Higgs diphoton rate enhancement from supersymmetric physics beyond the MSSM

Marcus Berg,[1,*] Igor Buchberger,[1,†] D. M. Ghilencea,[2,3,‡] and Christoffer Petersson[4,5,§]

[1]*Department of Physics, Karlstad University, 651 88 Karlstad, Sweden*
[2]*Theory Division, CERN, 1211 Geneva 23, Switzerland*
[3]*Theoretical Physics Department, National Institute of Physics and Nuclear Engineering (IFIN-HH), Bucharest MG-6, 077125 Romania*
[4]*Physique Théorique et Mathématique, Université Libre de Bruxelles, C.P. 231, 1050 Bruxelles, Belgium*
[5]*International Solvay Institutes, 1050 Brussels, Belgium*

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We show that supersymmetric "new physics" beyond the minimal supersymmetric standard model can naturally accommodate a Higgs mass near 126 GeV and enhance the signal rate in the $h \to \gamma\gamma$ channel, while the signal rates in all the other Higgs decay channels coincide with Standard Model expectations, except possibly the $h \to Z\gamma$ channel. The new physics that corrects the relevant Higgs couplings can be captured by two supersymmetric effective operators. We provide a simple example of an underlying model in which these operators are simultaneously generated. The scale of new physics that generates these operators can be around 5 TeV or larger, and outside the reach of the LHC.

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I. INTRODUCTION

The ATLAS and CMS collaborations at CERN recently presented strong experimental evidence for a Higgs-like resonance around 126 GeV [1], marking a historic achievement in particle physics. The signal rates in the ZZ^* and WW^* channels are in good agreement with the Standard Model (SM) predictions. The $b\bar{b}$ and $\tau^+\tau^-$ signal rates are also compatible with the SM, although with substantial error bars. At the time of writing, the signal rate in the $h \to \gamma\gamma$ channel is about 1.5–2 times larger than the SM prediction. This discrepancy is only at the level of two standard deviations, and there are theoretical uncertainties [2]. Nevertheless, it is interesting to contemplate whether physics beyond the SM can be responsible for this discrepancy (for recent work in this direction, see Ref. [3]).

In this work we shall assume that the excess in the diphoton channel is due to beyond the SM physics that has a negligible effect on channels other than $h \to \gamma\gamma$. Indeed, this channel is sensitive to new physics, since it is a loop-level process in the SM. With this in mind, we shall focus on the minimal supersymmetric (SUSY) extension of the SM (MSSM) close to the "decoupling limit" [4], in which the lightest neutral *CP*-even Higgs boson h is SM-like.

It is surprising that, despite its large number of parameters, the MSSM has difficulties in accommodating an enhancement of the $h \to \gamma\gamma$ partial decay width $\Gamma_{h\gamma\gamma}$ without affecting the other partial decay widths. In fact this requirement seems to single out loop-induced contributions from very light color singlet superpartners with a significant coupling to the Higgs, meaning strongly mixed light stau sleptons, at around 100 GeV [5]. However, this introduces issues with vacuum stability and may even be possible to rule out at the LHC. In addition, large radiative corrections are needed to obtain a mass of $m_h \approx 126$ GeV for the lightest Higgs of the MSSM. This requires large supersymmetry-breaking terms, such as TeV stop masses and/or a large top A term. The lack of evidence for superpartners in the direct SUSY searchers at the LHC also indicates that soft terms should be large. However, large supersymmetry-breaking terms lead to severe fine-tuning [6,7] in most versions of the MSSM[1] (with or without universal gaugino masses or Higgs soft masses at the UV scale different from the other scalar soft masses). This situation suggests that a solution to the problems of the Higgs mass and the diphoton rate is not in the SUSY-breaking sector but rather in the one that preserves supersymmetry.

Following this idea, the purpose of this work is to answer whether one can have a minimal *supersymmetric* extension of the Higgs sector that allows for $m_h \approx 126$ GeV *without* undue fine-tuning and a *simultaneous* $h \to \gamma\gamma$ enhancement, no changing of the other Higgs branching ratios, while also complying with the negative SUSY searches so far. To this end, one way to proceed is suggested by minimal extensions of the MSSM Higgs sector, like the next-to-MSSM (NMSSM) model which contains an additional singlet chiral superfield (see Ref. [9] for a review), and by models where the soft $B\mu$ term is promoted to a SUSY operator [10]. In the NMSSM an enhancement of the $h \to \gamma\gamma$ branching ratio is possible, but unfortunately this also alters other couplings beyond the SM level [11]. Further, the NMSSM also remains badly fine-tuned

*marcus.berg@kau.se
†igorbuch@kau.se
‡dumitru.ghilencea@cern.ch
§christoffer.petersson@ulb.ac.be

[1]For a discussion of the negative impact of the electroweak fine tuning on the χ^2 fit of such models, see also Ref. [8].

(fine-tuning $\Delta > 200$) for $m_h \approx 126$ GeV [12]. There are known ways to bypass this problem such as in the so-called "generalized" version of NMSSM with a superpotential mass term for the singlet superfield, where the electroweak fine-tuning is significantly reduced to more acceptable levels ($\Delta \approx 30$) [12–14]. These are examples of explicit supersymmetry-preserving modifications of the MSSM that can render more natural the interpretation of the resonance at 126 GeV as the lightest Higgs.

In this work, instead of considering such specific models, we shall relax the rigid, minimal structure of the MSSM Higgs sector and perform an effective field theory analysis of the most relevant SUSY-preserving operators in this sector. This approach should recover, in a particular region of the parameter space, scenarios such as those presented above (generalized NMSSM, etc). We show that it is possible to *naturally* accommodate a Higgs mass of 126 GeV, an enhanced Higgs coupling to photons, and *simultaneously* SM-like Higgs couplings to the other particles, using only a few supersymmetry-preserving operators with small coefficients.

In general there is a large set of operators that one could consider in the Higgs sector [15]. Regarding the Higgs mass, it is known that the presence of the effective dimension-five superpotential operator [16,17]

$$\frac{1}{M}(H_u \cdot H_d)^2 \tag{1.1}$$

can accommodate $m_h \approx 126$ GeV without undue fine-tuning [13]. The suppression scale M represents the mass scale of the SUSY degrees of freedom that have been integrated out to generate Eq. (1.1). Concerning $\Gamma_{h\gamma\gamma}$, from the list of effective operators of dimensions $d = 5$ and $d = 6$ in the Higgs sector [13,18], one notices the presence of a SUSY effective operator

$$\frac{1}{M^2}(H_u \cdot H_d)\mathrm{Tr}(W^\alpha W_\alpha), \tag{1.2}$$

where W_α is the electroweak gauge field strength superfield. The operator (1.2) can significantly increase the $h \to \gamma\gamma$ partial width. Note that this increase, and eventually also the change in the $h \to Z\gamma$ partial width, can be accomplished without affecting the other partial widths.[2] Based on these observations, we intend to investigate closer the phenomenological impact of Eqs. (1.1) and (1.2).

As a simple example will show, these operators can be generated simultaneously by an underlying microscopic model, making their combination rather natural. The effect of both operators is maximized for small $\tan\beta$, where the MSSM tree-level contribution to the Higgs mass is minimized. This means that the impact of the $d = 5$ operator in accommodating an $m_h \approx 126$ GeV is rather significant. To our knowledge, the particular combination of the effective SUSY operators in Eqs. (1.1) and (1.2) has not been studied in the past for the problems we address.[3] One can also use this information to set bounds on the scale M of new physics.

The paper is organized as follows. In Sec. II we calculate the corrections from the effective operators to the Higgs mass and mixing angle. In Sec. III we discuss how these operators correct the Higgs couplings and signal rates, with a focus on the decoupling limit. The results in terms of the Higgs mass and the partial widths for the $h \to \gamma\gamma$ and $h \to Z\gamma$ channels are discussed in Sec. IV. Section V provides an example of the origin of the effective operators, and Sec. VI contains our conclusions, while some details concerning the on-shell Lagrangian are given in the Appendix.

II. CORRECTIONS FROM SUSY OPERATORS TO THE HIGGS COUPLINGS

The effective model we consider consists of the usual MSSM Higgs sector, extended by the operators discussed in the introduction. The relevant part of the Lagrangian is, in standard notation,

$$\mathcal{L} = \int d^4\theta \sum_{i=u,d}(1 - m_i^2\theta^2\bar{\theta}^2)H_i^\dagger e^{V_i}H_i + \left(\int d^2\theta\mu(1 + B\theta^2)H_d \cdot H_u + \mathrm{H.c.}\right) + \mathcal{O}_5 + \mathcal{O}_6, \tag{2.1}$$

where the chiral superfields have components $H_i \equiv (h_i, \psi_i, F_i)$, and m_i and B are the soft terms. $\mathcal{O}_5$ is the only operator of dimension five that one can write in the Higgs sector, up to nonlinear field redefinitions [20], and has the form

$$\mathcal{O}_5 = \frac{c_0}{M}\int d^2\theta (H_u \cdot H_d)^2 + \mathrm{H.c.} \tag{2.2}$$

For the component fields expression of $\mathcal{O}_5$, see Eq. (A3).

There is a long list of operators in the Higgs sector of dimension $d = 6$ [13,18,20]. A careful analysis of these operators shows that of all these, there is one of them that can couple, in a supersymmetric way, to two gauge bosons:

$$\mathcal{O}_6 = \frac{1}{M^2}\sum_{s=1,2}\frac{c_s}{16g_s^2\kappa_s}\int d^2\theta \mathrm{Tr}(W^\alpha W_\alpha)_s(H_u \cdot H_d) + \mathrm{H.c.} \tag{2.3}$$

Here g_1 and g_2 denote the $\mathrm{U}(1)_Y$ and $\mathrm{SU}(2)_L$ gauge couplings, respectively, and κ_s is a constant that cancels the trace factor. W^α is the SUSY field strength of the $\mathrm{U}(1)_Y$ ($\mathrm{SU}(2)_L$) vector superfield V_1 (V_2) of components (λ_s, $V_{s,\mu}$, $D_s/2$), $s = 1, 2$. $\mathcal{O}_5$ and $\mathcal{O}_6$ provide a minimal set of operators that is enough for our purposes. One can

[2] When W_α corresponds to the gauge field strength of $\mathrm{SU}(3)_C$ gauge group, the operator (1.2) can affect the $h \to gg$ partial width (not considered in this work).

