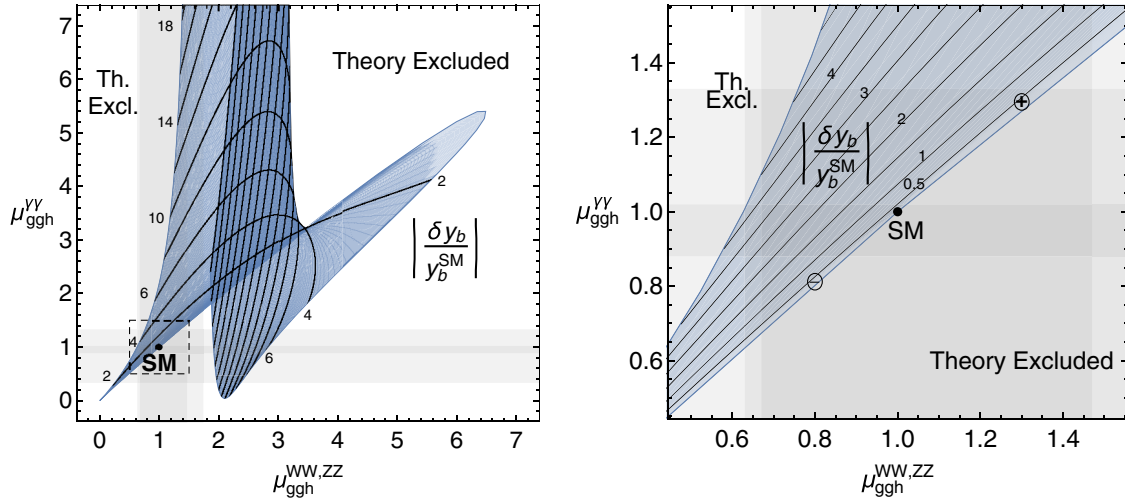


FIG. 2. Contours of the bottom quark Yukawa correction $|\frac{\delta y_b}{y_b^{\text{SM}}}|$ with respect to the gluon fusion signal strengths μ_{ggh}^{VV} and $\mu_{ggh}^{\gamma\gamma}$. The right panel zooms in on the region marked by a dashed box in the left panel. The horizontal and vertical gray bands represent the experimental bounds set by the LHC Run 1 (darker) and the preliminary data from LHC Run 2 (lighter). The “Theory Excluded” regions are points excluded by the brane Higgs limit constraint. Each contour is traced by varying the phase of δy_b and we include two parameter space points as example limits, marked by a \oplus and a \ominus , representing, respectively, an overall enhancement or suppression with respect to the SM predictions.

data, with improved constraints in the Higgs sector, such a study might become useful.³

Finally, flavor mixing between vectorlike quarks and the third generation can affect other flavor observables, particularly in B -physics. This was extensively discussed in the literature [22], where a suppression of $\text{BR}(B_s \rightarrow \mu\bar{\mu})$ and an enhancement in $\text{BR}(B_d \rightarrow \mu\bar{\mu})$ are shown to be most likely. Here we will simply ask that the mixing in the bottom sector remains small; i.e. we should consider parameter space points where the shift in the bottom quark Yukawa coupling δy_b is small. Again, a full flavor analysis should be addressed if the $t\bar{t}h$ enhanced signal is confirmed.

C. Higgs phenomenology

As we have seen earlier, the brane Higgs limit condition is quite predictive, and easily falsifiable in the near future from LHC Higgs data. The first important point is that within our minimal general setup, all signal strengths associated with the Higgs will deviate from the SM values *only* due to shifts in the top and bottom quark Yukawa couplings. This means that ratios of Higgs signal strengths involving electroweak production processes, and decays through the same channels “ ii ,” should be equal to 1, i.e.

$$\frac{\mu_{\text{VBF}}^{ii}}{\mu_{W_h}^{ii}} = \frac{\mu_{\text{VBF}}^{ii}}{\mu_{Z_h}^{ii}} = 1. \quad (42)$$

