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Higgs boson couplings to quarks with supersymmetric CP and flavor violations

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Abstract

In minimal supersymmetric model (SUSY) with a light Higgs sector, explicit CP violation and most general flavor mixings in the sfermion sector, integration of the superpartners out of the spectrum induces potentially large contributions to the Yukawa couplings of light quarks via those of the heavier ones. These corrections can be sizeable even for moderate values of $\tan\beta$, and remain nonvanishing even if all superpartners decouple. When the SUSY breaking scale is close to the electroweak scale, the Higgs exchange effects can compete with the gauge boson and box diagram contributions to rare processes, and their partial cancellations can lead to relaxation of the existing bounds on flavor violation sources. In this case there exist sizeable enhancements in flavor-changing Higgs decays. When the superpartners completely decouple, however, the Higgs mediation becomes the dominant SUSY contribution to rare processes the saturation of which, without a strong suppression of the flavor mixings, prefers large $\tan\beta$ and certain ranges for the CP-odd phases. The decay rate of the lightest Higgs into light down quarks becomes comparable with that into the bottom quark. Moreover, the Higgs decay into the up quark is significantly enhanced. There are observable implications for rare processes, atomic electric dipole moments, and collider searches for Higgs bosons.

The standard model of electroweak interactions (SM) has been extremely successful in explaining all the available data. The least understood aspects of the model concern the breaking of *gauge*, *CP* and *flavor* symmetries. Indeed, the Higgs boson mass and various parameters in the Yukawa matrices are left to experimental determination. Though the indications at LEP for a light Higgs boson of mass ~ 115 GeV are encouraging a full construction of the symmetry-breaking sector, including possibly its CP properties, is to wait for the upgraded Tevatron or LHC. On the other hand, existing as well as future data to come from the experiments on kaon, beauty and charmed hadrons will determine the structure of CP and flavor violations.

The scalar sector, which is responsible for breaking the gauge symmetry, is quadratically sensitive to the UV cut-off and hence the model must be embedded into a UV-safe extension beyond the TeV scale. Supersymmetry (SUSY) is the only weak-scale extension which stabilizes the Higgs sector against quadratic divergences and unifies the gauge couplings at high energies in agreement with the electroweak precision data. Quite generically, the SUSY models bring about novel sources for CP and flavor violations through the soft breaking masses. The

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main reason for SUSY flavor violation is that the fermions and sfermions are misaligned in the flavor space, and even if the flavor violation in the fermion sector is reduced to that of the CKM matrix the sfermion sector maintains its non-CKM structure.

The LR and RL $=$ LR† blocks of the sfermion mass-squared matrices are generated after the electroweak breaking with the maximal size $\mathcal{O}(m_t M_{\rm SUSY})$. The nontrivial flavor structures of these blocks are dictated by the Yukawa couplings $\mathbf{Y}_{u,d}$ and by the trilinear coupling matrices $\mathbf{Y}^A_{u,d}$ with

$$\left(\mathbf{Y}^A_u\right)_{ij} = (\mathbf{Y}_u)_{ij}(A_u)_{ij} \quad \text{and} \quad \left(\mathbf{Y}^A_d\right)_{ij} = (\mathbf{Y}_d)_{ij}(A_d)_{ij}, \tag{1}$$

where $A_{u,d}$ are not necessarily unitary so that even their diagonal entries contribute to CP-violating observables. The LL and RR blocks are insensitive to electroweak breaking, and their texture is determined by the SUSY breaking pattern. In minimal SUGRA and its nonuniversal variants with CP violation, for instance, size and structure of flavor and CP violation in LL and RR blocks are dictated by the CKM matrix [1]. On the other hand, in SUSY GUTs with Yukawa unification, e.g., SO(10), implementation of the see-saw mechanism for neutrino masses implies sizeable flavor violation in the RR block, given the large mixings observed in atmospheric neutrino data [2]. Independent of specific realizations, the squark mass-squared matrices can be parameterized as

