Radiative corrections to Higgs boson masses for the MSSM Higgs potential with dimension-six operators

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In the framework of the effective field theory approach to heavy supersymmetry radiative corrections in the Higgs sector of the minimal supersymmetric standard model (MSSM) for the effective potential decomposition up to the dimension-six operators are calculated. Symbolic expressions for the threshold corrections induced by F- and D-soft supersymmetry breaking terms are derived, and the Higgs boson mass spectrum respecting the condition $m_h = 125$ GeV for the lightest CP-even scalar is evaluated.

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I. INTRODUCTION

The absence of a signal of supersymmetric partners at the LHC up to the mass range of 1–2 TeV [1] increased an interest in the "heavy supersymmetry" scenarios [2] of the minimal supersymmetric standard model (MSSM), where the condition $m_h = 125$ GeV for the lightest CP-even scalar state, perhaps, observed by the ATLAS and CMS Collaborations [3] is respected explicitly in the MSSM parameter space. Large radiative corrections to the MSSM two-Higgs doublet sector which raise up m_h from the maximal tree-level value of m_Z to the observable value of 125 GeV appear due to large values of soft supersymmetry breaking parameters, which are associated with large masses of third generation quark supersymmetric partners, associated with large mixing of supersymmetric partners, and restricted from the above by the availability of the perturbative regime. For this reason, acceptable domains of the MSSM parameter space are rather limited [4] although there are several variants of such "fine-tuning." To ease tensions of parametric scenarios of the MSSM, two ways of action are appropriate: first, more precise calculations of radiative corrections at higher loops/decomposition of the effective potential in the higher inverse powers of M_S [i.e., including effective operators $1/M_S^n O(\Phi^{n+4})$ in the decomposition of the Coleman-Weinberg type potential]; second, the transition to extensions of the MSSM. For example, extensions of the MSSM where the superpotential includes an additional chiral singlet field [next to minimal supersymmetric standard model (NMSSM) [5]], or more chiral fields, are known. It is assumed that some new physics beyond the MSSM exists at an energy scale that is not too far away. Probably, such a scale of the order of 10^1 TeV is somewhat higher than the mass scale of superpartners M_S.

In the framework of a picture where MSSM is a low-energy limit of an extended theory (not only NMSSM, for example, supersymmetric grand unification models [6] or supersymmetric left-right models [7]) all possible effective operators of higher dimension should be introduced with the following separation of the observables which are sensitive to effects of the extended theory for phenomenological analysis. The effective Lagrangian of the MSSM extension can be written as a sum of operators suppressed by inverse powers of the new physics scale M^{-1} and M^{-2}, each of which is $SU(3)_c \times SU(2)_L \times U(1)_Y$ invariant and respects R parity. In the extended theory, such operators are generated either at the tree level or at the loop level. It was expected that contributions of the tree-level operators to the specific observables were more important because the loop-level operators have additional suppression factors proportional to $1/16\pi^2$. However, additional enhancements by large MSSM parameters (such as $\tan\beta = v_2/v_1$, which can compensate also for an extra power of the mass scale M) make the situation with various contributions rather nontrivial. A number of studies prior to the Higgs boson discovery can be found in the literature. A complete list of the tree-level dimension-five and dimension-six effective operators can be found in [8]. Note that supersymmetry (SUSY) restricts possible effective operator categories; for example, no operators of dimension-five involving Higgs-Higgsino supermultiplet and gauge-gaugino supermultiplet exist since no gauge invariant form can be constructed using three MSSM chiral superfields. Analogously, no operators of dimension-six involving Higgs-Higgsino supermultiplet exist because operators of this type must contain five chiral superfields; apparently, such forms violate gauge invariance. Various aspects related to extensions by the dimension-five/dimension-six operators were systematically analyzed in [9].

As mentioned above, radiative corrections coming from the loop diagrams with top quark and top superpartner are very important [10,11] for both large $\tan\beta$ and small $\tan\beta$

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parameters. The tree-level mass of the lightest CP-even state h is maximized at large $\tan\beta$. For small trilinear parameters $A_{t,b}$ and large stop mass scale M_S when A_t/M_S and μ/M_S are less than one (in other words, in the case of moderate stop mixing parameter $X_t = A_t - \mu\cot\beta$), the correction to m_h at the one loop is controlled by the logarithm $\log M_S/m_{\rm top}$ which is large enough for M_S of the order of 10–100 TeV. For large trilinear parameters $A_{t,b}$ (or in the case of the large stop mixing parameter) the correction is maximized at $A_t = M_S\sqrt{6}$ (so-called "maximal mixing scenario" at the one loop), and much smaller M_S values of the order of 1 TeV are appropriate. At $\tan\beta \sim 1$ or even smaller, large mixing may appear due to large Higgs superfield mass parameter μ of about 10 TeV. The nontrivial interplay of $A_{t,b}, \mu, M_S$, and $\tan\beta$ parameters at the level of the one-loop resummed Higgs potential was analyzed in detail for the potential decomposition in the inverse powers of M_S up to operators of dimension-four. The case of small mass splittings for quark superpartners [12] was generalized for the situation when each stop and sbottom is independently decoupled at its specific mass scale [13] for some special MSSM effective potentials. Note that the two-loop effects may be included by using a renormalization group improvement of the effective potential [14]. The scale dependence of the one-loop result is reduced if the two-loop renormalization group improvement of the one-loop effective potential is accounted for [12,15].

In this paper the effective MSSM Higgs potential decomposition up to operators of the dimension-six involving scalars only is considered. The contribution of the dimension-six operators to observables can be separated insofar, as already mentioned above, the dimension-six operators involving only scalar isodoublets appear at the loop level only. In Sec. II the mass basis for the extended Higgs sector is constructed. In Sec. III analytical expressions for the threshold corrections are derived, and some numerical evaluations for the mass spectra are performed.

II. MASS BASIS FOR THE CASE OF EFFECTIVE POTENTIAL WITH THE DIMENSION-SIX TERMS

In this section we construct the basis for the mass states of physical scalars following [16], where the case of dimension-four operators has been considered. Two Higgs doublets of the form

$$\Phi_i = \begin{pmatrix} \phi_i^+(x) \\ \phi_i^0(x) \end{pmatrix} = \begin{pmatrix} -i\omega_i^+ \\ \frac{1}{\sqrt{2}}(v_i + \eta_i + i\chi_i) \end{pmatrix}, \qquad i = 1, 2, \tag{1}$$

are used to define the general two-Higgs doublet potential. Calculation of quantum corrections to the Higgs potential requires a resummation of Feynman diagrams to all orders of perturbation theory. Because of loop graphs, self-interactions of Higgs fields acquire additional higher-order terms. At the one loop, the resummed potential can be written as

$$U = U^{(2)} + U^{(4)} + U^{(6)} + \cdots, \tag{2}$$

where the upper index shows the operator dimension in fields,

$$U^{(2)} = -\mu_1^2(\Phi_1^\dagger\Phi_1) - \mu_2^2(\Phi_2^\dagger\Phi_2) - [\mu_{12}^2(\Phi_1^\dagger\Phi_2) + \text{H.c.}], \tag{3}$$

$$\begin{aligned} U^{(4)} = {}& \lambda_1(\Phi_1^\dagger\Phi_1)^2 + \lambda_2(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) \\ & + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + [\lambda_5/2(\Phi_1^\dagger\Phi_2)(\Phi_1^\dagger\Phi_2) \\ & + \lambda_6(\Phi_1^\dagger\Phi_1)(\Phi_1^\dagger\Phi_2) + \lambda_7(\Phi_2^\dagger\Phi_2)(\Phi_1^\dagger\Phi_2) + \text{H.c.}], \end{aligned} \tag{4}$$