[3] The operator in Eq. (1.2) was separately discussed in Ref. [19].

also consider SUSY-breaking effects associated to these operators [see the Appendix, Eq. (A3)], but we only seek supersymmetric solutions to our problem. The effective expansion is reliable when $c_{0,1,2} = \mathcal{O}(1)$ and M is the largest scale in the theory. One can choose one of $c_{0,1,2}$, for example c_0, and set it to $c_0 = 1$ by redefining M. But it is useful to keep c_0 to easily trace or turn off the effects of $\mathcal{O}_5$. Also, to modify the diphoton rate c_1 or c_2 (or a combination thereof) is enough, so together with the scale M, we effectively have only two parameters. Keeping both $c_{1,2}$ generates an additional interesting coupling; see later in the text.

Additional operators of $d = 6$ can be present. Although they could have an impact on the Higgs mass [18], they have an additional scale suppression relative to[4] $\mathcal{O}_5$. There is an operator similar to $\mathcal{O}_6$ but involving instead the $SU(3)_C$ gauge group, which we do not consider here; this would change dramatically the Higgs decay rate to gluons, away from the SM values. We take the good agreement with the SM in most channels as evidence that if present, the coefficient of this operator must be small. Finally, another reason to restrict our analysis to $\mathcal{O}_{5,6}$ is that, as discussed later, they can be simultaneously generated by underlying physics.

A. On-shell Lagrangian

The calculation of the on-shell Higgs Lagrangian extended by $\mathcal{O}_5$ and $\mathcal{O}_6$ is detailed in the Appendix. The result is

$$\begin{aligned}
\mathcal{L} = &-\frac{1}{2}\Big[D_2^a D_2^a\Big(1+\frac{c_2}{2M^2}(h_u\cdot h_d+\text{H.c.})\Big)+(2\to 1)\Big]\\
&-\Big|\mu+2\frac{c_0}{M}h_d\cdot h_u\Big|^2(|h_d|^2+|h_u|^2)+\Big[\frac{\mu}{4}\Big(\frac{c_2}{M^2}\lambda_2^a\lambda_2^a+\frac{c_1}{M^2}\lambda_1^2\Big)(|h_d|^2+|h_u|^2)+\text{H.c.}\Big]\\
&+\Big\{\frac{c_2}{4M^2}(h_u\cdot h_d)[i(\lambda_2^a\sigma^\mu\mathcal{D}_\mu\bar\lambda_2^a-\mathcal{D}_\mu\bar\lambda_2^a\bar\sigma^\mu\lambda_2^a)]+\text{H.c.}+(2\to 1)\Big\}\\
&+\frac{c_0}{M}[2(h_u\cdot h_d)(\psi_d\cdot\psi_u)-(h_u\cdot\psi_d+\psi_u\cdot h_d)^2]+\text{H.c.}\\
&+\Big\{\frac{c_2}{4M^2}\Big[-\frac{1}{2}(h_u\cdot h_d)\Big(F_2^{a\mu\nu}F_{2\mu\nu}^a+\frac{i}{2}\epsilon^{\mu\nu\rho\sigma}F_{2\mu\nu}^aF_{2\rho\sigma}^a\Big)\\
&-\sqrt{2}(h_u\cdot\psi_d+\psi_u\cdot h_d)\sigma^{\mu\nu}\lambda_2^aF_{2\mu\nu}^a-\psi_u\cdot\psi_d\lambda_2^a\lambda_2^a\Big]+(2\to 1)+\text{H.c.}\Big\}\\
&+[\mu B(h_d\cdot h_u)+\text{H.c.}]-\tilde m_d^2|h_d|^2-\tilde m_u^2|h_u|^2,
\end{aligned}\tag{2.4}$$

where $\tilde m_i^2 = m_i^2+|\mu|^2$, $i=u,d$. For the explicit form of $D_2^aD_2^a$ and D_1^2, see Eqs. (A7) and (A8).

Equation (2.4) contains all the information one needs to extract the corrections to the Higgs masses and couplings. In particular, notice the presence of new, supersymmetric couplings:

$$-\frac{1}{8}(h_u\cdot h_d)\Big(\frac{c_2}{M^2}\,\text{Tr}F_2^2+\frac{c_1}{M^2}\,\text{Tr}F_1^2\Big)-\Big|\mu+2\frac{c_0}{M}h_d\cdot h_u\Big|^2(|h_d|^2+|h_u|^2)+\text{H.c.},\tag{2.5}$$

which are important below. There are also direct Higgs-Higgsino and Higgsino-gaugino couplings that can be relevant for dark matter models. From Eq. (2.4) we find the Higgs scalar potential V_h:

$$\begin{aligned}
V_h = &\ \tilde m_d^2|h_d|^2+\tilde m_u^2|h_u|^2-[\mu B h_d\cdot h_u+\text{H.c.}]+\frac{g_2^2}{2}|h_d^\dagger h_u|^2\Big[1+\frac{c_2}{2M^2}(h_d\cdot h_u+\text{H.c.})\Big]\\
&+\frac{1}{8}(|h_d|^2-|h_u|^2)^2\Big[g^2+\Big[(h_d\cdot h_u)\Big(\frac{g_1^2c_1}{M^2}+\frac{g_2^2c_2}{M^2}\Big)+\text{H.c.}\Big]\Big]+4\Big|\frac{c_0}{M}\Big|^2|h_d\cdot h_u|^2(|h_d|^2+|h_u|^2)\\
&+\Big[\Big(2\frac{c_0}{M}\mu^*\Big)(|h_d|^2+|h_u|^2)(h_d\cdot h_u)+\text{H.c.}\Big],\qquad (g^2\equiv g_1^2+g_2^2),
\end{aligned}\tag{2.6}$$

which depends on two parameters: c_0 from the effective dimension-five operator and the combination $(g_1^2c_1 + g_2^2c_2)$ from the effective dimension-six operator. Note that last term in the first line above does not contribute to the neutral Higgs sector masses.

We also include dominant loop corrections, although they do not play the same crucial role they do in the MSSM. In the small $\tan\beta$ regime and for dominant top Yukawa coupling, the one-loop and leading two-loop correction to V_h is [21]

[4]Strictly speaking, this is true for the small $\tan\beta$ region, which will actually be the relevant region in our case.

$$\Delta V_h = \frac{g^2}{8}\delta|h_u|^4, \tag{2.7}$$

where

$$\delta \equiv \frac{3h_t^4}{g^2\pi^2}\left[\ln\frac{M_{\tilde{t}}}{m_t} + \frac{X_t}{4} + \frac{1}{32\pi^2}(3h_t^2 - 16g_3^2) \times \left(X_t + 2\ln\frac{M_{\tilde{t}}}{m_t}\right)\ln\frac{M_{\tilde{t}}}{m_t}\right]$$

$$X_t \equiv \frac{2(A_t - \mu\cot\beta)^2}{M_{\tilde{t}}^2}\left(1 - \frac{(A_t - \mu\cot\beta)^2}{12M_{\tilde{t}}^2}\right), \tag{2.8}$$

with $M_{\tilde{t}}^2 \equiv m_{\tilde{t}_1}m_{\tilde{t}_2}$, and g_3 is the QCD coupling.

B. Higgs mass and mixing angle

The scalars receive mass corrections from the usual one-loop radiative corrections but now also from the effective operators. Here we take the parameters c_0, c_1, c_2 to be real. We find the following result for the mass of the lightest Higgs scalar h:

$$m_h^2 = \frac{1}{2}\{m_A^2 + m_Z^2 + \delta m_Z^2\sin^2\beta - \sqrt{w}\} + \Delta m_h^2, \tag{2.9}$$

where

$$w \equiv [(m_A^2 - m_Z^2)\cos 2\beta + \delta m_Z^2 \sin^2\beta]^2 + \sin^2 2\beta(m_A^2 + m_Z^2)^2 \tag{2.10}$$

and where Δm_h^2 is the contribution due to the higher-dimensional operators:

$$\Delta m_h^2 = \left(2\mu\frac{c_0}{M}\right)s_1 + \left(2\mu\frac{c_0}{M}\right)^2 s_2 + \left(\frac{g_1^2c_1}{M^2} + \frac{g_2^2c_2}{M^2}\right)s_3 + \mathcal{O}\left(\frac{1}{M^3}\right), \tag{2.11}$$

with

$$s_1 = v^2\sin 2\beta\left\{1 + \frac{(m_A^2 + m_Z^2)}{\sqrt{w}}\right\}$$

$$s_2 = \frac{v^4}{4\mu^2}\sin^2 2\beta + \frac{v^4}{\sqrt{w}}\left\{-1 + \frac{1}{2\mu^2}(m_A^2 + m_Z^2)\sin^2 2\beta\right\} + \frac{1}{w^{3/2}}(m_A^2 + m_Z^2)^2 v^4 \sin^2 2\beta$$

$$s_3 = \frac{v^4}{32}\sin 2\beta + \frac{v^4\sin 2\beta}{128\sqrt{w}}[8m_A^2 - (4 + 3\delta)m_Z^2 + 6\delta m_Z^2\cos 2\beta + 3(4m_A^2 - \delta m_Z^2)\cos 4\beta], \tag{2.12}$$

where we kept (small) effects from the interplay between the effective operators and the one-loop correction to V_h. The mass of the CP-odd Higgs boson is

$$m_A^2 = \frac{2B\mu}{\sin 2\beta} - \frac{2v^2}{\sin 2\beta}\left(\frac{c_0}{M}\mu\right) - \frac{v^4}{32}\frac{\cos^2 2\beta}{\sin 2\beta}\left(\frac{g_1^2c_1}{M^2} + \frac{g_2^2c_2}{M^2}\right) + \mathcal{O}\left(\frac{1}{M^3}\right), \tag{2.13}$$

which does not receive one-loop corrections. The mixing angle α is given by

$$\tan 2\alpha = -\frac{1}{\mathcal{D}}\left[(m_A^2 + m_Z^2)\tan 2\beta - \frac{2v^2}{\cos 2\beta}\left(2\mu\frac{c_0}{M}\right) - \left(2\frac{c_0}{M}\right)^2 v^4\tan 2\beta + \left(\frac{g_1^2c_1}{M^2} + \frac{g_2^2c_2}{M^2}\right)\frac{v^4(m_Z^2 f_Z - m_A^2 f_A)}{32\mathcal{D}\cos^2 2\beta}\right], \tag{2.14}$$

with

$$\mathcal{D} = m_A^2 - m_Z^2 + (\sec 2\beta - 1)\delta m_Z^2/2$$

$$f_Z = 4\cos 2\beta - (2 - 5\cos 2\beta + 6\cos 4\beta - 3\cos 6\beta)\delta/4$$

$$f_A = \cos 2\beta + 3\cos 6\beta. \tag{2.15}$$

One can see the corrections to $\tan 2\alpha$ due to the effective operators, which are used below.