$$\left(M^2_D\right)_{LL} = \begin{pmatrix} M^2_{\tilde{d}_L} & M^2_{\tilde{d}_L\tilde{s}_L} & M^2_{\tilde{d}_L\tilde{b}_L} \\ M^2_{\tilde{s}_L\tilde{d}_L} & M^2_{\tilde{s}_L} & M^2_{\tilde{s}_L\tilde{b}_L} \\ M^2_{\tilde{b}_L\tilde{d}_L} & M^2_{\tilde{b}_L\tilde{s}_L} & M^2_{\tilde{b}_L} \end{pmatrix}, \qquad \left(M^2_D\right)_{RR} = \begin{pmatrix} M^2_{\tilde{d}_R} & M^2_{\tilde{d}_R\tilde{s}_R} & M^2_{\tilde{d}_R\tilde{b}_R} \\ M^2_{\tilde{s}_R\tilde{d}_R} & M^2_{\tilde{s}_R} & M^2_{\tilde{s}_R\tilde{b}_R} \\ M^2_{\tilde{b}_R\tilde{d}_R} & M^2_{\tilde{b}_R\tilde{s}_R} & M^2_{\tilde{b}_R} \end{pmatrix} \tag{2}$$

in the bases $\{\tilde{d}_L, \tilde{s}_L, \tilde{b}_L\}$ and $\{\tilde{d}_R, \tilde{s}_R, \tilde{b}_R\}$, respectively. The same structure repeats for the up sector. The hermiticity of the mass matrices, $(M^2_D)_{LL,RR} = (M^2_D)^\dagger_{LL,RR}$, allows CP violation only in the off-diagonal entries.

In comparison to the SM amplitudes, the virtual effects of sparticles on the rare processes scale as $M_W/M_{\rm SUSY}$ to appropriate power due to either their derivative coupling to the vector bosons or the sensitivity of the particular amplitude to the electroweak breaking [3,4]. Similarly, the hadronic and leptonic dipole moments scale as (fermion mass)$/M^2_{\rm SUSY}$. In this sense, various bounds on SUSY flavor and CP violation sources from the current experimental data depend on how close $M_{\rm SUSY}$ is the electroweak scale. Looking from a different channel, the FCNC couplings of Z boson to fermions scale as $M^2_Z/M^2_{\rm SUSY}$ for Z boson decays [5], and SUSY effects become transparent only at collider energies $E \sim M_{\rm SUSY}$. This decoupling property of the SUSY effects does not hold for interactions of Higgs bosons with fermions as their couplings to sfermions are dictated by the soft-breaking sector. Consequently, gauge and Higgs bosons, considering their decays and productions as well as the FCNC processes they mediate, possess essential differences concerning their sensitivity to the SUSY breaking scale. Indeed, the contributions of the sparticles, even if they are too heavy to be produced directly at near-future colliders, to gauge (Higgs) boson couplings to fermions are (are not) suppressed by $1/M_{\rm SUSY}$. This nondecoupling property of the Higgs bosons persists unless the Higgs sector itself enters the decoupling regime in which case the SM results are recovered [6]. Therefore, when the SUSY Higgs sector is stabilized at the weak scale the Higgs boson interactions with the standard matter provides a direct access to SUSY even if it can be in the decoupling regime. The Higgs-mediated FCNC becomes sizeable when the vacuum expectation values (VEV) of the Higgs fields are hierarchically split. This regime of the parameter space is motivated by LEP constraints on the SUSY parameter space and by the Yukawa-unified models like SO(10) [7–9]. Depending on what sparticles are contained in the light spectrum the weak-scale effective theory can vary from a two-doublet model to a full SUSY model as two extremes. The EDM and FCNC constraints on Higgs mediation can be strong for the former [8–10] whereas they can be milder for the latter [11,12].