$$\begin{aligned} U^{(6)} = {}& \kappa_1(\Phi_1^\dagger\Phi_1)^3 + \kappa_2(\Phi_2^\dagger\Phi_2)^3 + \kappa_3(\Phi_1^\dagger\Phi_1)^2(\Phi_2^\dagger\Phi_2) + \kappa_4(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2)^2 \\ & + \kappa_5(\Phi_1^\dagger\Phi_1)(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + \kappa_6(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) \\ & + [\kappa_7(\Phi_1^\dagger\Phi_2)^3 + \kappa_8(\Phi_1^\dagger\Phi_1)^2(\Phi_1^\dagger\Phi_2) + \kappa_9(\Phi_1^\dagger\Phi_1)(\Phi_1^\dagger\Phi_2)^2 \\ & + \kappa_{10}(\Phi_1^\dagger\Phi_2)^2(\Phi_2^\dagger\Phi_2) + \kappa_{11}(\Phi_1^\dagger\Phi_2)^2(\Phi_2^\dagger\Phi_1) + \kappa_{12}(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_2)^2 \\ & + \kappa_{13}(\Phi_1^\dagger\Phi_1)(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_2) + \text{H.c.}], \end{aligned} \tag{5}$$

so the parameters μ_1, μ_2, and μ_{12} are dimension of mass, $\lambda_i, i = 1, \ldots, 7$ are dimensionless, and the dimension of $\kappa_i, i = 1, \ldots, 13$ is of inverse mass squared. In the general case μ_1^2, μ_2^2, $\lambda_1, \ldots, \lambda_4$ and $\kappa_1, \ldots, \kappa_6$ are real, and all other parameters can be complex. In this section the mass basis for the general case of explicitly CP-violating potential [16,17] with nonzero imaginary parts of μ, λ, and κ parameters will be constructed. Transformations of the $SU(2)$ states $\eta_{1,2}, \chi_{1,2}, \omega_{1,2}^\pm$, Eq. (1), to the mass states $h, H, A, H^\pm, G^0, G^\pm$ can be performed using two orthogonal rotations

$$\begin{pmatrix} \eta_1 \\ \eta_2 \end{pmatrix} = \mathcal{O}_\alpha \begin{pmatrix} H \\ h \end{pmatrix}, \qquad \begin{pmatrix} \chi_1 \\ \chi_2 \end{pmatrix} = \mathcal{O}_\beta \begin{pmatrix} G^0 \\ A \end{pmatrix},$$
$$\begin{pmatrix} \omega_1^\pm \\ \omega_2^\pm \end{pmatrix} = \mathcal{O}_\beta \begin{pmatrix} G^\pm \\ h^\pm \end{pmatrix}, \tag{6}$$

where

$$\mathcal{O}_X = \begin{pmatrix} \cos X & -\sin X \\ \sin X & \cos X \end{pmatrix}, \qquad X = \alpha, \beta \tag{7}$$

(in the following we denote $\cos X = c_X, \sin X = s_X$, etc.) and the Higgs potential (2) up to I_6 terms takes the form

$$U = c_0 A + c_1 hA + c_2 HA + \frac{m_h^2}{2}h^2 + \frac{m_H^2}{2}H^2 + \frac{m_A^2}{2}A^2 + m_{H^\pm}^2 H^+H^- + I_3 + I_4 + I_5 + I_6. \tag{8}$$

Here $I_{3,4,5,6}$ denote the interaction terms of physical scalars and the coefficients c_i, $i = 0, 1, 2$, which are dependent on the imaginary parts of λ_i, κ_i,

$$\begin{aligned} c_1 &= v^2(-1/2\cdot \mathrm{Im}\lambda_5 c_{\alpha+\beta} + \mathrm{Im}\lambda_6 s_\alpha c_\beta - \mathrm{Im}\lambda_7 c_\alpha s_\beta) \\ &\quad + \frac{v^4}{4}\{-c_{\alpha+\beta}s_{2\beta}(3\mathrm{Im}\kappa_7 + \mathrm{Im}\kappa_{11} + \mathrm{Im}\kappa_{13}) \\ &\quad + 4(s_\alpha c_\beta^3 \mathrm{Im}\kappa_8 - c_\alpha s_\beta^3 \mathrm{Im}\kappa_{12}) \\ &\quad + 2[s_\beta^2(-3c_\alpha c_\beta + s_\alpha s_\beta)\mathrm{Im}\kappa_{10} \\ &\quad - c_\beta^2(c_\alpha c_\beta - 3s_\alpha s_\beta)\mathrm{Im}\kappa_9]\}, \end{aligned} \tag{9}$$

$$\begin{aligned} c_2 &= -v^2\left(1/2\cdot \mathrm{Im}\lambda_5 s_{\alpha+\beta} + \mathrm{Im}\lambda_6 c_\beta c_\alpha + \mathrm{Im}\lambda_7 s_\beta s_\alpha\right) - \frac{v^4}{2}[2\mathrm{Im}\kappa_8 c_\beta^3 c_\alpha + \mathrm{Im}\kappa_9 c_\beta^2(s_{\alpha+\beta} + 2c_\alpha s_\beta) \\ &\quad + \mathrm{Im}\kappa_{10}s_\beta^2(s_{\alpha+\beta} + 2c_\beta s_\alpha) + 2\mathrm{Im}\kappa_{12}s_\beta^3 s_\alpha + \frac{1}{2}(3\mathrm{Im}\kappa_7 + \mathrm{Im}\kappa_{11} + \mathrm{Im}\kappa_{13})s_{2\beta}s_{\alpha+\beta}] \end{aligned} \tag{10}$$

are equal to zero in the mass basis. In a local minimum where derivatives of the potential in the fields are zero, μ_1^2 and μ_2^2 can be expressed as

$$\begin{aligned} \mu_1^2 &= -\mathrm{Re}\mu_{12}^2 t_\beta + \frac{v^2}{4}[4\lambda_1 c_\beta^2 + 3\mathrm{Re}\lambda_6 s_{2\beta} + 2s_\beta^2(\lambda_{345} + \mathrm{Re}\lambda_7 t_\beta)] \\ &\quad + \frac{v^4}{4}\{3\kappa_1 c_\beta^4 + 5\mathrm{Re}\kappa_8 c_\beta^3 s_\beta + 3(\mathrm{Re}\kappa_7 + \mathrm{Re}\kappa_{11} + \mathrm{Re}\kappa_{13})c_\beta s_\beta^3 \\ &\quad + [\mathrm{Re}\kappa_9 + (\kappa_3 + \kappa_5)/2]s_{2\beta}^2 + (\kappa_4 + \kappa_6 + 2\mathrm{Re}\kappa_{10} + \mathrm{Re}\kappa_{12}t_\beta)s_\beta^4\}, \end{aligned} \tag{11}$$

$$\begin{aligned} \mu_2^2 &= -\mathrm{Re}\mu_{12}^2 \cot\beta + \frac{v^2}{4}[4\lambda_2 s_\beta^2 + 3\mathrm{Re}\lambda_7 s_{2\beta} + 2c_\beta^2(\lambda_{345} + \mathrm{Re}\lambda_6 \cot\beta)] \\ &\quad + \frac{v^4}{4}\{3\kappa_2 s_\beta^4 + 5\mathrm{Re}\kappa_{12}s_\beta^3 c_\beta + 3(\mathrm{Re}\kappa_7 + \mathrm{Re}\kappa_{11} + \mathrm{Re}\kappa_{13})s_\beta c_\beta^3 \\ &\quad + [\mathrm{Re}\kappa_{10} + (\kappa_4 + \kappa_6)/2]s_{2\beta}^2 + (\kappa_3 + \kappa_5 + 2\mathrm{Re}\kappa_9 + \mathrm{Re}\kappa_8 \cot\beta)c_\beta^4\}. \end{aligned} \tag{12}$$