Also signal strengths involving decays into WW should be equal to signals with decays into ZZ , i.e.

$$\frac{\mu_{ggh}^{WW}}{\mu_{ggh}^{ZZ}} = \frac{\mu_{t\bar{t}h}^{WW}}{\mu_{t\bar{t}h}^{ZZ}} = \frac{\mu_{Vh}^{WW}}{\mu_{Vh}^{ZZ}} = \frac{\mu_{\text{VBF}}^{WW}}{\mu_{\text{VBF}}^{ZZ}} = 1. \quad (43)$$

These are strong model dependent predictions, likely testable at the present Run 2 at the LHC.

Now, more specific to our setup, and as seen from Eqs. (32)–(35), the corrections to all of the Higgs signal strengths depend only on four parameters, i.e. the absolute values of the relative top and bottom Yukawa coupling deviations $|\delta y_t|$ and $|\delta y_b|$, and their two phases. Moreover, only the $t\bar{t}h$ signal strengths depend on all four parameters. We thus start exploring the dominant Higgs production mechanism, the gluon fusion process, paying particular attention to the signal strengths $\mu_{ggh}^{\gamma\gamma}$ and $\mu_{ggh}^{WW,ZZ}$. These depend only on the deviation of the bottom quark coupling (magnitude and phase). It is therefore possible to study the relationship between these two signal strengths, for different values of δy_b . This is plotted in Fig. 2, where we show that only a specific region in the $(\mu_{ggh}^{\gamma\gamma}, \mu_{ggh}^{WW,ZZ})$ plane is allowed, due to the brane Higgs limit constraint. The horizontal and vertical gray bands correspond to limits set by LHC Run 1 and preliminary LHC Run 2 data, as summarized in Table I. In the right panel of that figure, we zoom in on the square enclosed by dashed lines in the left panel to consider signal strengths close to the SM model value, and we can see the region where $\mu_{ggh}^{\gamma\gamma} < \mu_{ggh}^{WW,ZZ}$ is not allowed, thus providing a very simple and strong prediction of the scenario. Corrections in the direction $\mu_{ggh}^{\gamma\gamma} > \mu_{ggh}^{WW,ZZ}$ are possible, but require increasingly large deviations in the bottom Yukawa coupling. For relatively

³Note that we are still assuming that first and second quark generations have highly suppressed Yukawa couplings with the top and bottom vectorlike partners.

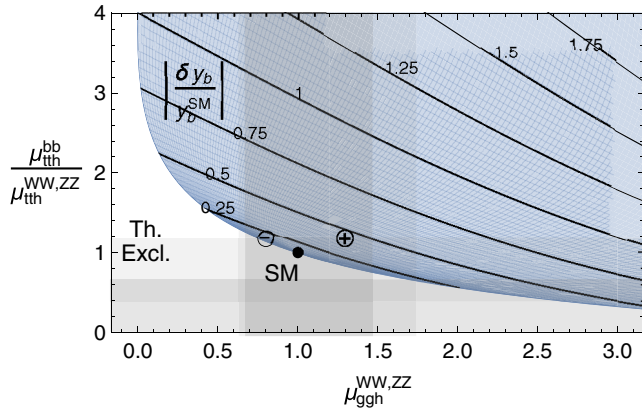


FIG. 3. Contours of the bottom quark Yukawa correction $|\frac{\delta y_b}{y_b^{\text{SM}}}|$ with respect to the gluon fusion signal strength μ_{ggh}^{VV} and the ratio of the $t\bar{t}h$ signal strengths $\frac{\mu_{t\bar{t}h}^{bb}}{\mu_{t\bar{t}h}^{VV}}$. The horizontal and vertical gray bands represent the experimental bounds set by the LHC Run 1 (darker) and the preliminary data from LHC Run 2 (lighter). The “Th. Excl.” region comprises all points excluded by the brane Higgs limit constraint. Each contour is traced by varying the phase of δy_b and we included the same two parameter space points, marked by \oplus and \ominus , as in the previous figure.

small values of δy_b , one can still obtain important deviations in the signals if one moves along the $\mu_{ggh}^{\gamma\gamma} \approx \mu_{ggh}^{WW,ZZ}$ diagonal line. For future use, we choose two points along that line, close to the boundaries set by the LHC constraints. We denote them with a \oplus and \ominus , and they represent either an overall enhancement in the ggh signal

strengths, or an overall suppression, with respect to the SM predictions.