We also note that the new operators correct the gauge field kinetic terms when the Higgs fields receive vacuum expectation values (VEVs). The corrected gauge couplings are the ones that appear in the following.

III. CORRECTIONS TO THE PARTIAL WIDTHS OF $h \to \gamma\gamma$ AND $h \to Z\gamma$

In this section we study how the new operators correct the Higgs couplings to the SM particles. To this end, we parametrize these corrections in terms of the usual MSSM Higgs couplings.

A. Higgs couplings and signal rates

The renormalizable part of the Lagrangian for the lightest neutral CP-even Higgs scalar h can be written [22] as

$$\mathcal{L}_{\rm ren} = -c_t\frac{m_t}{v}ht\bar{t} - c_c\frac{m_c}{v}hc\bar{c} - c_b\frac{m_b}{v}hb\bar{b} - c_\tau\frac{m_\tau}{v}h\tau^+\tau^- + c_Z\frac{m_Z^2}{v}hZ^\mu Z_\mu + c_W\frac{2m_W^2}{v}hW^{+\mu}W^-_\mu, \tag{3.1}$$

where the dimensionless coefficients are given by

$$c_t = c_c = \frac{\cos\alpha}{\sin\beta}, \qquad c_b = c_\tau = -\frac{\sin\alpha}{\cos\beta}, \qquad c_Z = c_W = \sin(\beta - \alpha), \tag{3.2}$$

where the mixing angle α is given in Eq. (2.14). In the scenario under consideration, all loop corrections to the tree-level coefficients in Eq. (3.2) are negligible. The usual

SM values for the couplings in Eqs. (3.1) and (3.2) are obtained in the decoupling limit, in which $\alpha \to \beta - \pi/2$, implying that $\cos\alpha \to \sin\beta$, $\sin\alpha \to -\cos\beta$, and hence $c_i \to c_i^{\rm SM} = 1$, where $i = t, c, b, \tau, Z, W$.

We work in the limit when loop contributions from superpartners and other Higgs scalars are negligible. The dimension-five part of the Higgs Lagrangian, which takes into account one-loop contributions from SM particles as well as the contributions from the effective operators in Eq. (2.4) can be written as

$$\mathcal{L}_{\rm dim\,5} = c_g^{\rm loop}\frac{\alpha_{\rm S}}{12\pi v}h{\rm Tr}G^{\mu\nu}G_{\mu\nu} + (c_\gamma^{\rm loop} + c_\gamma^{\rm BMSSM}) \times \frac{\alpha_{\rm EM}}{8\pi v}hF^{\mu\nu}F_{\mu\nu} + (c_{\gamma Z}^{\rm loop} + c_{\gamma Z}^{\rm BMSSM}) \times \frac{\alpha_{\rm EM}}{4\pi\sin\theta_w v}hZ^{\mu\nu}F_{\mu\nu}. \tag{3.3}$$

The one-loop contributions to these coefficients are given by[5] [24]

$$\begin{aligned} c_g^{\rm loop} &= c_t\mathcal{A}_g^{(t)} + c_b\mathcal{A}_g^{(b)} \approx 1.03c_t - (0.05 + 0.07i)c_b \\ c_\gamma^{\rm loop} &= c_W\mathcal{A}_\gamma^{(W)} + c_t\mathcal{A}_\gamma^{(t)} \approx -8.36c_W + 1.84c_t \\ c_{Z\gamma}^{\rm loop} &= c_W\mathcal{A}_{Z\gamma}^{(W)} + c_t\mathcal{A}_{Z\gamma}^{(t)} \approx 5.80c_W - 0.31c_t, \end{aligned} \tag{3.4}$$

where, in the last steps, we have inserted $m_h = 126$ GeV in the one-loop form factors $\mathcal{A}$, for which the explicit expressions are given in Appendix A 2. In the decoupling limit, where $c_i \to c_i^{\rm SM} = 1$ in Eq. (3.2), the $c^{\rm loop}$ coefficients in Eq. (3.4) approach the values they have in the SM, which, for $m_h = 126$ GeV, follow trivially from Eq. (3.4),

$$c_g^{\rm loop} \to c_g^{\rm SM} \approx 0.98 + 0.07i \qquad c_\gamma^{\rm loop} \to c_\gamma^{\rm SM} \approx -6.52 \qquad c_{\gamma Z}^{\rm loop} \to c_{\gamma Z}^{\rm SM} \approx 5.49. \tag{3.5}$$

In order to obtain the $c_\gamma^{\rm BMSSM}$ and $c_{\gamma Z}^{\rm BMSSM}$ coefficients in Eq. (3.3), we extract the relevant component interactions from the operators in Eq. (2.3), written in Eqs. (2.4) and (2.5),

$$\mathcal{O}_6 \supset \frac{v\cos(\beta+\alpha)}{8M^2}([c_1\cos^2\theta_w + c_2\sin^2\theta_w]hF^{\mu\nu}F_{\mu\nu} + 2(c_2 - c_1)\sin\theta_w\cos\theta_w hF^{\mu\nu}Z_{\mu\nu}),$$

where we have used

$$\begin{aligned} &h_u\cdot h_d = h_u^+h_d^- - h_u^0h_d^0, \quad h_i^0 = \frac{1}{\sqrt{2}}(v_i + {\rm Re}h_i^0 + i{\rm Im}h_i^0), \\ &{\rm Re}h_d^0 = -\sin\alpha h + \cos\alpha H, \quad {\rm Re}h_u^0 = \cos\alpha h + \sin\alpha H, \\ &A_{1\mu} = \cos\theta_w A_\mu - \sin\theta_w Z_\mu, \quad A_{2\mu}^{(3)} = \sin\theta_w A_\mu + \cos\theta_w Z_\mu \end{aligned} \tag{3.6}$$

and $v_d = v\cos\beta$, $v_u = v\sin\beta$, with $v = 246$ GeV. Moreover, the hypercharge gauge boson $A_{1\mu}$ and the (third component of the) $SU(2)_L$ gauge boson $A_{2\mu}^{(3)}$ have been rewritten in terms of the photon A_μ and the Z boson Z_μ. Note that there is also a dimension-five operator generated from Eq. (2.3) that involves the Higgs scalar h and two field strengths of the Z boson (as well as an analogous operator involving two field strengths of the W boson). However, since these operators will have couplings comparable to the $\gamma\gamma$ or $Z\gamma$ couplings, but strongly phase space suppressed, we expect them to be irrelevant with respect to the usual dimension-three Higgs coupling to the Z and W bosons in Eq. (3.1). Therefore, we do not consider them.

The contributions to Eq. (3.3) from Eq. (3.6) are given by

$$\begin{aligned} c_\gamma^{\rm BMSSM} &= \frac{\pi v^2\cos(\beta+\alpha)}{M^2\alpha_{\rm EM}}(c_1\cos^2\theta_w + c_2\sin^2\theta_w) \\ c_{\gamma Z}^{\rm BMSSM} &= \frac{\pi v^2\cos(\beta+\alpha)}{M^2\alpha_{\rm EM}}(c_2 - c_1)\sin^2\theta_w\cos\theta_w. \end{aligned} \tag{3.7}$$

In the decoupling limit, where $\cos(\beta+\alpha) \to \sin 2\beta$, we see that the coefficients in Eq. (3.7) are maximized for small $\tan\beta$.