The real part of μ_{12}^2 is fixed by zero eigenvalue of the mass matrix (which ensures a massless Goldstone boson state and defines the CP-odd scalar mass m_A^2)

$$\begin{aligned} \mathrm{Re}\mu_{12}^2 &= s_\beta c_\beta\left[m_A^2 + \frac{v^2}{2}(2\mathrm{Re}\lambda_5 + \mathrm{Re}\lambda_6\cot\beta + \mathrm{Re}\lambda_7\tan\beta)\right] \\ &\quad + v^4\{\mathrm{Re}\kappa_9 c_\beta^3 s_\beta + \mathrm{Re}\kappa_{10}c_\beta s_\beta^3 \\ &\quad + \frac{1}{4}[\mathrm{Re}\kappa_8 c_\beta^4 + \mathrm{Re}\kappa_{12}s_\beta^4 \\ &\quad + (9\mathrm{Re}\kappa_7 + \mathrm{Re}\kappa_{11} + \mathrm{Re}\kappa_{13})s_\beta^2 c_\beta^2]\}. \end{aligned} \tag{13}$$

The requirement $c_0 = 0$ in Eq. (8) fixes the imaginary part of μ_{12}^2,

$$\begin{aligned} \mathrm{Im}\mu_{12}^2 &= \frac{v^2}{2}(s_\beta c_\beta \mathrm{Im}\lambda_5 + c_\beta^2\mathrm{Im}\lambda_6 + s_\beta^2\mathrm{Im}\lambda_7) \\ &\quad + \frac{v^4}{4}[\mathrm{Im}\kappa_8 c_\beta^4 + 2\mathrm{Im}\kappa_9 c_\beta^3 s_\beta \\ &\quad + (3\mathrm{Im}\kappa_7 + \mathrm{Im}\kappa_{11} + \mathrm{Im}\kappa_{13})c_\beta^2 s_\beta^2 \\ &\quad + 2\mathrm{Im}\kappa_{10}c_\beta s_\beta^3 + \mathrm{Im}\kappa_{12}s_\beta^4]. \end{aligned} \tag{14}$$

Minimization conditions above must be performed for a generic two-doublet potential. In the following, the case of the MSSM potential will be analyzed. The one-loop resummed MSSM potential at the renormalization scale m_{top} using dimensional reduction and the $\overline{\mathrm{MS}}$ scheme can be written in the form

$$U_{\mathrm{eff}} = U^0 + \frac{3}{32\pi^2}\mathrm{tr}\mathcal{M}^4\left(\ln\frac{\mathcal{M}^2}{m_{\mathrm{top}}^2} - \frac{3}{2}\right), \tag{15}$$

where U^0 is a tree-level potential at the scale M_S,

$$U^0 = -\mu_1^2(\Phi_1^\dagger\Phi_1) - \mu_2^2(\Phi_2^\dagger\Phi_2) - [\mu_{12}^2(\Phi_1^\dagger\Phi_2) + \text{H.c.}] + \frac{g_1^2+g_2^2}{8}[(\Phi_1^\dagger\Phi_1)^2 + (\Phi_2^\dagger\Phi_2)^2] + \frac{g_2^2-g_1^2}{4}(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) - \frac{g_2^2}{4}(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1), \quad (16)$$

and $\mathcal{M}^2 = \mathcal{M}_M^2 + \mathcal{M}_\Gamma^2 + \mathcal{M}_\Lambda^2$ is the squark mass matrix squared (see Appendix). At the mass scale of quark superpartners, the mass matrix elements are

$$\mathcal{M}_{11}^2 = m_A^2 s_\beta^2 + m_Z^2 c_\beta^2, \qquad \mathcal{M}_{22}^2 = m_A^2 c_\beta^2 + m_Z^2 s_\beta^2,$$
$$\mathcal{M}_{12}^2 = -s_\beta c_\beta (m_A^2 + m_Z^2). \quad (17)$$

Radiative corrections to these tree-level expressions are parametrized using

$$\lambda_i(M) = \lambda_i^{\text{tree}}(M_S) - \Delta\lambda_i(M)/2, \qquad i = 1, 2,$$
$$\lambda_i(M) = \lambda_i^{\text{tree}}(M_S) - \Delta\lambda_i(M), \qquad i = 3, \ldots, 7, \quad (18)$$

where $\lambda_{1,2}^{\text{tree}} = \frac{g_1^2+g_2^2}{8}$, $\lambda_3^{\text{tree}} = \frac{g_2^2-g_1^2}{4}$, $\lambda_4^{\text{tree}} = -\frac{g_2^2}{2}$, $\lambda_{5,6,7}^{\text{tree}} = 0$, $\kappa_i^{\text{tree}} = 0$, $i = 1, \ldots, 13$, so corrections to the matrix elements of CP-even states mass matrix are

$$\Delta\mathcal{M}_{11}^2 = -v^2(\Delta\lambda_1 c_\beta^2 + \text{Re}\Delta\lambda_5 s_\beta^2 + \text{Re}\Delta\lambda_6 s_{2\beta}) + v^4[3\kappa_1 c_\beta^4 + 4\text{Re}\kappa_8 c_\beta^3 s_\beta + (\kappa_3 + \kappa_5 + 3\text{Re}\kappa_9)c_\beta^2 s_\beta^2 + (3\text{Re}\kappa_7 + \text{Re}\kappa_{11} + \text{Re}\kappa_{13})c_\beta s_\beta^3 + \text{Re}\kappa_{10} s_\beta^4], \quad (19)$$

$$\Delta\mathcal{M}_{22}^2 = -v^2(\Delta\lambda_2 s_\beta^2 + \text{Re}\Delta\lambda_5 c_\beta^2 + \text{Re}\Delta\lambda_7 s_{2\beta}) + v^4[\text{Re}\kappa_9 c_\beta^4 + (3\text{Re}\kappa_7 + \text{Re}\kappa_{11} + \text{Re}\kappa_{13})c_\beta^3 s_\beta + (\kappa_4 + \kappa_6 + 3\text{Re}\kappa_{10})c_\beta^2 s_\beta^2 + 4\text{Re}\kappa_{12} c_\beta s_\beta^3 + 3\kappa_2 s_\beta^4], \quad (20)$$

$$\Delta\mathcal{M}_{12}^2 = -v^2(\Delta\lambda_{34} s_\beta c_\beta + \text{Re}\Delta\lambda_6 c_\beta^2 + \text{Re}\Delta\lambda_7 s_\beta^2) + v^4[\text{Re}\kappa_8 c_\beta^4 + (\kappa_3 + \kappa_5 + \text{Re}\kappa_9)c_\beta^3 s_\beta + 2(\text{Re}\kappa_{11} + \text{Re}\kappa_{13})c_\beta^2 s_\beta^2 + (\kappa_4 + \kappa_6 + \text{Re}\kappa_{10})c_\beta s_\beta^3 + \text{Re}\kappa_{12} s_\beta^4]. \quad (21)$$