Once the gluon fusion signals have been fixed, we can study the effect on other signal strengths which receive corrections only through the bottom quark Yukawa coupling. In particular we can explore how the ratios $\frac{\mu_{t\bar{t}h}^{bb}}{\mu_{t\bar{t}h}^{VV}} = \frac{\mu_{Vh}^{bb}}{\mu_{Vh}^{VV}}$ behave as a function of μ_{ggh}^{VV} (all top quark Yukawa dependence cancels out in the ratio). This is shown in Fig. 3, where we consider variations of the ratio $\frac{\mu_{t\bar{t}h}^{bb}}{\mu_{t\bar{t}h}^{VV}}$ (with the corresponding LHC bounds represented by the horizontal gray bands), with respect to the gluon fusion strength $\mu_{t\bar{t}h}^{bb}$. As we can see, the current experimental data tend to prefer values for that ratio close to 1 or less, therefore putting some pressure on the allowed parameter space. We can see that if the $\mu_{t\bar{t}h}^{bb}$ signal strength is smaller than the $\mu_{t\bar{t}h}^{VV}$ one (both in $t\bar{t}h$ production), then the data prefer a slight enhancement in the gluon fusion production strength. Conversely, if the $\mu_{t\bar{t}h}^{bb}$ signal is enhanced, then gluon fusion signals should be suppressed. Overall, the deviations on the bottom quark Yukawa coupling must be kept small, unless the $\mu_{t\bar{t}h}^{bb}$ signal happens to be very much larger than the $\mu_{t\bar{t}h}^{VV}$ signal. The chosen example points \oplus and \ominus stay within a ratio of $\mu_{t\bar{t}h}$ production signals close to 1.

Once we analyzed the restriction on the deviations from bottom quark Yukawa couplings, we can investigate the signals that do depend on the top Yukawa coupling deviations. In Fig. 4 we choose to study the variation of $\mu_{t\bar{t}h}^{VV}$ with respect to the top Yukawa deviation $|\delta y_t|$. The rest