We can now define the relevant Higgs partial decay widths, normalized to the corresponding SM value, in terms of the dimensionless c coefficients in Eqs. (3.2), (3.4), (3.5), and (3.7),

$$\begin{aligned} &\frac{\Gamma_{hii}}{\Gamma_{hii}^{\rm SM}} = |c_i|^2, \qquad \frac{\Gamma_{hgg}}{\Gamma_{hgg}^{\rm SM}} = \left|\frac{c_g^{\rm loop}}{c_g^{\rm SM}}\right|^2, \\ &\frac{\Gamma_{h\gamma\gamma}}{\Gamma_{h\gamma\gamma}^{\rm SM}} = \left|\frac{c_\gamma^{\rm loop} + c_\gamma^{\rm BMSSM}}{c_\gamma^{\rm SM}}\right|^2, \\ &\frac{\Gamma_{h\gamma Z}}{\Gamma_{h\gamma Z}^{\rm SM}} = \left|\frac{c_{\gamma Z}^{\rm loop} + c_{\gamma Z}^{\rm BMSSM}}{c_{\gamma Z}^{\rm SM}}\right|^2 \end{aligned} \tag{3.8}$$

as well as the corresponding branching ratios (BRs),

$$\begin{aligned} &\frac{{\rm BR}_{hii}}{{\rm BR}_{hii}^{\rm SM}} = \left|\frac{c_i}{c_{\rm tot}}\right|^2, \qquad \frac{{\rm BR}_{hgg}}{{\rm BR}_{hgg}^{\rm SM}} = \left|\frac{c_g^{\rm loop}}{c_g^{\rm SM}c_{\rm tot}}\right|^2 \\ &\frac{{\rm BR}_{h\gamma\gamma}}{{\rm BR}_{h\gamma\gamma}^{\rm SM}} = \left|\frac{c_\gamma^{\rm loop} + c_\gamma^{\rm BMSSM}}{c_\gamma^{\rm SM}c_{\rm tot}}\right|^2, \\ &\frac{{\rm BR}_{h\gamma Z}}{{\rm BR}_{h\gamma Z}^{\rm SM}} = \left|\frac{c_{\gamma Z}^{\rm loop} + c_{\gamma Z}^{\rm BMSSM}}{c_{\gamma Z}^{\rm SM}c_{\rm tot}}\right|^2. \end{aligned} \tag{3.9}$$

The coefficient $c_{\rm tot}$ in Eq. (3.9) can be written as

$$|c_{\rm tot}|^2 = \sum_{i=t,c,b,\tau,Z,W}|c_i|^2{\rm BR}_{hii}^{\rm SM} + \left|\frac{c_g^{\rm loop}}{c_g^{\rm SM}}\right|^2{\rm BR}_{hgg}^{\rm SM}, \tag{3.10}$$

where we have neglected the contributions from, for example, $h \to \gamma\gamma$ and $h \to Z\gamma$ as well as possible invisible decays. Let us now define the inclusive as well as the individual gluon-gluon fusion (ggF), vector boson fusion (VBF), and vector boson associated (VH) production cross sections, normalized with respect to the corresponding SM values,

[5] See also Ref. [23] for additional studies of $h \to Z\gamma$.

$$\frac{\sigma_{\rm incl}}{\sigma_{\rm incl}^{\rm SM}} = \frac{|c_g^{\rm loop}/c_g^{\rm SM}|^2\sigma_{\rm ggF}^{\rm SM} + |c_V|^2(\sigma_{\rm VBF}^{\rm SM} + \sigma_{\rm VH}^{\rm SM})}{\sigma_{\rm ggF}^{\rm SM} + \sigma_{\rm VBF}^{\rm SM} + \sigma_{\rm VH}^{\rm SM}},$$
$$\frac{\sigma_{\rm ggF}}{\sigma_{\rm ggF}^{\rm SM}} = |c_g^{\rm loop}/c_g^{\rm SM}|^2, \qquad \frac{\sigma_{\rm VBF}}{\sigma_{\rm VBF}^{\rm SM}} = \frac{\sigma_{\rm VH}}{\sigma_{\rm VH}^{\rm SM}} = |c_V|^2, \tag{3.11}$$

where we have denoted $c_V = c_Z = c_W$, since the Higgs couplings to Z and W bosons coincide in Eq. (3.2). We can now write, for example, the signal rates in the inclusive and dijet channels of the $h \to \gamma\gamma$ decay mode, again normalized with respect to the SM,

$$R_{\gamma\gamma}^{\rm incl} = \frac{\sigma_{\rm incl}}{\sigma_{\rm incl}^{\rm SM}}\frac{{\rm BR}_{h\gamma\gamma}}{{\rm BR}_{h\gamma\gamma}^{\rm SM}} \qquad R_{\gamma\gamma}^{\rm dijet} = \frac{\epsilon_{\rm ggF}^{\gamma}|c_g^{\rm loop}/c_g^{\rm SM}|^2\sigma_{\rm ggF}^{\rm SM} + \epsilon_{\rm VBF}^{\gamma}|c_V|^2\sigma_{\rm VBF}^{\rm SM} + \epsilon_{\rm VH}^{\gamma}|c_V|^2\sigma_{\rm VH}^{\rm SM}}{\epsilon_{\rm ggF}^{\gamma}\sigma_{\rm ggF}^{\rm SM} + \epsilon_{\rm VBF}^{\gamma}\sigma_{\rm VBF}^{\rm SM} + \epsilon_{\rm VH}^{\gamma}\sigma_{\rm VH}^{\rm SM}}\frac{{\rm BR}_{h\gamma\gamma}}{{\rm BR}_{h\gamma\gamma}^{\rm SM}}, \tag{3.12}$$

where the ϵ^γ coefficients are the selection efficiencies for the different production modes in the dijet-tag category of final states.

B. Decoupling limit

Let us now take the decoupling limit, in which

$$\frac{c_i}{c_i^{\rm SM}} = \frac{c_g^{\rm loop}}{c_g^{\rm SM}} = \frac{c_\gamma^{\rm loop}}{c_\gamma^{\rm SM}} = \frac{c_{\gamma Z}^{\rm loop}}{c_{\gamma Z}^{\rm SM}} = 1, \tag{3.13}$$

where $i = t, c, b, \tau, Z, W$. This implies that $|c_{\rm tot}| = 1$ in Eq. (3.10) and that

$$\frac{\Gamma_{hii}}{\Gamma_{hii}^{\rm SM}} = \frac{{\rm BR}_{hii}}{{\rm BR}_{hii}^{\rm SM}} = \frac{\Gamma_{hgg}}{\Gamma_{hgg}^{\rm SM}} = \frac{{\rm BR}_{hgg}}{{\rm BR}_{hgg}^{\rm SM}} = 1, \tag{3.14}$$

whereas

$$\frac{\Gamma_{h\gamma\gamma}}{\Gamma_{h\gamma\gamma}^{\rm SM}} = \frac{{\rm BR}_{h\gamma\gamma}}{{\rm BR}_{h\gamma\gamma}^{\rm SM}} = \left|1 + \frac{c_{\gamma,\rm dec}^{\rm BMSSM}}{c_\gamma^{\rm SM}}\right|^2,$$
$$\frac{\Gamma_{h\gamma Z}}{\Gamma_{h\gamma Z}^{\rm SM}} = \frac{{\rm BR}_{h\gamma Z}}{{\rm BR}_{h\gamma Z}^{\rm SM}} = \left|1 + \frac{c_{\gamma Z,\rm dec}^{\rm BMSSM}}{c_{\gamma Z}^{\rm SM}}\right|^2 \tag{3.15}$$

for which the coefficients in Eq. (3.7) are given by, in the decoupling limit,

$$c_{\gamma,\rm dec}^{\rm BMSSM} = \frac{\pi v^2 \sin 2\beta}{M^2\alpha_{\rm EM}}(c_1\cos^2\theta_w + c_2\sin^2\theta_w)$$
$$c_{\gamma Z,\rm dec}^{\rm BMSSM} = \frac{\pi v^2 \sin 2\beta}{M^2\alpha_{\rm EM}}(c_2 - c_1)\sin^2\theta_w\cos\theta_w. \tag{3.16}$$

In the decoupling limit, the production cross sections in Eq. (3.11) are all equal to their SM corresponding SM value,

$$\frac{\sigma_{\rm incl}}{\sigma_{\rm incl}^{\rm SM}} = \frac{\sigma_{\rm ggF}}{\sigma_{\rm ggF}^{\rm SM}} = \frac{\sigma_{\rm VBF}}{\sigma_{\rm VBF}^{\rm SM}} = \frac{\sigma_{\rm VH}}{\sigma_{\rm VH}^{\rm SM}} = 1. \tag{3.17}$$

Thus, all signal rates, for any production mode, associated with the channels $h \to ii$, for $i = t, c, b, \tau, Z, W$, as well as $h \to gg$, will be equal to their corresponding SM value. In the $h \to \gamma\gamma$ channel, we see that the signal rates in Eq. (3.12) (as well as any other signal rate in the $h \to \gamma\gamma$ channel) will be given by the corresponding normalized partial width,

$$R_{\gamma\gamma} = R_{\gamma\gamma}^{\rm incl} = R_{\gamma\gamma}^{\rm dijet} = \frac{\Gamma_{h\gamma\gamma}}{\Gamma_{h\gamma\gamma}^{\rm SM}} = \left|1 + \frac{c_{\gamma,\rm dec}^{\rm BMSSM}}{c_\gamma^{\rm SM}}\right|^2, \tag{3.18}$$

and the same for the $h \to Z\gamma$ channel,

$$R_{Z\gamma} = R_{Z\gamma}^{\rm incl} = R_{Z\gamma}^{\rm dijet} = \frac{\Gamma_{hZ\gamma}}{\Gamma_{hZ\gamma}^{\rm SM}} = \left|1 + \frac{c_{Z\gamma,\rm dec}^{\rm BMSSM}}{c_{Z\gamma}^{\rm SM}}\right|^2. \tag{3.19}$$

In summary, in the decoupling limit, all the partial decay widths, except for $\Gamma_{h\gamma\gamma}$ and $\Gamma_{hZ\gamma}$, are all equal to their corresponding SM value. This implies that all the production cross sections, as well as the signal rates in all other channels, are equal to their SM values. Moreover, as seen in Eqs. (3.18) and (3.19), the partial decay widths for $h \to \gamma\gamma$ and $h \to Z\gamma$, normalized with respect to the SM values, coincide with the corresponding signal rates. Hence, in this limit, $\Gamma_{h\gamma\gamma}/\Gamma_{h\gamma\gamma}^{\rm SM}$ and $\Gamma_{hZ\gamma}/\Gamma_{hZ\gamma}^{\rm SM}$ can be compared directly to the measured signal rates. From Eq. (3.16) notice that if $c_1 = c_2$, one can change $\Gamma_{h\gamma\gamma}$ without affecting $\Gamma_{hZ\gamma}$.

IV. RESULTS

We can now evaluate the effect of the operators in Eqs. (2.2) and (2.3) on the mass of the lightest neutral *CP*-even Higgs particle h and on the partial decay widths $\Gamma_{h\gamma\gamma}$ and $\Gamma_{hZ\gamma}$, which directly correspond to the rates in the decoupling limit, as discussed in the previous section.