Then the masses of CP-even scalars can be expressed as

$$m_{H,h}^2 = \frac{1}{2}\left(m_A^2 + m_Z^2 + \Delta\mathcal{M}_{11}^2 + \Delta\mathcal{M}_{22}^2 \pm \sqrt{m_A^4 + m_Z^4 - 2m_A^2 m_Z^2 c_{4\beta} + C}\right), \quad (22)$$

where

$$C = 4\Delta\mathcal{M}_{12}^4 + (\Delta\mathcal{M}_{11}^2 - \Delta\mathcal{M}_{22}^2)^2 - 2(m_A^2 - m_Z^2)(\Delta\mathcal{M}_{11}^2 - \Delta\mathcal{M}_{22}^2)c_{2\beta} - 4(m_A^2 + m_Z^2)\Delta\mathcal{M}_{12}^2 s_{2\beta}, \quad (23)$$

and the mixing angle α is defined by

$$\tan 2\alpha = \frac{2\Delta\mathcal{M}_{12}^2 - (m_Z^2 + m_A^2)s_{2\beta}}{(m_Z^2 - m_A^2)c_{2\beta} + \Delta\mathcal{M}_{11}^2 - \Delta\mathcal{M}_{22}^2}. \quad (24)$$

The CP-odd scalar mass m_A can be expressed through m_h. Using Eq. (22) one can define m_A as an internal model parameter if the numerical value of the Higgs mass $m_h = 125$ GeV is fixed,

$$m_A^2 = \frac{m_h^2(C_1 - m_h^2) + m_Z^2(C_2 - C_3) - \Delta\mathcal{M}_{11}^2\Delta\mathcal{M}_{22}^2 + \Delta\mathcal{M}_{12}^4}{C_1 - C_2 - C_3 + m_Z^2 c_{2\beta}^2}, \quad (25)$$

where

$$C_1 = \Delta\mathcal{M}_{11}^2 + \Delta\mathcal{M}_{22}^2,$$
$$C_2 = m_h^2 - \Delta\mathcal{M}_{12}^2 s_{2\beta},$$
$$C_3 = \Delta\mathcal{M}_{11}^2 s_\beta^2 + \Delta\mathcal{M}_{22}^2 c_\beta^2.$$

The mass of the charged Higgs boson in the form

$$m_{H^\pm}^2 = m_W^2 + m_A^2 - \frac{v^2}{2}(\text{Re}\Delta\lambda_5 - \Delta\lambda_4) + \frac{v^4}{4}[c_\beta^2(2\text{Re}\kappa_9 - \kappa_5) + s_\beta^2(2\text{Re}\kappa_{10} - \kappa_6) - s_{2\beta}(\text{Re}\kappa_{11} - 3\text{Re}\kappa_7)] \quad (26)$$

can be obtained diagonalizing the corresponding mass matrix. Two important conditions that restrict implicitly the MSSM parameter space follow from Eq. (22):

$$m_A^4 + m_Z^4 - 2m_A^2 m_Z^2 c_{4\beta} + C \geq 0,$$
$$m_A^4 + m_Z^4 + \Delta\mathcal{M}_{11}^2 + \Delta\mathcal{M}_{22}^2 - 2m_h^2 \geq 0. \quad (27)$$

III. SYMBOLIC EXPRESSIONS FOR κ_i AND NUMERICAL RESULTS

The one-loop expressions for parameters κ_i in front of the dimension-six terms can be obtained decomposing the effective potential (15) in the inverse powers of M_S in the approximation of degenerate squark masses [12,18]. In the following we are using the notation $M_S = M_{\tilde{Q},\tilde{U},\tilde{D}}$ (see Appendix). Effective potential terms of the dimension-six in the decomposition are

$$U_{\mathrm{eff}}^{(6)} = \frac{3}{32M_S^2\pi^2}\left\{\frac{1}{3}\mathrm{tr}(\mathcal{M}_\Lambda^2)^3 - \frac{1}{2M_S^2}\mathrm{tr}[(\mathcal{M}_\Gamma^2)^2(\mathcal{M}_\Lambda^2)^2] + \frac{1}{6M_S^4}\mathrm{tr}[(\mathcal{M}_\Gamma^2)^4\mathcal{M}_\Lambda^2] - \frac{1}{60M_S^6}\mathrm{tr}(\mathcal{M}_\Gamma^2)^6\right\}. \tag{28}$$

Given the Lagrangian of the Higgs boson–squarks interaction (see Appendix), squark mass matrices $\mathcal{M}$ can be calculated and κ_i factors in front of the dimension-six terms can be derived. For example, factors κ_1 and κ_2 written in the form that uses powers of (μ/M_S) and (A/M_S) are

$$\begin{aligned}\kappa_1 = {}& \frac{h_D^6}{32M_S^2\pi^2}\left(2 - \frac{3|A_D|^2}{M_S^2} + \frac{|A_D|^4}{M_S^4} - \frac{|A_D|^6}{10M_S^6}\right) - h_D^4\frac{g_1^2+g_2^2}{128M_S^2\pi^2}\left(3 - 3\frac{|A_D|^2}{M_S^2} + \frac{|A_D|^4}{2M_S^4}\right)\\ &+ \frac{h_D^2}{512M_S^2\pi^2}\left(\frac{5}{3}g_1^4 + 2g_1^2g_2^2 + 3g_2^4\right)\left(1 - \frac{|A_D|^2}{2M_S^2}\right) - h_U^6\frac{|\mu|^6}{320M_S^8\pi^2} + h_U^4\frac{(g_1^2+g_2^2)|\mu|^4}{256M_S^6\pi^2}\\ &- h_U^2\frac{(17g_1^4 - 6g_1^2g_2^2 + 9g_2^4)|\mu|^2}{3072M_S^4\pi^2} + \frac{g_1^2}{1024M_S^2\pi^2}(g_1^4 - g_2^4),\end{aligned} \tag{29}$$

$$\begin{aligned}\kappa_2 = {}& -h_D^6\frac{|\mu|^6}{320M_S^8\pi^2} + h_D^4\frac{(g_1^2+g_2^2)|\mu|^4}{256M_S^6\pi^2} - h_D^2\frac{(5g_1^4 + 6g_1^2g_2^2 + 9g_2^4)|\mu|^2}{3072M_S^4\pi^2}\\ &- \frac{h_U^6}{32M_S^2\pi^2}\left(-2 + \frac{3|A_U|^2}{M_S^2} - \frac{|A_U|^4}{M_S^4} + \frac{|A_U|^6}{10M_S^6}\right) - h_U^4\frac{g_1^2+g_2^2}{128M_S^2\pi^2}\left(3 - 3\frac{|A_U|^2}{M_S^2} + \frac{|A_U|^4}{2M_S^4}\right)\\ &+ \frac{h_U^2}{3072M_S^2\pi^2}(17g_1^4 - 6g_1^2g_2^2 + 9g_2^4)\left(2 - \frac{|A_U|^2}{M_S^2}\right) - \frac{g_1^2}{1024M_S^2\pi^2}(g_1^4 - g_2^4).\end{aligned} \tag{30}$$