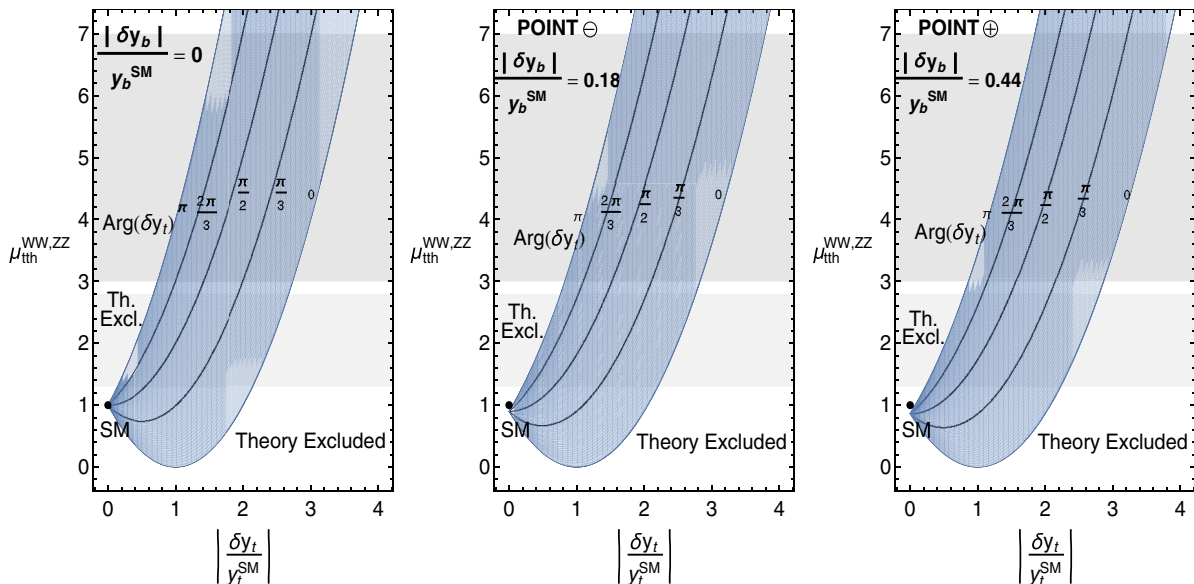


FIG. 4. Higgs signal strength $\mu_{t\bar{t}h}^{VV}$ with respect to the top quark Yukawa coupling $\mu_{t\bar{t}h}^{VV}$ correction $|\frac{\delta y_t}{y_t^{\text{SM}}}|$. The contours trace points with constant value of the phase of δy_t , and the horizontal gray bands represent the experimental bounds set by the LHC Run 1 (darker) and the preliminary data from LHC Run 2 (lighter). The “Theory Excluded” regions are excluded by the brane Higgs limit constraint. In the left panel, we consider a parameter space point where the bottom Yukawa coupling is SM-like. In the central panel the bottom quark Yukawa coupling is corrected by 18% (point \ominus) and the right panel has a bottom quark Yukawa correction of 44% (point \oplus).

of the $t\bar{t}h$ signal strengths can be obtained from ratios of other Higgs production signals strengths, since for example $\mu_{t\bar{t}h}^{\gamma\gamma} = \frac{\mu_{\text{qgh}}^{\gamma\gamma}}{\mu_{\text{qgh}}^{\text{VV}}} \mu_{t\bar{t}h}^{\text{VV}}$. We fix the values of the bottom quark Yukawa coupling in three limits, i.e. when y_b is SM-like ($\delta y_b = 0$), when it has a 18% correction ($|\delta y_b/y_b^{\text{SM}}| = 0.18$, corresponding to the point \ominus), and when it has a 44% correction ($|\delta y_b/y_b^{\text{SM}}| = 0.44$, corresponding to the point \oplus). As can be seen in Fig. 4, for moderate values of δy_b there is very mild dependence on δy_b , so that the three panels show very similar behavior of the signal strength as a function of the deviations in top quark Yukawa couplings. The parameter space region is a diagonal band, and we show contours of the phase of the top Yukawa shift δy_t , tracing the band diagonally. The dependence is very sensitive to variations in the phase of the shift of the top Yukawa coupling. We can clearly see that if the magnitude of the top Yukawa deviation is less than 1 (the natural expectation for heavy KK top partners), in order to obtain a signal enhancement (as hinted at by LHC data), the phase must be close to π . This is in agreement with the perturbative expressions obtained earlier for the Yukawa shifts and it corresponds to values of the mass matrix phase θ'_2 close to 0.

III. BRANE HIGGS LIMIT IN RANDALL-SUNDRUM MODELS

In this section we describe briefly how to reproduce the previous phenomenological scenario within the context of the Randall-Sundrum model (RS) [23]. Consider a sector of a 5D scenario with a 5D top quark, i.e. a doublet fermion $Q(x, y)$ and a singlet $T(x, y)$ defined by the following action:

$$\begin{aligned}
S = & -i \int d^4x dy \sqrt{-g} [\bar{Q} \not{D} Q + c_q \sigma' \bar{Q} Q + \bar{T} \not{D} T + c_t \sigma' \bar{T} T \\
& + \delta(y - y_1) (\alpha_{qL} \bar{Q}_L \not{\partial} Q_L + \alpha_{qR} \bar{Q}_R \not{\partial} Q_R \\
& + \alpha_{tL} \bar{T}_L \not{\partial} T_L + \alpha_{tR} \bar{T}_R \not{\partial} T_R) \\
& + \delta(y - y_1) (Y_1 H \bar{Q}_L T_R + Y_2 H \bar{Q}_R T_L + \text{H.c.})] \quad (44)
\end{aligned}$$