The Higgs mass in Eq. (2.9) as a function of $\tan\beta$ is displayed in Fig. 1. It is well known that we can accommodate a Higgs mass at 126 GeV by tuning the soft parameters in the loop correction (2.7), but this usually demands a large $\tan\beta$, which we do not consider here (since then additional Yukawa couplings that we do not include become important). With the $d = 5$ operator (2.2), one can easily obtain a value of $m_h \approx 126$ GeV (see Fig. 1); this has an acceptable fine-tuning $\Delta < 30$ [13], even for small $\tan\beta < 10$, which is an otherwise very fine-tuned region of the MSSM. Therefore, the impact of the $d = 5$ operator in reducing fine-tuning is more important than usually thought.

The dimensionless parameter $\epsilon \equiv c_0\mu/M$ measures the extent to which the contribution from $\mathcal{O}_5$ to the mass can be considered perturbative, and for the given numbers, it is below $\epsilon < 0.06$. The mass contributions from $\mathcal{O}_6$ rarely have any visible effect on curves such as those as in Fig. 1,

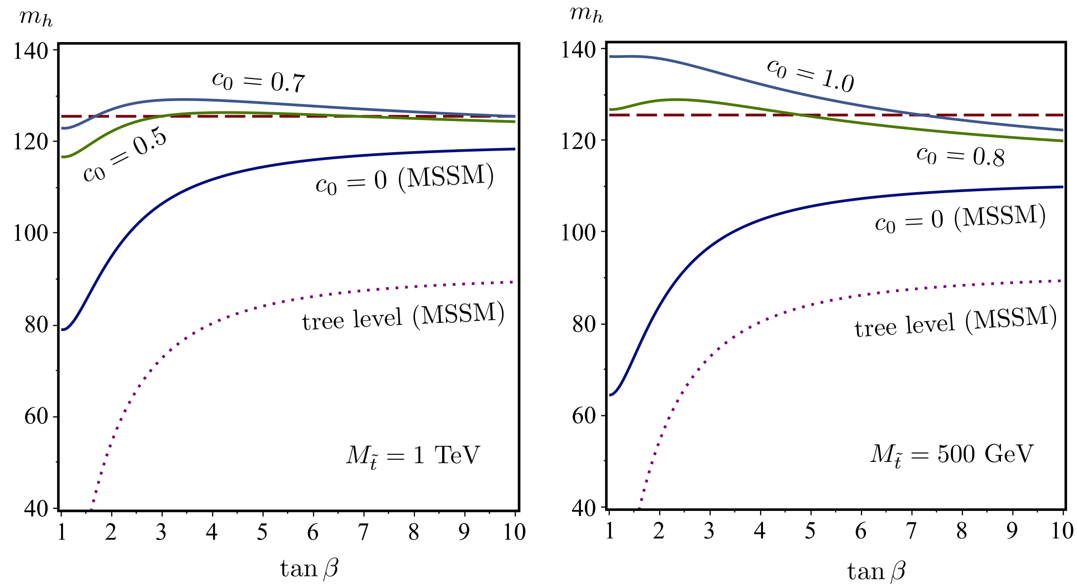

FIG. 1 (color online). The mass of the lightest Higgs particle h as a function of $\tan\beta$, for $M = 5$ TeV, $\mu = 300$ GeV, $m_A = 1$ TeV, no mixing ($X_t = 0$), $M_{\tilde{t}} = 1$ TeV (left panel), $M_{\tilde{t}} = 500$ GeV (right panel), $c_1 = c_2 = -1$, and for various values of the $\mathcal{O}_5$ coefficient c_0. The solid curves include the MSSM loop corrections.

but they are included for completeness. We included only a subset of the loop corrections, which is relevant at low $\tan\beta$, so we expect the curves to differ from the complete result by few GeV only, which we confirmed in our examples using FeynHiggs [25].

Setting the scale M of new physics around 5 TeV as in the example above (which means that the new physics may not be discovered at the LHC), one would like to examine the signal rates for $h \to \gamma\gamma$ and $h \to Z\gamma$. These signal rates $R_{\gamma\gamma}$ and $R_{Z\gamma}$ in the decoupling limit, given in Eqs. (3.18) and (3.19), are shown in Fig. 2 as functions of the coefficients c_1 and c_2 of the operators in Eqs. (2.2) and (2.3). Concerning the $h \to \gamma\gamma$ channel, from the dependence on c_1 and c_2 in $c_{\gamma,\mathrm{dec}}^{\mathrm{BMSSM}}$ in Eq. (3.16), and from the fact that c_{γ}^{SM} is negative in Eq. (3.5), we see that the maximal enhancement of $R_{\gamma\gamma}$ in Eq. (3.18) is obtained for negative values of both c_1 and c_2. In contrast, for positive (and not too large) values of the two coefficients, the $h \to \gamma\gamma$ signal is depleted with respect to the SM prediction, as can be seen in Fig. 2. We emphasize again that we actually do not need both coefficients c_1 and c_2 [corresponding to the $\mathrm{U}(1)_Y$ and $\mathrm{SU}(2)_L$ operators in

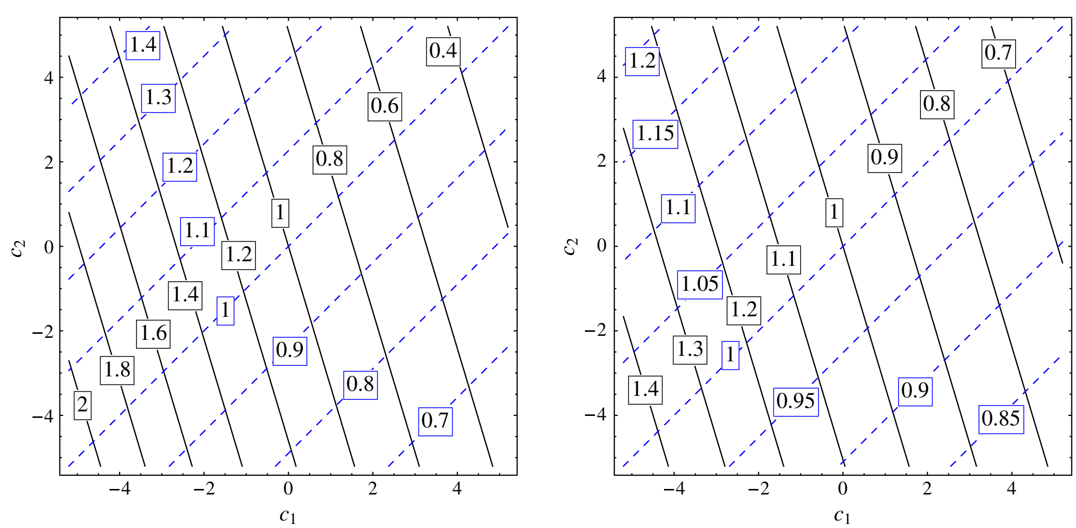

FIG. 2 (color online). The $h \to \gamma\gamma$ signal rate $R_{\gamma\gamma}$ (solid black lines) in Eq. (3.18) and the $h \to Z\gamma$ rate $R_{Z\gamma}$ (dashed blue lines) in Eq. (3.19) are shown as functions of the coefficients c_1 and c_2 of the operators in Eqs. (2.2) and (2.3). In the plots we have set $\tan\beta = 3$ (left panel) and $\tan\beta = 7$ (right panel), $M = 5$ TeV, and we have taken the decoupling limit.

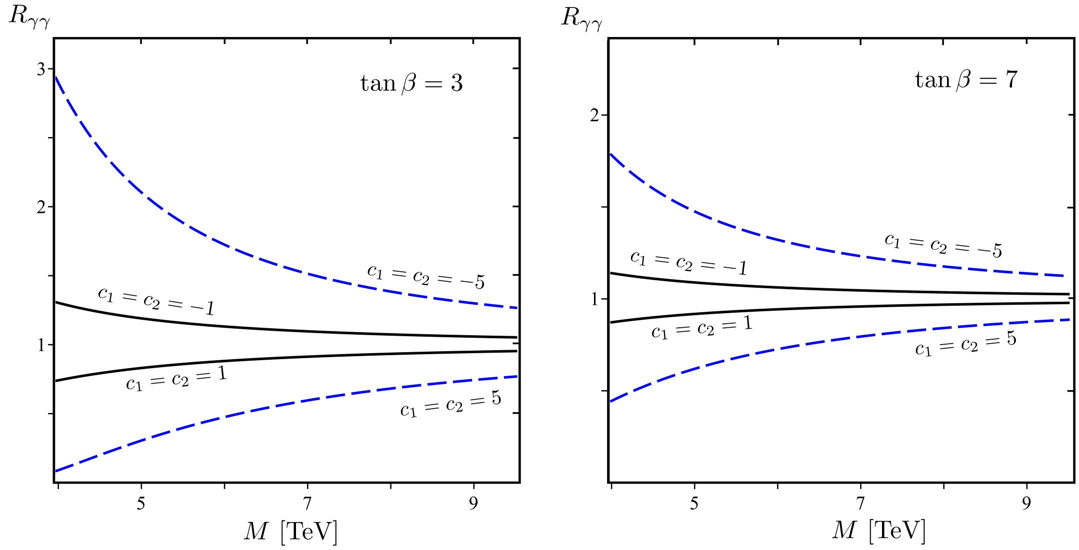

FIG. 3 (color online). The $h \to \gamma\gamma$ signal rate $R_{\gamma\gamma}$ from Eq. (3.18) as functions of the scale M, where we have set the coefficients of the two operators in Eq. (2.3) *equal*, $c_1 = c_2$. We have set $M_{\tilde{t}} = 1$ TeV, $\mu = 300$ GeV, $m_A = 1$ TeV, and $X_t = 0$. Dashed blue lines provide rough estimates of the range of validity of the effective field theory. To contrast with Fig. 2, note that Fig. 2 has a fixed value M = 5 TeV.