In a more compact notation κ_i, $i = 1, \ldots, 13$ can be rewritten using gauge coupling dependent factors G_i, $i = 1, \ldots, 4$, and parameter dependent factors A_j, B_k, and C_l,

$$\kappa_1 = h_D^6C_9^D - h_D^4G_4C_8^D + h_D^2G_2B_1^D + h_U^6A_1 + h_U^4G_4A_2 - h_U^2G_3A_3 + G_1, \tag{31}$$

$$\kappa_2 = h_D^6A_1 + h_D^4G_4A_2 - h_D^2G_2A_3 + h_U^6C_9^U - h_U^4G_4C_8^U + h_U^2G_3B_1^U - G_1, \tag{32}$$

$$\kappa_3 = h_D^6C_7^D + h_D^4G_4B_3^D - h_D^2G_2(2B_1^D + A_3) + h_U^6C_1^U - h_U^4G_4B_4^U|\mu|^2 + h_U^2G_3(B_1^U + 2A_3) - 3G_1, \tag{33}$$

$$\kappa_4 = h_D^6C_1^D - h_D^4G_4B_4^D|\mu|^2 + h_D^2G_2(B_1^D + 2A_3) + h_U^6C_7^U + h_U^4G_4B_3^U - h_U^2G_3(2B_1^U + A_3) + 3G_1, \tag{34}$$

$$\kappa_5 = h_D^6C_7^D + h_D^4G_4B_3^D - h_D^2G_2(2B_1^D + A_3) + h_U^6C_1^U - h_U^4G_4B_4^U|\mu|^2 + h_U^2G_3(B_1^U + 2A_3) - 3G_1, \tag{35}$$

$$\kappa_6 = h_D^6C_1^D - h_D^4G_4B_4^D|\mu|^2 + h_D^2G_2(B_1^D + 2A_3) + h_U^6C_7^U + h_U^4G_4B_3^U - h_U^2G_3(2B_1^U + A_3) + 3G_1, \tag{36}$$

$$\kappa_7 = \frac{\mu^3}{320M_S^8\pi^2}(A_D^3h_D^6 + A_U^3h_U^6), \tag{37}$$

$$\kappa_8 = h_D^6C_6^D + 2h_D^4G_4C_4^D + h_D^2G_2A_7^D + h_U^6A_2^U + h_U^4G_4A_5^U + h_U^2G_3A_7^U, \tag{38}$$

$$\kappa_9 = h_D^6C_2^D - h_D^4G_4A_6^D + h_U^6A_4^U + h_U^4G_4A_6^U, \tag{39}$$

$$\kappa_{10} = h_D^6A_4^D + h_D^4G_4A_6^D + h_U^6C_2^U - h_U^4G_4A_6^U, \tag{40}$$

$$\kappa_{11} = h_D^6C_3^D + h_D^4G_4C_5^D - 2h_D^2G_2A_7^D + h_U^6C_3^U + h_U^4G_4C_5^U - 2h_U^2G_3A_7^U, \tag{41}$$

$$\kappa_{12} = h_D^6A_2^D + h_D^4G_4A_5^D + h_D^2G_2A_7^D + h_U^6C_6^U + 2h_U^4G_4C_4^U + h_U^2G_3A_7^U, \tag{42}$$

$$\kappa_{13} = h_D^6C_3^D + h_D^4G_4C_5^D - 2h_D^2G_2A_7^D + h_U^6C_3^U + h_U^4G_4C_5^U - 2h_U^2G_3A_7^U, \tag{43}$$

where $(X = U,\ D)$

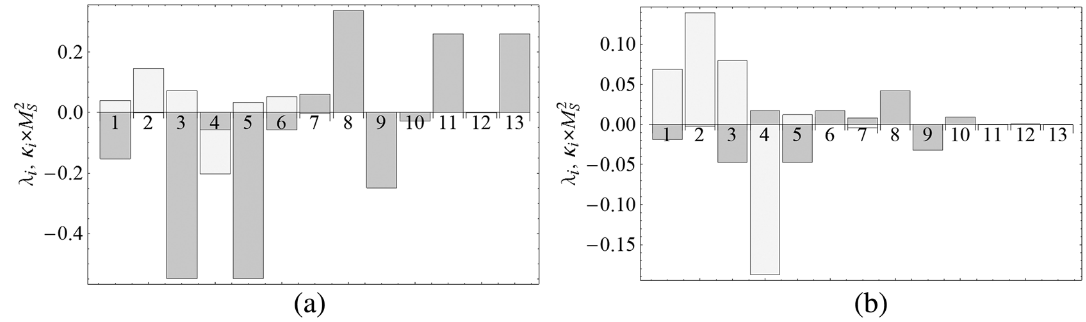

FIG. 1. The dimensionless parameters λ_i, $i = 1, \ldots, 7$ (light gray) calculated using the analytical results of [16] and $\kappa_j \cdot M_S^2$, $j = 1, \ldots, 13$ (dark gray) calculated at the squark mass scale (a) $M_S = 5$ TeV and (b) $M_S = 7$ TeV for $A_t = A_b = 10$ TeV, $\mu = 14$ TeV, and $\tan\beta = 5$.

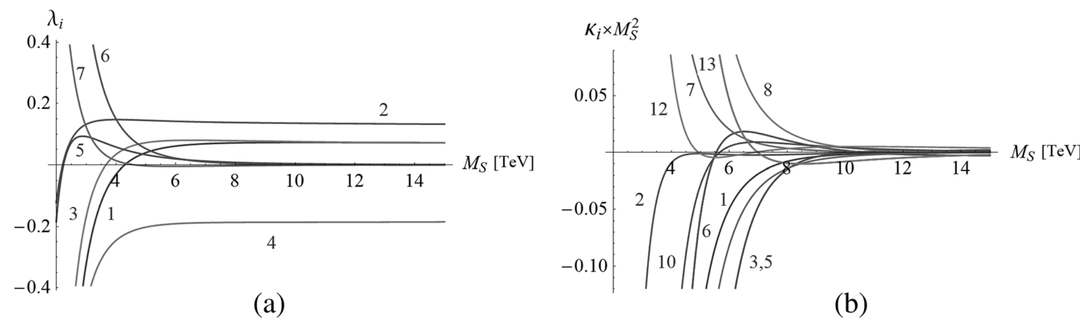

FIG. 2. The dimensionless parameters (a) λ_i and (b) $\kappa_i \cdot M_S^2$ as a function of M_S for $A_t = A_b = 10$ TeV, $\mu = 14$ TeV, $\tan\beta = 5$. λ_i are evaluated using analytical formulas from [16], where the contribution of nonleading D terms is accounted for.

$$G_1 = \frac{1}{M_S^2}\frac{g_1^2(g_1^4 - g_2^4)}{1024\pi^2}, \qquad G_2 = \frac{5g_1^4 + 6g_1^2g_2^2 + 9g_2^4}{3072\pi^2}, \qquad G_3 = \frac{17g_1^4 - 6g_1^2g_2^2 + 9g_2^4}{3072\pi^2}, \qquad G_4 = \frac{g_1^2 + g_2^2}{256\pi^2}, \tag{44}$$

$$\begin{aligned} A_1 &= -\frac{|\mu|^6}{320M_S^8\pi^2}, \qquad A_2 = \frac{|\mu|^4}{M_S^6}, \qquad A_3 = \frac{|\mu|^2}{M_S^4}, \qquad A_2^X = \frac{3A_X\mu|\mu|^4}{320M_S^8\pi^2}, \\ A_4^X &= -\frac{3A_X^2\mu^2|\mu|^2}{320M_S^8\pi^2}, \qquad A_5^X = -\frac{2A_X\mu|\mu|^2}{M_S^6}, \qquad A_6^X = \frac{A_X^2\mu^2}{M_S^6}, \qquad A_7^X = \frac{\mu A_X}{M_S^4}, \end{aligned} \tag{45}$$