where $\not{D} = \gamma_A e_A^M D_M$ and $\not{\partial} = \gamma_a e_a^\mu \partial_\mu$, with γ_A the 5D gamma matrices, e_A^M the vielbein, and $D_M = (\partial_M + \Gamma_M)$ the 5D covariant derivative involving the spin connection Γ_M , with $\Gamma_\mu = \frac{1}{2} \gamma_5 \gamma_\mu \sigma'$ and $\Gamma_5 = 0$. The fifth dimension is understood as an interval, with the boundary terms fixing the boundary conditions of the fields. We have added a set of fermion kinetic terms localized at the boundary $y = y_1$. Other boundary fermion kinetic terms, involving y -derivatives, are allowed but we leave them out for simplicity. Also note that we should only consider positive brane kinetic term coefficients α_i , in order to avoid tachyons and/or ghosts [24,25].

We also consider Higgs localized Yukawa couplings on the same boundary. Note that the doublet Q is vectorlike in 5D, and we define $Q = Q_L + Q_R$ and $T = T_L + T_R$ where Q_L, T_L and Q_R, T_R are the left- and right-handed components.

The background spacetime metric is assumed to take the form

$$ds^2 = e^{-2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2 \quad (45)$$

where $\sigma(y) = ky$ is known as the warp factor [note that the signature is $(-, +, +, +, +)$] and with $k \sim M_{\text{Pl}}$ as the 5D curvature. We assume that $\sigma(y_0) = 1$ and $\sigma(y_1) \simeq 34$ such that there are some 15 orders of magnitude of scale hierarchy between both boundaries.

In the absence of fermion brane kinetic terms (proportional to α_i 's), this setup produces a tower of Kaluza-Klein (KK) modes, such that the lowest lying modes of the doublet and singlet fields have wave functions exponentially localized towards either of the boundaries [26]. The localization depends on the value of the 5D fermion mass parameters c_q and c_t . When $c_q < 1/2$ and $c_t > -1/2$, the zero modes of Q and T will be localized near the $y = y_1$ boundary, and will be identified as the two chiral components of the SM top quark. The rest of SM quarks will be obtained in a similar way, but the value of their bulk mass will localize them towards the $y = y_0$ boundary. Because the Higgs boson is by construction located at the $y = y_1$ boundary, the top quark will be ‘‘naturally’’ heavy (coupled strongly to the Higgs) whereas the rest of the quarks are lighter, since they couple to the Higgs weakly due to their geographical separation. On the other hand, the excited modes of all fermions will be very heavy and localized towards the $y = y_1$ boundary; their typical KK masses are of order $M_{\text{Pl}} e^{-ky_1} \sim \text{TeV}$ and they will also couple strongly with the Higgs.

The scenario that we call the brane Higgs limit requires the presence of *only* the top quark and bottom quark heavy partners, which means the rest of the KK partners should be decoupled (i.e. much heavier). For this we turn on the fermion brane kinetic terms (the α_i 's) of the 5D top and bottom quarks. There will still be massless fermion modes (associated to the SM top and bottom quarks⁴), but it now becomes possible to lower the masses of the KK top and bottom modes. In general, it is possible to obtain analytically the associated KK spectrum (before electroweak symmetry breaking) in terms of Bessel functions. Nevertheless, since we are mainly interested in the top quark, it is much simpler and transparent to treat the special case where $c_q = c_t = 0$. These simple bulk masses are perfectly toplike, and they have the advantage of producing very simple equations of motion. The usual dimensional reduction procedure involves a mixed separation of variables performed on the 5D fermions, i.e.