The purpose of this work is to compute the couplings of Higgs bosons to quarks in the presence of SUSY CP and flavor violation effects within the minimal SUSY model. It will be shown that there are parametrically sizeable corrections to light quark Yukawas which imply novel properties: (i) the present constraints from non-Higgs contributions to FCNC processes [3,4] can be modified, (ii) the EDMs can probe CP violation from both flavor-blind and flavor-sensitive SUSY phases, (iii) the flavor-violating decay rates of the Higgs bosons can be comparable with the flavor-conserving ones, and (iv) the Higgs bosons can turn out to be totally blind to all quarks but the charm and the top. These phenomena have observable signatures for experiments at meson factories as well as Higgs searches at colliders.

In general, in models with two or more Higgs doublets suppression of the tree level FCNC is accomplished by imposing certain symmetries. In minimal SUSY, it is a U(1) symmetry under which all fields are neutral except for $\mathbf{d}_R$ and H_d which have identical charges. This implies that the Higgs doublet H_u (H_d) couples only to up (down) type quarks. However, the symmetry under concern is broken at the loop level due to the soft SUSY-breaking masses [7,8]. Thus, the effective Lagrangian describing the Higgs–quark interactions below $M_{\rm SUSY}$ may be written as

$$-\mathcal{L} = \bar{\mathbf{d}}_R\big[\mathbf{Y}_d - \gamma^d\big]H_d^0\mathbf{d}_L + \bar{\mathbf{d}}_R\Gamma^d H_u^{0\star}\mathbf{d}_L + \bar{\mathbf{u}}_R\big[\mathbf{Y}_u + \gamma^u\big]H_u^0\mathbf{u}_L - \bar{\mathbf{u}}_R\Gamma^u H_d^{0\star}\mathbf{u}_L + \text{h.c.}, \tag{3}$$

where, at tree level, flavor and CP violations are entirely determined by the Yukawa matrices $\mathbf{Y}_d$ and $\mathbf{Y}_u$ whose simultaneous diagonalization leads to the CKM matrix as the only observable effect. Therefore, without loss of generality, one can choose an appropriate basis for $\mathbf{Y}_{d,u}$ such as the down quark diagonal one

$$\mathbf{Y}_d = \begin{pmatrix} h_d & 0 & 0 \\ 0 & h_s & 0 \\ 0 & 0 & h_b \end{pmatrix}, \qquad \mathbf{Y}_u = \begin{pmatrix} h_u & 0 & 0 \\ 0 & h_c & 0 \\ 0 & 0 & h_t \end{pmatrix} \cdot V^0 \tag{4}$$

where h_i and V^0 are tree-level Yukawa couplings and the CKM matrix, respectively.

The nonholomorphic Yukawa structures $\gamma^{u,d}$ and $\Gamma^{u,d}$ in (3) result from integrating out the heavy degrees of freedom which may include the entire sparticle spectrum or part of it. The dominant contributions to these SUSY threshold effects can be gathered by employing the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ symmetric limit and neglecting their gauge couplings (cf. [9] for a discussion the electroweak breaking effects). Then the electroweak breaking occurs after integrating out the sparticles. In this limit, the LR and RL blocks of the sfermion mass matrices vanish so do the self-energy corrections on the quark lines. Hence, $\gamma^{u,d}$ and $\Gamma^{u,d}$ are generated by the vertex diagrams meditated by gluino–squark and Higgsino–squark loops. Using the Yukawa bases (4) in the trilinear couplings (1) and relabelling the quarks and squarks as $\{d, s, b\} \equiv \{d^1, d^2, d^3\}$ and $\{u, c, t\} \equiv \{u^1, u^2, u^3\}$, one finds