$$B_1^X = -\frac{|A_X|^2}{M_S^4} + \frac{2}{M_S^2}, \qquad B_2^X = -\frac{4|A_X|^2}{M_S^6} + \frac{6}{M_S^4}, \qquad B_3^X = C_8^X + |\mu|^2 B_2^X, \qquad B_4^X = \frac{|\mu|^2}{M_S^6} + B_2^X, \tag{46}$$

$$\begin{aligned} C_1^X &= \frac{|\mu|^4}{320\pi^2}\left(-\frac{9|A_X|^2}{M_S^8} + \frac{10}{M_S^6}\right), \qquad C_2^X = \frac{A_X^2\mu^2}{320\pi^2}\left(-\frac{3|A_X|^2}{M_S^8} + \frac{10}{M_S^6}\right), \\ C_3^X &= \frac{A_X\mu|\mu|^2}{320\pi^2}\left(\frac{9|A_X|^2}{M_S^8} - \frac{20}{M_S^6}\right), \qquad C_4^X = A_X\mu\left(\frac{|A_X|^2}{M_S^6} - \frac{3}{M_S^4}\right), \\ C_5^X &= -2A_X\mu\left(\frac{|A_X|^2 - |\mu|^2}{M_S^6} - \frac{3}{M_S^4}\right), \qquad C_6^X = \frac{A_X\mu}{320\pi^2}\left(\frac{3|A_X|^4}{M_S^8} - \frac{20|A_X|^2}{M_S^6} + \frac{30}{M_S^4}\right), \\ C_7^X &= -\frac{|\mu|^2}{320\pi^2}\left(\frac{9|A_X|^4}{M_S^8} - \frac{40|A_X|^2}{M_S^6} + \frac{30}{M_S^4}\right), \qquad C_8^X = \frac{|A_X|^4}{M_S^6} - \frac{6|A_X|^2}{M_S^4} + \frac{6}{M_S^2}, \\ C_9^X &= -\frac{1}{320\pi^2}\left(\frac{|A_X|^6}{M_S^8} - \frac{10|A_X^4|}{M_S^6} + \frac{30|A_X|^2}{M_S^4} - \frac{20}{M_S^2}\right). \end{aligned} \tag{47}$$

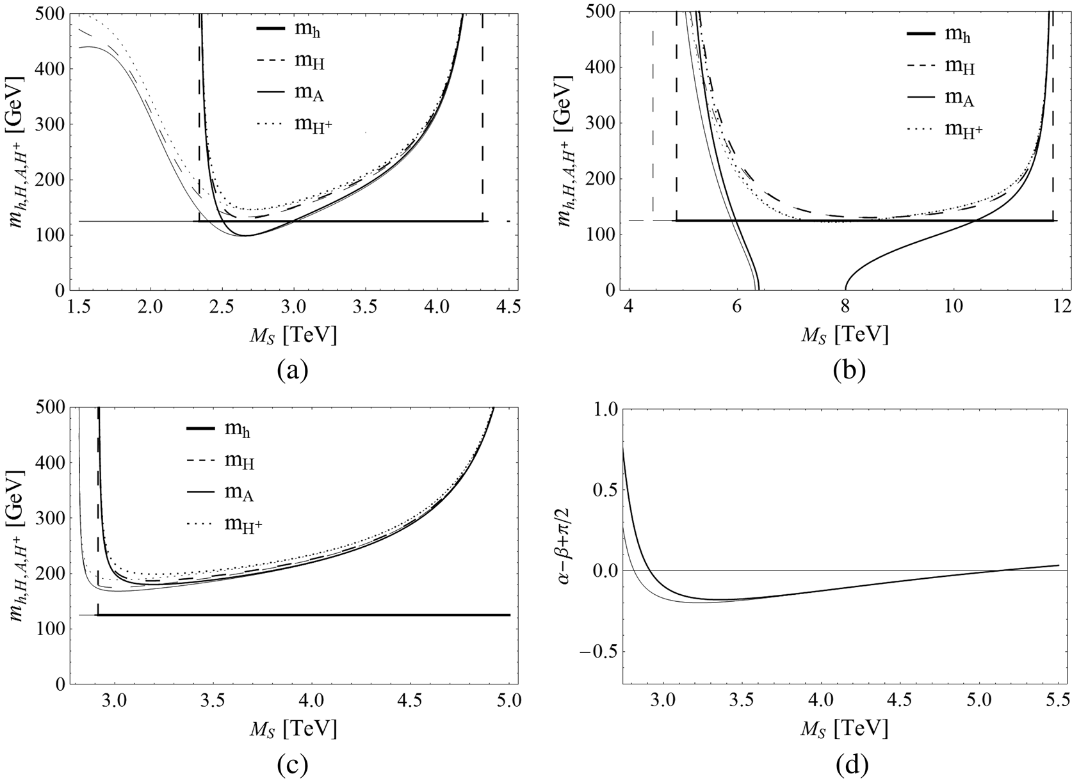

FIG. 3. Higgs boson masses and mixing angles combination as functions of the squark mass scale M_S. Thin lines correspond to the effective potential $U^{(4)}$ including terms with the maximal dimension four in the fields, and thick lines are the results for masses calculated with the effective potential $U^{(6)}$ including the dimension-six operators. The MSSM parameter sets: (a) $\tan\beta = 4$, $A = 10$ TeV, $\mu = 8$ TeV; (b) $\tan\beta = 8$, $A = 25$ TeV, $\mu = 30$ TeV; (c), (d) $\tan\beta = 5$, $A = 10$ TeV, and $\mu = 5$ TeV. The discontinuity in (c) at M_S of about 3 TeV corresponds to the zero denominator of Eq. (25).

Meaningful numerical results following from the effective potential expansions in the inverse powers of M_S are using the assumption of small mass splitting among the squark mass eigenstates (or simultaneous decoupling of squark fields). In the literature it is usually considered that the expansion is valid if $(m^2_{\tilde{t}_1} - m^2_{\tilde{t}_2})/(m^2_{\tilde{t}_1} + m^2_{\tilde{t}_2}) < 0.5$ where $m_{\tilde{t}_{1,2}}$ are the stop masses. Then M_S^2 can be defined as the average $(m^2_{\tilde{t}_1} + m^2_{\tilde{t}_2})/2$.[1] The contribution of dimension-six operators is small in the phase with softly broken symmetry if at least $2|m_{\rm top}A_t| < M_S^2$ and $2|m_{\rm top}\mu| < M_S^2$ [12]. However, the dimension-six terms may play an important role in the A, μ parameter range of about/of the order of 10^1 TeV and moderate M_S. For example, values of κ_i evaluated for $A = 10$ TeV, $\mu = 14$ TeV, $\tan\beta = 5$ are shown in Fig. 1, where the dimensionless couplings $\kappa_i \cdot M_S^2$ are depicted for M_S values of 5 and 7 TeV. The behavior of λ_i and $\kappa_i \cdot M_S^2$ as a function of M_S at the multi-TeV energy scale is shown in Fig. 2. One can see that significant values of $\kappa_i \cdot M_S^2$ are observed in the M_S range less than 8 TeV.