⁴Which acquire their SM masses after electroweak symmetry breaking, like in the SM.

$$Q_L(x, y) = Q_L(y)t_L(x) \quad (46)$$

$$Q_R(x, y) = Q_R(y)t_R(x) \quad (47)$$

$$T_L(x, y) = T_L(y)t_L(x) \quad (48)$$

$$T_R(x, y) = T_R(y)t_R(x) \quad (49)$$

where $t_L(x)$ and $t_R(x)$ are the left- and right-handed components of 4D fermions (the lightest of which is the SM top quark). From there, one must solve for the KK profiles along the extra dimension $Q_L(y)$, $Q_R(y)$, $T_L(y)$ and $T_R(y)$. In the simple case of $c_q = c_t = 0$, and before electroweak symmetry breaking, the equation for the profile $\tilde{Q}_R(y) = e^{-2\sigma(y)}Q_R(y)$, for example, becomes

$$(e^{-ky}\tilde{Q}'_R)' + m^2 e^{ky}\tilde{Q}_R = 0 \quad (50)$$

with the Dirichlet boundary condition on the $y = 0$ boundary (since there are no kinetic terms there). The solution is simple,

$$\tilde{Q}_R(y) = N_Q \sin\left(\frac{m(e^{ky} - 1)}{k}\right), \quad (51)$$

which obviously vanishes at $y = 0$. The brane kinetic term on the $y = y_1$ boundary enforces a matching boundary condition at that location, and that boundary condition fixes the spectrum of the whole tower of KK modes. In this simple case, the KK spectrum of the 5D fermions $Q(x, y)$ and $T(x, y)$ is given by

$$\tan\left(\frac{me^{ky_1}}{k}\right) = -\sqrt{\frac{\alpha_{qL}}{\alpha_{qR}}} \tan\left(me^{ky_1}\sqrt{\alpha_{qR}\alpha_{qL}}\right) \quad (52)$$

and

$$\tan\left(\frac{me^{ky_1}}{k}\right) = -\sqrt{\frac{\alpha_{tL}}{\alpha_{tR}}} \tan\left(me^{ky_1}\sqrt{\alpha_{tR}\alpha_{tL}}\right), \quad (53)$$

in agreement with the flat metric limit considered in [24]. With a further simplification, taking $\alpha_{qR} = \alpha_{qL} = \alpha_q$ and $\alpha_{tR} = \alpha_{tL} = \alpha_t$, the conditions become

$$\tan\left(\frac{me^{ky_1}}{k}\right) = -\tan(me^{ky_1}\alpha_i) \quad (54)$$

with a spectrum given by

$$m_n = \frac{n\pi}{1 + k\alpha_i} ke^{-ky_1} \quad (55)$$

for $n = 0, 1, 2, 3, \dots$. This shows that, indeed, the spectrum of the KK tops can be significantly reduced in the presence of brane kinetic terms.

In the scenario we have in mind, only the 5D top and bottom quarks have large brane kinetic terms (without further justification) and therefore their associated KK modes can be much lighter than the rest, maybe as light as 1 TeV. At the same time, the rest of the quarks and KK gauge bosons follow the usual RS pattern with KK masses maybe an order of magnitude larger (~ 10 TeV). In this limit, flavor and precision electroweak bounds are much safer and the main phenomenological effects of the model may occur within the Higgs sector of the scenario.

If we decouple the up, down, strange and charm heavy KK quarks, the fermion mass matrices will involve only SM quarks along with KK tops and KK bottoms. Mixing between light quarks localized near $y = 0$ and heavier quarks localized near $y = y_1$ is going to be CKM suppressed, as usual in RS, and therefore the mass matrices to consider have the same form of those in Eqs. (2) and (3), but with the phases $\theta_1^t = \theta_1^b = 0$. Indeed, the values of the off-diagonal terms in the mass matrices are now associated to the 5D Yukawa interactions localized at the $y = y_1$ brane and in this case are such that