Eq. (2.3)] to achieve enhancement; we could set e.g. $c_2 = 0$, but the flexibility this additional parameter affords is useful in the next figure. The maximum allowed enhancement continues to decrease for larger values of $\tan\beta$, unless of course if we simultaneously lower M.

From Eq. (3.16) we see that $c^{\mathrm{BMSSM}}_{\gamma Z,\mathrm{dec}}$ is maximized when c_1 and c_2 have opposite signs. Moreover, since $c^{\mathrm{SM}}_{\gamma Z}$ is positive in Eq. (3.5), in order to achieve an enhancement of $R_{Z\gamma}$ in Eq. (3.18), it is required that $(c_2 - c_1) > 0$. This is seen in Fig. 2, where $R_{Z\gamma}$ is maximized for large positive values for c_2 and large negative values for c_1. Notice that the dependence on the sign of c_1 and c_2 for $R_{\gamma\gamma}$ and $R_{Z\gamma}$ is not specific to this scenario or SUSY. It simply follows from electroweak symmetry breaking, as can be seen in Eq. (3.6).

In Figs. 3 and 4, we show a different representation of the same physics as in Fig. 2, where we fix the coefficients c_1 and c_2 in each curve, and instead vary the overall scale of new physics M. As expected, the effect of the higher-dimensional operators decreases with increasing M, but even for M approaching 10 TeV, there can be some small

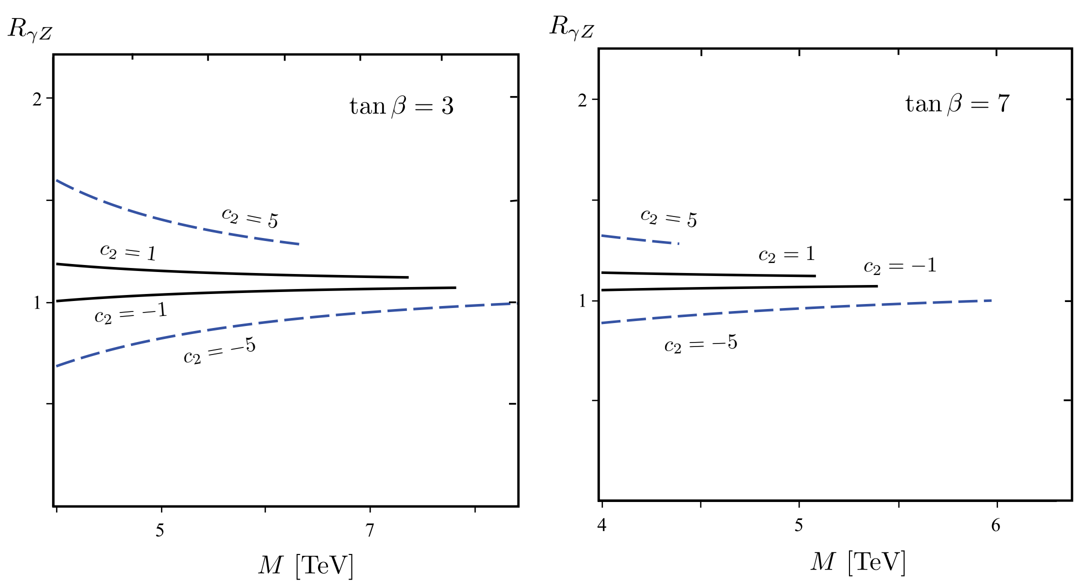

FIG. 4 (color online). The $h \to Z\gamma$ rate $R_{Z\gamma}$ from Eq. (3.19) for various values of the coefficients c_1 and c_2 of the operators in Eqs. (2.2) and (2.3), if we require $R_{\gamma\gamma} = 1.3$ and vary $-5 < c_1 < 5$ for fixed c_2. The curve ends when c_1 goes out of range for $R_{\gamma\gamma} = 1.3$ with the given parameters. We have set $M_{\tilde{t}} = 1$ TeV, $\mu = 300$ GeV, $m_A = 1$ TeV, and $X_t = 0$. Dashed blue lines provide rough estimates of the range of validity of the effective field theory expansion. Again, to contrast with Fig. 2, note that Fig. 2 has a fixed value M = 5 TeV.

effect. This perhaps somewhat counterintuitive behavior is simply because the relevant SM couplings are small to begin with, as emphasized in the introduction. Since the "new physics" that generated these operators comes from a scale around or larger than 5 TeV, it will not be within easy reach of the LHC.

In Fig. 4 we have illustrated the behavior of $R_{\gamma Z}$ if we *require*, for example, $R_{\gamma\gamma} = 1.3$ (i.e. interpret the diphoton excess as signal) and vary $-5 < c_1 < 5$, so each curve represents a particular value of c_2 and ends at some upper bound value of M where it is no longer possible to achieve the prescribed value of $R_{\gamma\gamma} = 1.3$. We see that above $\tan\beta = 5$ or so, one would have to rely on the scale M being not too far above 5 TeV.

With hindsight, it may appear that the analysis of the effects of $\mathcal{O}_5$ in Eq. (2.2) and $\mathcal{O}_6$ in Eq. (2.3) could have been performed mostly independently of each other. To be clear, we did not assume this; as a matter of principle, we always include the contribution to, e.g., the mixing angle α from Eq. (2.2) when computing the effects of Eq. (2.3). But we emphasize that the $R_{\gamma\gamma}$ contributions arising from the dimension-six operators in Eq. (2.3) are maximized for small $\tan\beta$; see Eq. (3.16). Therefore, since the usual MSSM tree-level contribution to the Higgs mass is minimized for small $\tan\beta$, the contribution from the dimension-five operator to the tree-level Higgs mass is crucial in order to accommodate a 126 GeV Higgs mass, as is seen in Fig. 1.

V. GENERATING THE EFFECTIVE OPERATORS FROM UNDERLYING PHYSICS

The natural question is then what new physics could generate the effective operators discussed. In this section we discuss a simple example of an underlying model from which both $\mathcal{O}_5$ and $\mathcal{O}_2$ arise simultaneously in the low-energy effective theory, upon integrating out some massive supersymmetric degrees of freedom. Consider a model that contains the following superpotential, involving a massive gauge singlet chiral superfield[6] Σ,

$$W \supset (\mu + \lambda\Sigma) H_d \cdot H_u + \frac{1}{2}\mu_S \Sigma^2, \qquad (5.1)$$

and a gauge kinetic function τ that depends on Σ,

$$\tau(\Sigma)\mathrm{Tr}(W_\alpha W^\alpha) \quad \text{with} \quad \tau(\Sigma) \supset \frac{\Sigma}{\Lambda}, \qquad (5.2)$$

where Λ is a dimensionful suppression scale. If the SUSY mass μ_S is sufficiently large with respect to the energy scale under consideration, then Σ can be integrated out supersymmetrically via its holomorphic equation of motion, which sets

$$\Sigma = -\frac{\lambda}{\mu_S} H_d \cdot H_u - \frac{1}{\mu_S \Lambda}\mathrm{Tr}(W_\alpha W^\alpha) + \cdots, \qquad (5.3)$$

where the dots stand for higher-dimensional terms (further suppressed by μ_S, Λ). By inserting this solution back into the original Lagrangian, we obtain the following terms:

$$\int d^2\theta \Big(\mu H_d \cdot H_u - \frac{\lambda^2}{2\mu_S}(H_d \cdot H_u)^2 - \frac{\lambda}{\mu_S \Lambda} H_d \cdot H_u \mathrm{Tr}(W_\alpha W^\alpha) \Big) + \int d^4\theta \Big(\Big| \frac{\lambda}{\mu_S} \Big|^2 (H_d \cdot H_u)^\dagger (H_d \cdot H_u) \Big), \qquad (5.4)$$

where we have included operators up to dimension six. We see that operators $\mathcal{O}_5$, $\mathcal{O}_6$ of Eqs. (2.2) and (2.3) were simultaneously generated as a consequence of integrating out Σ.

The dimension-six Kähler potential operator in the second line of Eq. (5.4) gives corrections to the quartic Higgs scalar potential and hence to the tree-level Higgs mass. However, in comparison to our dimension-five operator in the first line of Eq. (5.4), this operator is suppressed by a higher power of μ_S. As long as μ_S is sufficiently large, the corrections from this dimension-six operator will be smaller in size with respect to the corrections from the dimension-five operator.

It should be acknowledged that this example is not renormalizable since the gauge kinetic term (5.2) has dimension $d = 5$. To have a renormalizable microscopic model, one should also specify the degrees of freedom responsible for generating this $d = 5$ operator. Nevertheless, this operator with a moduli-dependent gauge kinetic function is generically present in models derived from supergravity or string theory.

Finally in order to connect with the discussion in the rest of the paper, we should assume that the dimension-five operators in Eq. (5.2) only involve the $\mathrm{U}(1)_Y$ and $\mathrm{SU}(2)_L$ gauge field strength, and not also the $\mathrm{SU}(3)_C$ one that could in principle be also present. In string models, something along these lines could be achieved by, for example, considering a brane model in which the dimensionality of the branes that give rise to the $\mathrm{U}(1)_Y$ and $\mathrm{SU}(2)_L$ gauge fields is different from those that give rise to the $\mathrm{SU}(3)_C$ gauge field. In this way, since the $\mathrm{U}(1)_Y$ and $\mathrm{SU}(2)_L$ branes and the $\mathrm{SU}(3)_C$ branes, respectively, wrap different cycles of the internal geometry, they depend on different gauge singlet moduli fields associated with the different cycles.

VI. CONCLUSIONS

Recent LHC data on the Higgs mass and its couplings to the photon, and the negative SUSY searches, present increasing difficulties for MSSM-like models to naturally accommodate a Higgs mass near 126 GeV, without undue fine-tuning, and a potential enhancement of the $h \to \gamma\gamma$

[6]Gauge singlet fields with a supersymmetric mass term appear in general versions of the NMSSM [12–14].

partial decay rate, without affecting other partial decay widths. Motivated by these observations, in this work we investigated whether supersymmetric effects beyond the MSSM could simultaneously accommodate these results. Using an effective approach, we identified two effective operators of dimensions $d = 5$ and $d = 6$ that can address these problems and give the leading-order corrections to the Higgs quartic coupling and the Higgs coupling to photons, respectively.