The Higgs boson masses $m_{H,A,H^\pm}$ evaluated for two ($\tan\beta$, A, μ) parameter sets at fixed values of the lightest CP-even state mass $m_h = 125$ GeV and large X_t mixing parameter of the order of 10 TeV are shown in Fig. 3 as a function of the squark mass scale M_S. The CP-odd scalar mass m_A is calculated using Eq. (25), where m_h is an input parameter with fixed value. A pole of $m_A^2(M_S)$ may take place when the denominator in Eq. (25) is zero. In the unphysical region of M_S, for example, to the left of the pole in Fig. 3(c), the restrictions imposed by Eq. (27) are not respected. The contribution of the dimension-six terms $U^{(6)}$ to masses of scalars is very small in comparison with the dimension-four terms $U^{(4)}$ for moderate M_S [$M_S \geq 3$ TeV, Fig. 3(a), and $M_S \geq 7$ TeV, Fig. 3(b)], but for smaller M_S corrections

[1]Besides the above-mentioned approach developed in [13], recent direct comparison of results for the one-loop MSSM amplitudes ggh and $\gamma\gamma h$ obtained by means of the diagrammatic calculation and the covariant derivative expansion method [19] for the case of either degenerate or nondegenerate stop mass spectrum can be found in [20]. For large $\tan\beta$ the approximation of (almost) degenerate stop masses is not satisfactory at $m_{\tilde{t}} < 0.5$ TeV and large X_t mixing parameter values of a few TeV; however, $m_h = 125$ GeV is mostly available.

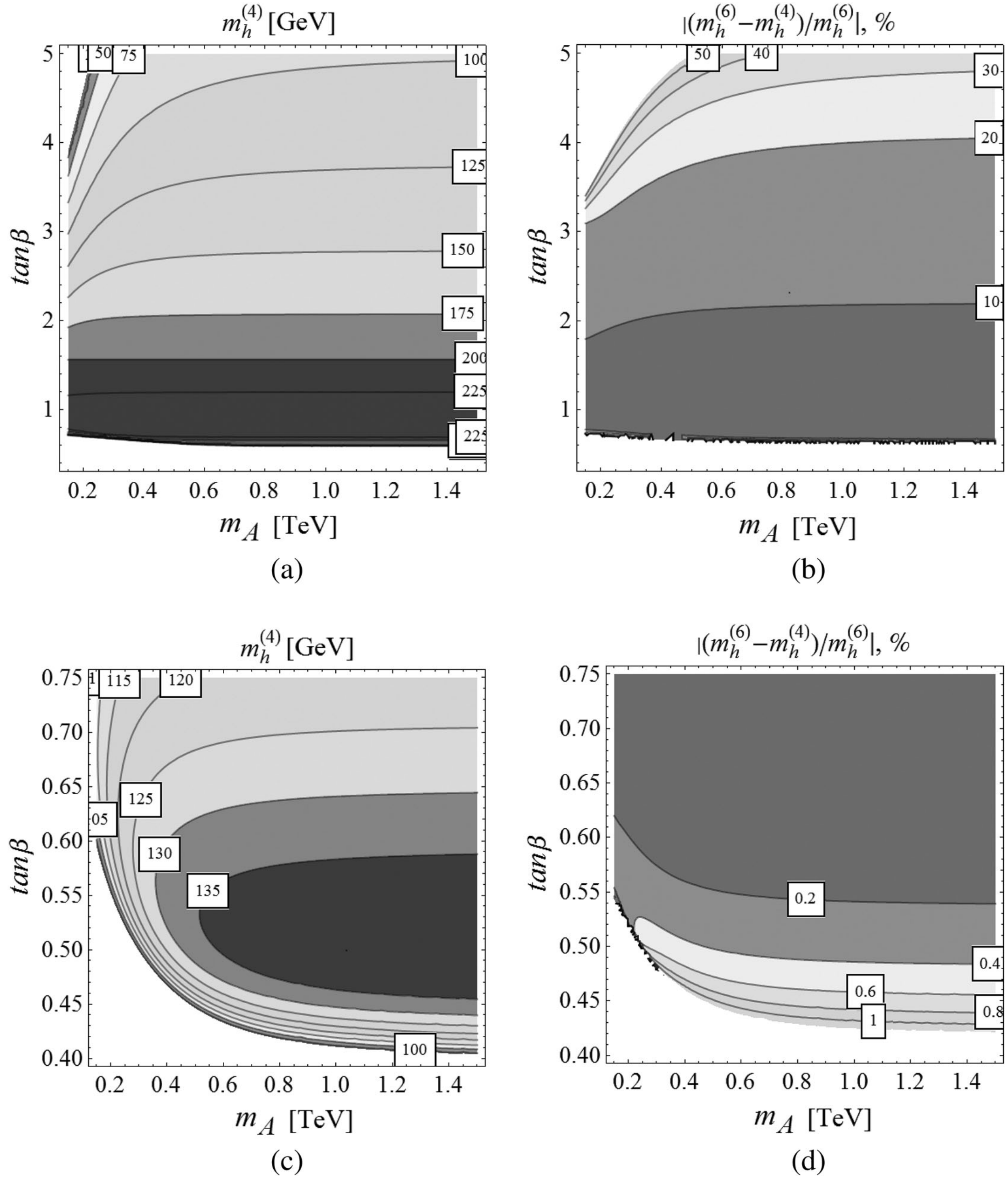

FIG. 4. Left panel: Contours for the Higgs boson mass $m_h^{(4)}$ calculated with the dimension-four potential terms; right panel: the relative difference in percent between $m_h^{(6)}$ and $m_h^{(4)}$ masses; the parameter set $A = 10$ TeV, $\mu = 8.3$ TeV, $M_S = 2$ TeV (a), (b) and $M_S = 5$ TeV (c), (d).

are very important. In Fig. 3(a) the physical region of $m_h^2 > 0$ indicated by vertical lines narrows to 2.3 TeV (lower bound). In Fig. 3(b), the CP-odd scalar mass squared is not positively defined for the M_S range from 6.3 to 8 TeV. At moderate $\tan\beta \approx 10$ positively defined masses squared of H, A, and $H^\pm$ consistent with the input $m_h = 125$ GeV are not possible for M_S greater than 12 TeV. Note that a nonstandard mass spectrum with an extremely light pseudoscalar is available in this case. At fixed $m_h = 125$ GeV the CP-odd state A can be as light as 25–30 GeV with H and $H^\pm$ states in the decoupling regime or with masses of the order of the electroweak scale. For example, Higgs masses for the Fig. 3(b) parameter set and $M_S \simeq 6.3$ TeV are $m_h = 125$ GeV, $m_H = 190$ GeV, $m_A = 27$ GeV, $m_{H^\pm} = 170$ GeV. The alignment limit [21] when $\alpha \approx \beta - \pi/2$ takes place for set 3(b) in the vicinity of $M_S = 5.5$ TeV; it is possible for $A, H, H^\pm$ in the decoupling regime only. The regime of alignment without decoupling without small m_A is available if $\tan\beta = 5$, $A = 10$ TeV, and $\mu = 5$ TeV. For this parameter set [see Fig. 3(c)],

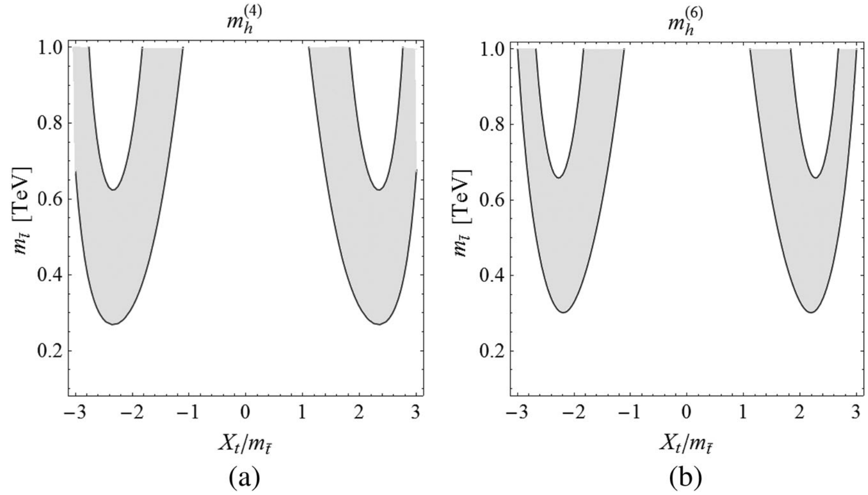

FIG. 5. Domains of the Higgs boson mass $m_h = 125 \pm 3$ GeV for $m_A = 300$ GeV and the Higgs superfield mass parameter μ equal to zero, $\tan\beta = 20$, calculated with (a) the dimension-four operators and (b) the dimension-six operators.