$$Y_t^0 = Y_{33}^{5D} f_{q_L} f_{t_R} \quad (56)$$

$$(Y_1^t)_{nm} = Y_{33}^{5D} f_{Q_L^n} f_{T_R^m} \quad (57)$$

$$(Y_{qT})_n = Y_{33}^{5D} f_{q_L} f_{T_R^n} \quad (58)$$

$$(Y_{Qt})_n = Y_{33}^{5D} f_{Q_L^n} f_{t_R} \quad (59)$$

where Y_{33}^{5D} is the 5D top Yukawa coupling, and where the f_i 's are the wave functions evaluated at the $y = y_1$ brane,⁵ with q_L and t_R being zero modes and Q_L^n and T_R^n representing the n th KK modes. We can see that all these terms share the same phase, so that we can set $\theta_1^t = \theta_1^b = 0$.

Now, consider first a mass matrix with only one KK level ($n = 1$), so that the matrices are exactly the same as before and thus the corresponding effect in Higgs production will come from the sum

$$\sum_n \left(\frac{y_{nn}^u}{m_n^u} \right) = \frac{1}{v} \frac{1 + 3\varepsilon_{Q_t} \varepsilon_T \frac{|Y_1^t|}{|Y_1^b|} e^{i\theta_1^t} \left(1 - \frac{|Y_1^t| |Y_1^0|}{|Y_{Q_t}| |Y_{qT}|}\right)}{1 + \varepsilon_{Q_t} \varepsilon_T \frac{|Y_1^t|}{|Y_1^b|} e^{i\theta_1^t} \left(1 - \frac{|Y_1^t| |Y_1^0|}{|Y_{Q_t}| |Y_{qT}|}\right)}. \quad (60)$$

It becomes then apparent that $(1 - \frac{|Y_1^t| |Y_1^0|}{|Y_{Q_t}| |Y_{qT}|}) = 0$ due to the structure of the 5D couplings. It is important that the zero

⁵Note that another effect of the brane kinetic terms is to suppress the value of wave functions through normalization, due to the new brane localized kinetic terms. Nevertheless, in order to obtain the top quark mass, the 5D coupling Y_{33} must be enhanced accordingly and thus the wave function suppression is compensated by a coupling enhancement, while remaining in a perturbative regime [25].

modes (SM top and bottom) come from the same 5D fermion as the KK modes, since the cancellation will only happen if they all share the same 5D Yukawa coupling. It turns out that it is simple to prove that if we take into account the complete towers of KK tops we still have

$$\sum_n^{\infty} \left(\frac{y_{nn}^u}{m_n^u} \right) = \frac{1}{v} \quad (61)$$

and similarly for the bottom quarks, and thus the Higgs phenomenology of this scenario is indeed the same as in the brane Higgs limit introduced earlier in a bottom-up approach, since relative corrections due to the KK modes of the up, down, strange and charm quarks will scale as $(m_{\text{KK}}^{\text{tops}}/m_{\text{KK}}^{\text{rest}})^2 \sim 1\%$ (assuming that the rest of the KK quarks are an order of magnitude heavier than the KK tops/bottoms). Note that if the Higgs boson is not exactly localized at the boundary, then the cancellation will not be exact and new corrections will arise.

The contributions of this RS scenario to flavor and precision electroweak observables will be limited to effects due to the mixing of top and bottom with the vectorlike partners, since we are considering very heavy KK gauge bosons (~ 10 TeV). As pointed out earlier, Yukawa coupling mixing effects can lead to deviations in $|V_{tb}|$ which can easily be kept under control. Also effects can appear in the couplings $Z \rightarrow b_R \bar{b}_R$ and in $Z \rightarrow b_L \bar{b}_L$ [27], but since we consider the contribution from heavy KK gauge bosons to be suppressed, only Yukawa coupling mixings contribute, limiting the correction. In the usual RS scenario, it was already possible to find points in the (c_b, c_{q_3}) plane, such that $Zb\bar{b}$ couplings remain within experimental bounds, with all the SM masses and mixings correctly obtained [28]. In our scenario, finding parameter points safe from precision tests will be even easier, since the source of corrections is further limited.