We showed that the MSSM with small, supersymmetric corrections due to these effective operators can simultaneously naturally accommodate a Higgs boson with a mass near 126 GeV, an enhanced Higgs coupling to photons (and also $Z\gamma$) relative to the SM expectation, and finally SM-like Higgs couplings to the other SM particles. The scale of the supersymmetric effective operators is in the region of 5 TeV or even larger and is therefore possibly not within the LHC reach. The corrections from the dimension-six operators to the Higgs coupling to photons (and $Z\gamma$) are maximized for small $\tan\beta$ which is also the region where the dimension-five operator produces the most dramatic effect relative to the MSSM. This suggests that it is natural to consider these operators together, and this is further supported by the fact that both of them can be generated simultaneously by an underlying model, as we showed.

There remains the question of how the existence of these operators can be tested. Let us assume that the signal rate in the $h \to \gamma\gamma$ channel is confirmed to be higher than the SM expectation while the signal rates in all the other channels coincide with the SM values. If at the same time, one can rule out light stau sleptons in the mass range (of 150 GeV or so) needed in order to enhance the diphoton signal with the correct amount, this would cause a real problem for the MSSM, and physics beyond the MSSM will be required. Should the excess go away when further data is analyzed, our results will remain useful to provide bounds on the scale M of supersymmetric new physics beyond the MSSM. Either way, this suggests that the diphoton rate is a useful, sensitive probe in this context.

ACKNOWLEDGMENTS

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Note added.—After this paper was finished, new experimental analyses appeared [26] that leave the diphoton excess uncertain. At this point the CMS Collaboration no longer reports any significant excess, but the ATLAS Collaboration still reports an excess. The situation with respect to this decay channel remains uncertain, and it would be interesting to perform a more extensive analysis of our results when the experiments begin again.

APPENDIX

1. Details concerning the Lagrangian

Here we derive the Lagrangian of Sec. II. Unlike in the text, we also include SUSY-breaking effects associated with operators $\mathcal{O}_5$, $\mathcal{O}_6$, by using the spurion field. The starting point is

$$\mathcal{L} = \int d^4\theta \sum_{i=u,d} (1 - m_i^2\theta^2\bar{\theta}^2) H_i^\dagger e^{V_i} H_i + \left\{ \int d^2\theta \mu (1 + B\theta^2) H_d \cdot H_u + \text{H.c.} \right\} + \mathcal{O}_5 + \mathcal{O}_6. \tag{A1}$$

The superfield components are $V_i = (\lambda_i, V_{i,\mu}, D_i^a/2)$, $H_i = (h_i, \psi_i, F_i)$. Also $H_d \cdot H_u = \epsilon^{ij} H_d^i H_u^j$, with $\epsilon^{ij}\epsilon^{kj} = \delta^{ik}$; $\epsilon^{ij}\epsilon^{kl} = \delta^{ik}\delta^{jl} - \delta^{il}\delta^{jk}$, $\epsilon^{12} = 1$, $h_d \cdot h_u = h_d^0 h_u^0 - h_d^- h_u^+$. Further,

$$\begin{aligned}\mathcal{O}_5 &= \frac{1}{M} \int d^2\theta (c_0 + c_0'\theta^2)(H_u \cdot H_d)^2 + \text{H.c.} \\ &= \frac{c_0}{M}[2(h_u \cdot h_d)(h_u \cdot F_d + F_u \cdot h_d - \psi_u \cdot \psi_d) \\ &\quad - (h_u \cdot \psi_d + \psi_u \cdot h_d)^2] + \frac{c_0'}{M}(h_u \cdot h_d)^2 + \text{H.c.},\end{aligned} \tag{A2}$$

and

$$\begin{aligned}\mathcal{O}_6 &= \frac{1}{M^2} \sum_{s=1,2} \frac{1}{16 g_s^2 \kappa} \int d^2\theta (c_s + c_s'\theta^2) \mathrm{Tr}(W^\alpha W_\alpha)_s (H_u \cdot H_d) + \text{H.c.} \\ &= \sum_{s=1,2} \frac{c_s}{4M^2} \left\{ (h_u \cdot h_d) \left[i(\lambda_s^a \sigma^\mu \mathcal{D}_\mu \bar{\lambda}_s^a - \mathcal{D}_\mu \bar{\lambda}_s^a \bar{\sigma}^\mu \lambda_s^a) + D_s^a D_s^a - \frac{1}{2}\left(F_s^{a\mu\nu} F_{s\mu\nu}^a + \frac{i\epsilon^{\mu\nu\rho\sigma}}{2} F_{s\mu\nu}^a F_{s\rho\sigma}^a \right) \right] \right. \\ &\quad \left. - \sqrt{2}(h_u \cdot \psi_d + \psi_u \cdot h_d)(\lambda_s^a D_s^a + \sigma^{\mu\nu}\lambda_s^a F_{s\mu\nu}^a) + (h_u \cdot F_d + F_u \cdot h_d - \psi_u \cdot \psi_d)\lambda_s^a \lambda_s^a \right\} + \frac{c_s'}{4M^2}(h_u \cdot h_d)(\lambda_s^a \lambda_s^a) + \text{H.c.}\end{aligned} \tag{A3}$$

Above we introduced $\mathcal{D}_\mu \bar{\lambda}^a = \partial_\mu \bar{\lambda}^a - g t^{abc} V_\mu^b \bar{\lambda}^c$ for covariant derivatives of the gauginos.

From $\mathcal{L}$ one finds the equations of motion for the auxiliary fields of Higgs superfields:

$$\begin{aligned} F_d^{*q} &= -\epsilon^{qp} h_u^p \Big[\mu + 2\frac{c_0}{M}(h_d \cdot h_u) - \frac{c_2}{4M^2}\lambda_2^a \lambda_2^a - \frac{c_1}{4M^2}\lambda_1^2 \Big] \\ F_u^{*q} &= -\epsilon^{pq} h_d^p \Big[\mu + 2\frac{c_0}{M}(h_d \cdot h_u) - \frac{c_2}{4M^2}\lambda_2^a \lambda_2^a - \frac{c_1}{4M^2}\lambda_1^2 \Big], \end{aligned} \tag{A4}$$

where q is a $SU(2)_L$ doublet index. For the auxiliary fields of the vector superfields, we find

$$D_2^a = -\Big[g_2(h_d^\dagger T^a h_d + h_u^\dagger T^a h_u)\Big(1 - \frac{c_2}{2M^2}(h_u \cdot h_d + \mathrm{H.c.})\Big) - \frac{\sqrt{2}c_2}{4M^2}((h_u \cdot \psi_d + \psi_u \cdot h_d)\lambda_2^a + \mathrm{H.c.})\Big] \tag{A5}$$

$$D_1 = -\Big[g_1\Big(h_d^\dagger\Big(-\frac{1}{2}\Big)h_d + h_u^\dagger\Big(\frac{1}{2}\Big)h_u\Big)\Big(1 - \frac{c_1}{2M^2}(h_u \cdot h_d + \mathrm{H.c.})\Big) - \frac{\sqrt{2}c_1}{4M^2}(h_u \cdot \psi_d + \psi_u \cdot h_d)\lambda_1^a + \mathrm{H.c.})\Big], \tag{A6}$$

with $T^a = \sigma^a/2$. The squares become

$$\begin{aligned} D_2^a D_2^a = \frac{g_2^2}{4}\Big[1 - \Big(\frac{c_2}{2M^2}h_u \cdot h_d + \mathrm{H.c.}\Big)\Big]^2 [(|h_d|^2 - |h_u|^2)^2 + 4|h_d^\dagger h_u|^2] \\ - \frac{\sqrt{2}}{2} g_2 [h_d^\dagger T^a h_d + h_u^\dagger T^a h_u]\Big[\frac{c_2}{2M^2}(h_u \cdot \psi_d + \psi_u \cdot h_d)\lambda_2^a + \mathrm{H.c.}\Big], \end{aligned} \tag{A7}$$

$$\begin{aligned} D_1^2 = \frac{g_1^2}{4}\Big[1 - \Big(\frac{c_1}{2M^2}h_u \cdot h_d + \mathrm{H.c.}\Big)\Big]^2 (|h_d|^2 - |h_u|^2)^2 - \frac{\sqrt{2}}{2} g_1 \Big[h_d^\dagger\Big(-\frac{1}{2}\Big)h_d + h_u^\dagger\Big(\frac{1}{2}\Big)h_u\Big] \\ \times \Big[\frac{c_1}{2M^2}(h_u \cdot \psi_d + \psi_u \cdot h_d)\lambda_1 + \mathrm{H.c.}\Big]. \end{aligned} \tag{A8}$$

$\mathcal{O}_5$ and $\mathcal{O}_6$ and Eqs. (A4)–(A8) give the corrections to the MSSM Higgs Lagrangian. Using the corrected auxiliary fields in the usual MSSM Higgs Lagrangian, additional terms suppressed by $1/M$ and $1/M^2$ are generated. The full on-shell Lagrangian is then

$$\mathcal{L} = \mathcal{L}_D + \mathcal{L}_F + \mathcal{L}_1 + \mathcal{L}_2 + \mathcal{L}_3 + \mathcal{L}_{\mathrm{SSB}}. \tag{A9}$$

Eliminating the D-dependent terms in $\mathcal{L}$, one finds [see Eqs. (A5)–(A8)]

$$\mathcal{L}_D = \sum_{s=1,2} -\frac{1}{2} D_s^a D_s^a \Big[1 + \frac{c_s}{2M^2}(h_u \cdot h_d + \mathrm{H.c.})\Big] \tag{A10}$$

and uses Eq. (A8). Eliminating the F-dependent terms in $\mathcal{L}$ gives $\mathcal{L}_F$:

$$-\mathcal{L}_F \equiv |F_d|^2 + |F_u|^2 = \Big|\mu + 2\frac{c_0}{M}h_d \cdot h_u\Big|^2 (|h_d|^2 + |h_u|^2) + \Big[\mu\Big(-\frac{c_2}{4M^2}\lambda_2^a\lambda_2^a - \frac{c_1}{4M^2}\lambda_1^2\Big)(|h_d|^2 + |h_u|^2) + \mathrm{H.c.}\Big]. \tag{A11}$$

Apart from auxiliary field contributions, there are also terms in the Lagrangian with space-time derivatives, which contribute to the kinetic terms for Weyl fermions $\psi_{u,d}$, $\lambda_{1,2}^a$ when the neutral singlet $h_{u,d}^0$ components of $h_{u,d}$ acquire a VEV:

$$\mathcal{L}_1 = \frac{c_2}{4M^2}(h_u \cdot h_d)[i(\lambda_2^a \sigma^\mu \mathcal{D}_\mu \bar{\lambda}_2^a - \mathcal{D}_\mu \bar{\lambda}_2^a \bar{\sigma}^\mu \lambda_2^a)] + \mathrm{H.c.} + (2 \to 1). \tag{A12}$$

When the Higgs neutral singlets acquire a VEV, these terms produce wave function renormalization of Weyl kinetic terms and a threshold correction to gauge couplings g_1 and g_2.