when curves are more stable with respect to corrections, there are two alignment limits. In Fig. 3(d) the first alignment limit takes place at $M_S = 2.98$ TeV without decoupling, and the second limit at 5.1 TeV demonstrates decoupling of H, A, and $H^{\pm}$ states. Figure 4 illustrates an increasing role of corrections from the $U^{(6)}$ terms to m_h in the case of "low $\tan\beta$" scenarios [4], which are found to be about 1% at $M_S = 5$ TeV and A, μ more than 10 TeV and about 20% for the lower superpartner mass scale $M_S = 2$ TeV and A, μ less than 10 TeV. In Fig. 5 the condition $m_h = 125 \pm 3$ GeV is translated to the mixing parameter–quark superpartner mass plane $(X_t/m_{\tilde t}, m_{\tilde t})$, demonstrating sensitivity of the contours in the regime $\mu = 0$ (see also [20], where similar contours are reconstructed using the diagrammatic calculation [22]). Increasing the μ parameter of a few hundreds of GeV changes strongly these exclusion contours, leaving only a small acceptable domain in the left upper corner of the plot.

IV. SUMMARY

In the absence of direct evidence motivating extensions of the Standard Model-like Higgs sector, the effective field theory (EFT) approach is a convenient framework to describe possible new physics either in a model-independent or in a model-dependent way. In both cases, the MSSM Lagrangian is extended by higher-dimensional operators that are suppressed by the mass scale of new physics. In the model-dependent case of the MSSM when the resummed effective potential is expanded up to dimension-six operators induced by the soft supersymmetry breaking terms, we calculated symbolically corrections to the effective sextic couplings and used them to determine the post-Higgs discovery mass spectrum of the heavy MSSM Higgs bosons. An improved precision can be reached using such procedure, especially at the low EFT cutoff scale. Corrections to the mass spectrum depend strongly on the domain in the MSSM parameter space and are defined mainly by the quark superpartner mass scale and mixing in the sector of soft SUSY-breaking terms. Even at the moderate mixing parameter values significant contributions to the heavy scalar mass spectrum of the order of 10%–20% induced by the dimension-six operators are found at the squark mass scale $M_S \sim 2$–3 TeV. Thus, for moderately heavy supersymmetry additional corrections induced by higher-order terms in the expansion of the effective potential should be taken into account. One can observe that in a number of cases the restrictions on the MSSM parameter space are not so much a consequence of the condition $m_h = 125$ GeV as the presence of the mass basis for the five Higgs bosons, where mass hierarchy is acceptable from the experimental point of view and there are no tachyonic states.

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APPENDIX: HIGGS BOSON—SQUARKS INTERACTIONS

The most general scalar potential, including the Higgs boson and one generation of squarks, can be written as [18,23]

$$\mathcal{V}^0 = \mathcal{V}_M + \mathcal{V}_\Gamma + \mathcal{V}_\Lambda + \mathcal{V}_{\tilde Q}, \qquad \text{(A1)}$$

where $\mathcal{V}_M$ contains mass squark terms, $\mathcal{V}_\Gamma$–F terms, $\mathcal{V}_\Lambda$–D terms of Higgs-squark interactions and $\mathcal{V}_{\tilde Q}$–quartic squark interaction terms,

$$\mathcal{V}_M = -\mu_{ij}^2 \Phi_i^\dagger \Phi_j + M_{\tilde{Q}}^2 (\tilde{Q}^\dagger \tilde{Q}) + M_{\tilde{U}}^2 (\tilde{U}^* \tilde{U}) + M_{\tilde{D}}^2 (\tilde{D}^* \tilde{D}), \tag{A2}$$

$$\mathcal{V}_\Gamma = \Gamma_i^D (\Phi_i^\dagger \tilde{Q}) \tilde{D} + \Gamma_i^U (i\Phi_i^T \sigma_2 \tilde{Q}) \tilde{U} + \text{H.c.}, \tag{A3}$$

$$\begin{aligned}\mathcal{V}_\Lambda = {} & \Lambda_{ik}^{jl} (\Phi_i^\dagger \Phi_j)(\Phi_k^\dagger \Phi_l) \\ & + (\Phi_i^\dagger \Phi_j)[\Lambda_{ij}^Q (\tilde{Q}^\dagger \tilde{Q}) + \Lambda_{ij}^U (\tilde{U}^* \tilde{U}) + \Lambda_{ij}^D (\tilde{D}^* \tilde{D})] \\ & + \overline{\Lambda}_{ij}^Q (\Phi_i^\dagger \tilde{Q})(\tilde{Q}^\dagger \Phi_j) \\ & + \frac{1}{2}[\Lambda \epsilon_{ij} (i\Phi_i^T \sigma_2 \Phi_j) \tilde{D}^* \tilde{U} + \text{H.c.}],\end{aligned} \tag{A4}$$

and Γ, Λ are determined by the tree-level SUSY relations,

$$\Lambda^Q = \mathrm{diag}\left[\frac{1}{4}(g_2^2 - g_1^2 Y_Q), h_U^2 - \frac{1}{4}(g_2^2 - g_1^2 Y_Q)\right], \tag{A5}$$

$$\overline{\Lambda}^Q = \mathrm{diag}\left(h_D^2 - \frac{1}{2}g_2^2, \frac{1}{2}g_2^2 - h_U^2\right), \tag{A6}$$

$$\Lambda^U = \mathrm{diag}\left(-\frac{1}{4}g_1^2 Y_U, h_U^2 + \frac{1}{4}g_1^2 Y_U\right), \tag{A7}$$

$$\Lambda^D = \mathrm{diag}\left(h_D^2 - \frac{1}{4}g_1^2 Y_D, \frac{1}{4}g_1^2 Y_D\right), \tag{A8}$$

$$\Lambda = -h_U h_D, \tag{A9}$$

$$\Gamma_{1,2}^U = h_U(-\mu, A_U), \qquad \Gamma_{1,2}^D = h_D(A_D, -\mu), \tag{A10}$$

$g_{1,2}$ are couplings of $SU(2)_L \times U(1)_Y$, $Y_{Q,U,D} = \{\frac{1}{3}(-1), \frac{2}{3}(2), -\frac{4}{3}\}$–squark (slepton) hypercharges, $h_U = \frac{g_2 m_U}{\sqrt{2} m_W s_\beta}$, $h_D = \frac{g_2 m_D}{\sqrt{2} m_W c_\beta}$–Yukawa couplings, $A_{U,D}$–trilinear couplings, μ–Higgs superfield mass parameter.

The squark mass matrix is obtained by taking derivatives

$$\mathcal{M}_{a,b}^2 = \frac{\partial^2 \mathcal{V}^0}{\partial \Psi_a \partial \Psi_b^*}, \tag{A11}$$

where $\Psi = (\tilde{Q}, \tilde{U}^*, \tilde{D}^*)$.

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