IV. CONCLUSION

In this work, we presented a simple explanation of the possible enhancement in the $t\bar{t}h$ associated production seen at the LHC. We added one $SU(2)$ doublet and two $SU(2)$ singlet vectorlike quarks to the matter content of the SM, as partners of the third family, and allow significant mixing between these with the third family only. After electroweak symmetry breaking, Yukawa couplings induce off-diagonal terms into the fermion mass matrix and, once in the physical basis, the top and bottom Yukawa couplings and their corresponding masses lose their SM alignment. With the proper sign (or phase), this misalignment can induce an enhancement of the top quark Yukawa coupling, and thus increase the cross section for $t\bar{t}h$ production. But the mechanism should also affect other observables in the Higgs sector, in particular, the cross section for Higgs production through gluon fusion. This is the main production channel for the SM Higgs and, being a radiative effect, it

receives a contribution from all the fermions in the model. Each contribution is proportional to the ratio of the Yukawa fermion couplings to the fermion mass, so that the main contributions come from the top quark (with an enhanced Yukawa coupling), and from the new vectorlike quarks.

We showed that working in a particular limit of parameter space, the corrections to gluon fusion caused by the top Yukawa coupling enhancement are exactly offset by the contributions of the new top partners, so that the overall top sector of our scenario (top quark plus heavy partners) gives the same contribution as the single top quark contribution in the SM. We call this scenario the brane Higgs limit and it yields extremely predictive relationships between the production cross sections (gg and $t\bar{t}h$) and decay branching ratios for the Higgs bosons (into WW , ZZ , $b\bar{b}$ and $\gamma\gamma$), where the only free parameters are the absolute values of the shifts in the Yukawa couplings of the top and the bottom quarks, and their phases. For instance, in $t\bar{t}h$ production, if the branching ratio into $b\bar{b}$ is smaller than that into VV , then the gluon fusion production cross section must be also greater than its SM value, and conversely, an enhancement of the $b\bar{b}$ branching ratio in $t\bar{t}h$ production indicates a suppressed gluon fusion signal. Overall, the deviations in the Yukawa coupling of the bottom quark are constrained to be small, unless new data indicate a significant enhancement of the $b\bar{b}$ branching ratio. The scenario we consider predicts that any enhancement or suppression in the $\gamma\gamma$ signal should be matched with identical enhancement or suppression in $VV = WW, ZZ$ decays (for gluon fusion production), or at least remain always slightly higher than decays into $VV = WW, ZZ$, but never lower. Finally, a shift in the top quark Yukawa coupling will affect all $t\bar{t}h$ signals through the production cross section, and enhancement or suppression will depend on the phase of the shift.

The mixing in the top quark and bottom quark sectors should also have consequences in the CKM mixing matrix V_{CKM} , as well as in precise electroweak measurements. We briefly discussed how the scenario affects (and thus is constrained by) the V_{tb} entry of V_{CKM} and the decay $Z \rightarrow b\bar{b}(A_{FB}^b)$.

Finally we showed that the phenomenology we described here depends on a specific structure of the fermion mixing matrix, mixing top quark with its partners. In particular the Yukawa coupling matrix should have a vanishing determinant, and thus some mechanism or flavor symmetry should be invoked to realize the scenario. The required structure is naturally realized in a Randall-Sundrum model without a need for flavor symmetries. A key ingredient of this scenario is the presence of brane kinetic terms for the top and bottom, which can then result in lighter $n = 1$ KK modes for the top and bottom partners, but heavy masses for all other KK modes. If the Higgs is localized exactly at the boundary, the overall phenomenology of the simple model introduced here is essentially recovered (i.e. the cancellation of the terms happening in

the gluon fusion calculation occurs by construction, even if in this case the effect comes from a complete tower of KK states).

The model presented here thus has a simple theoretical realization, is highly predictable, and can be tested (or ruled out) by more precise measurements of the Higgs signal strength in Run 2 at LHC.

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