There are also terms contributing to fermion masses when the Higgs fields acquire VEVs,

$$\mathcal{L}_2 = \frac{c_2'}{4M^2}(h_u \cdot h_d)(\lambda_2^a \lambda_2^a) + \frac{c_1'}{4M^2}(h_u \cdot h_d)(\lambda_1 \lambda_1) + \frac{c_0}{M}[2(h_u \cdot h_d)(-\psi_u \cdot \psi_d) - (h_u \cdot \psi_d + \psi_u \cdot h_d)^2] + \mathrm{H.c.} \tag{A13}$$

Further, there are some interaction terms generated:

$$\mathcal{L}_3 = \left\{ \frac{c_2}{4M^2}\left[-\frac{1}{2}(h_u \cdot h_d)\left(F_2^{a\mu\nu}F^a_{2\mu\nu} + \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}F^a_{2\mu\nu}F^a_{2\rho\sigma}\right) - \sqrt{2}(h_u \cdot \psi_d + \psi_u \cdot h_d)\sigma^{\mu\nu}\lambda_2^a F^a_{2\mu\nu} - \psi_u \cdot \psi_d \lambda_2^a \lambda_2^a \right] + (2 \to 1) + \text{H.c.} \right\}. \tag{A14}$$

Finally, the Lagrangian contains (F- and D-independent) corrections from supersymmetry breaking due to spurion dependence in the dimension-5 operator as well as the usual soft terms of the MSSM. All these together give $\mathcal{L}_{\text{SSB}}$:

$$\mathcal{L}_{\text{SSB}} = -V_{\text{SSB}} = \left[\frac{c_0'}{M}(h_u \cdot h_d)^2 + \mu B(h_d \cdot h_u) + \text{H.c.} \right] - m_d^2|h_d|^2 - m_u^2|h_u|^2.$$

This concludes the presentation of the Lagrangian to $1/M^2$ order. From $\mathcal{L}$ we find the scalar potential V_h for the Higgs sector shown in the text, Eq. (2.6), in which as usual $h_{u,d}$ denote SU(2) doublets. From this one obtains

$$m_h^2 = m^2_{h,\text{loop}} + \Delta m_h^2, \tag{A15}$$

with

$$m^2_{h,\text{loop}} = \frac{1}{2}\{m_A^2 + m_Z^2 + \delta m_Z^2 \sin^2\beta - \sqrt{w}\} \qquad w \equiv [(m_A^2 - m_Z^2)\cos 2\beta + \delta m_Z^2 \sin^2\beta]^2 + \sin^2 2\beta (m_A^2 + m_Z^2)^2 \tag{A16}$$

and

$$\Delta m_h^2 = f_1\left(2\mu\frac{c_0}{M}\right) + f_2\left(-2\frac{c_0'}{M}\right) + f_3\left(2\mu\frac{c_0}{M}\right)^2 + f_4\left(-2\frac{c_0'}{M}\right)^2 + f_5\left(2\mu\frac{c_0}{M}\right)\left(-2\frac{c_0'}{M}\right) + f_6\left(g_1^2\frac{c_1}{M^2} + g_2^2\frac{c_2}{M^2}\right) + \mathcal{O}\left(\frac{1}{M^3}\right), \tag{A17}$$

where

$$\begin{aligned}
f_1 &= v^2 \sin 2\beta\left\{1 + \frac{(m_A^2 + m_Z^2)}{\sqrt{w}}\right\} \qquad f_2 = \frac{v^2}{2}\left\{1 - \frac{\cos 2\beta}{\sqrt{w}}[(m_A^2 - m_Z^2)\cos 2\beta + m_Z^2 \delta \sin^2\beta]\right\} \\
f_3 &= \frac{v^4}{4\mu^2}\sin^2 2\beta + \frac{v^4}{\sqrt{w}}\left\{-1 + \frac{1}{2\mu^2}(m_A^2 + m_Z^2)\sin^2 2\beta\right\} + \frac{1}{w^{3/2}}(m_A^2 + m_Z^2)^2 v^4 \sin^2 2\beta \\
f_4 &= -\frac{v^4}{16 w^{3/2}}(m_A^2 + m_Z^2)^2 \sin^2 4\beta \qquad f_5 = -\frac{v^4}{4w^{3/2}}(m_A^2 + m_Z^2)[\delta m_Z^2 + (2m_A^2 - (2+\delta)m_Z^2)\cos 2\beta]\sin 4\beta \\
f_6 &= \frac{v^4}{32}\sin 2\beta\left[1 + \frac{1}{4\sqrt{w}}[8m_A^2 - (4 + 3\delta)m_Z^2 + 6\delta m_Z^2 \cos 2\beta + 3(4m_A^2 - \delta m_Z^2)\cos 4\beta]\right]
\end{aligned} \tag{A18}$$

and finally

$$m_A^2 = \frac{2B\mu}{\sin 2\beta} - \frac{v^2}{\sin 2\beta}\left(2\mu\frac{c_0}{M}\right) + \left(2\frac{c_0'}{M}\right)v^2 - \frac{v^4}{32}\frac{\cos^2 2\beta}{\sin 2\beta}\left(\frac{g_1^2 c_1}{M^2} + \frac{c_2 g_2^2}{M^2}\right) + \mathcal{O}\left(\frac{1}{M^3}\right).$$

2. One-loop form factors

The form factors in Eq. (3.4) are given by

$$\begin{aligned}
&\mathcal{A}_g^{(t)} = \frac{3}{4}\mathcal{A}_{1/2}(\tau_t), \qquad \mathcal{A}_g^{(b)} = \frac{3}{4}\mathcal{A}_{1/2}(\tau_b), \qquad \mathcal{A}_\gamma^{(W)} = \mathcal{A}_1(\tau_W), \qquad \mathcal{A}_\gamma^{(t)} = N_c Q_t^2 \mathcal{A}_{1/2}(\tau_t), \\
&\mathcal{A}_{Z\gamma}^{(W)} = \cos\theta_w A_1(\tau_W, \lambda_W), \qquad \mathcal{A}_{Z\gamma}^{(t)} = N_c \frac{Q_t(2T_3^{(t)} - 4Q_t \sin^2\theta_w)}{\cos\theta_w} A_{1/2}(\tau_t, \lambda_t),
\end{aligned} \tag{A19}$$

where $\tau_i = 4m_i^2/m_h^2$, $\lambda_i = 4m_i^2/m_Z^2$, $N_c = 3$, $Q_t = 2/3$, $T_3^{(t)} = 1/2$ and

$$\mathcal{A}_{1/2}(\tau) = 2\tau^2[\tau^{-1} + (\tau^{-1} - 1)f(\tau^{-1})], \qquad \mathcal{A}_1(\tau) = -\tau^2[2\tau^{-2} + 3\tau^{-1} + 3(2\tau^{-1} - 1)f(\tau^{-1})],$$
$$\mathcal{A}_{1/2}(\tau, \lambda) = I_1(\tau, \lambda) - I_2(\tau, \lambda), \qquad \mathcal{A}_1(\tau, \lambda) = 4(3 - \tan^2\theta_w)I_2(\tau, \lambda) + [(1 + 2\tau^{-1})\tan^2\theta_w - (5 + 2\tau^{-1})]I_1(\tau, \lambda), \tag{A20}$$

where

$$\begin{aligned} I_1(\tau, \lambda) &= \frac{\tau\lambda}{2(\tau - \lambda)} + \frac{\tau^2\lambda^2}{2(\tau - \lambda)^2}[f(\tau^{-1}) - f(\lambda^{-1})] + \frac{\tau^2\lambda}{(\tau - \lambda)^2}[g(\tau^{-1}) - g(\lambda^{-1})], \\ I_2(\tau, \lambda) &= -\frac{\tau\lambda}{2(\tau - \lambda)}[f(\tau^{-1}) - f(\lambda^{-1})], \end{aligned} \tag{A21}$$

and

$$f(x) = \begin{cases} \arcsin^2\sqrt{x} & x \le 1 \\ -\frac{1}{4}\left[\log\frac{1+\sqrt{1-x^{-1}}}{1-\sqrt{1-x^{-1}}} - i\pi\right]^2 & x > 1, \end{cases} \tag{A22}$$

$$g(x) = \begin{cases} \sqrt{x^{-1} - 1}\arcsin\sqrt{x} & x \le 1 \\ \frac{\sqrt{1-x^{-1}}}{2}\left[\log\frac{1+\sqrt{1-x^{-1}}}{1-\sqrt{1-x^{-1}}} - i\pi\right]^2 & x > 1. \end{cases} \tag{A23}$$

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