$$\begin{aligned}
\gamma_{ii}^d = {} & \frac{2\alpha_s}{3\pi}\big(\mathbf{Y}_d^A\big)_{ii} M_g^\star I_3\Big(M^2_{\tilde{d}_L^i}, M^2_{\tilde{d}_R^i}, |M_g|^2\Big) \\
& + \frac{2\alpha_s}{3\pi}\sum_{j=1}^{3}\big(\mathbf{Y}_d^A\big)_{jj} M_g^\star M_{\tilde{D}}^4 I_5\Big(M^2_{\tilde{d}_L^i}, M^2_{\tilde{d}_L^j}, M^2_{\tilde{d}_R^j}, M^2_{\tilde{d}_R^i}, |M_g|^2\Big)\big(\delta_{ij}^d\big)_{RR}\big(\delta_{ji}^d\big)_{LL} \\
& + \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\sum_{j=1}^{3}\big(\mathbf{Y}_u^\dagger\big)_{ij}(\mathbf{Y}_u)_{ji}|\mu|^2 I_3\Big(M^2_{\tilde{u}_R^j}, M^2_{\tilde{u}_L^i}, |\mu|^2\Big), \\
\Gamma_{ii}^d = {} & \frac{2\alpha_s}{3\pi}(\mathbf{Y}_d)_{ii}\mu^\star M_g^\star I_3\Big(M^2_{\tilde{d}_L^i}, M^2_{\tilde{d}_R^i}, |M_g|^2\Big) \\
& + \frac{2\alpha_s}{3\pi}\sum_{j=1}^{3}(\mathbf{Y}_d)_{jj}\mu^\star M_g^\star M_{\tilde{D}}^4 I_5\Big(M^2_{\tilde{d}_L^i}, M^2_{\tilde{d}_L^j}, M^2_{\tilde{d}_R^j}, M^2_{\tilde{d}_R^i}, |M_g|^2\Big)\big(\delta_{ij}^d\big)_{RR}\big(\delta_{ji}^d\big)_{LL} \\
& + \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\sum_{j=1}^{3}\big(\mathbf{Y}_u^{A\dagger}\big)_{ij}(\mathbf{Y}_u)_{ji}\mu^\star I_3\Big(M^2_{\tilde{u}_R^j}, M^2_{\tilde{u}_L^i}, |\mu|^2\Big),
\end{aligned} \tag{5}$$

for the diagonal elements, and

$$\begin{aligned}
\gamma^d_{ij} &= \frac{2\alpha_s}{3\pi}\big(\mathbf{Y}^A_d\big)_{ii} M^\star_g M^2_{\tilde{D}} I_4\Big(M^2_{\tilde{d}^j_L}, M^2_{\tilde{d}^i_L}, M^2_{\tilde{d}^i_R}, |M_g|^2\Big)\big(\delta^d_{ij}\big)_{LL} \\
&+ \frac{2\alpha_s}{3\pi}\big(\mathbf{Y}^A_d\big)_{jj} M^\star_g M^2_{\tilde{D}} I_4\Big(M^2_{\tilde{d}^j_L}, M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^i_R}, |M_g|^2\Big)\big(\delta^d_{ij}\big)_{RR} \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_u\big)_{ij} (\mathbf{Y}_u)_{jj} |\mu|^2 I_3\Big(M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big) \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_u\big)_{ii} (\mathbf{Y}_u)_{ij} |\mu|^2 I_3\Big(M^2_{\tilde{u}^i_R}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big) \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_u\big)_{jj} (\mathbf{Y}_u)_{jj} |\mu|^2 M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big)\big(\delta^u_{ij}\big)_{LL} \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_u\big)_{ii} (\mathbf{Y}_u)_{jj} |\mu|^2 M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_R}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big)\big(\delta^u_{ij}\big)_{RR}, \\
\Gamma^d_{ij} &= \frac{2\alpha_s}{3\pi}\ (\mathbf{Y}_d)_{ii} \mu^\star M^\star_g M^2_{\tilde{D}}\ I_4\Big(M^2_{\tilde{d}^j_L}, M^2_{\tilde{d}^i_L}, M^2_{\tilde{d}^i_R}, |M_g|^2\Big)\big(\delta^d_{ij}\big)_{LL} \\
&+ \frac{2\alpha_s}{3\pi}\ (\mathbf{Y}_d)_{jj} \mu^\star M^\star_g M^2_{\tilde{D}} I_4\Big(M^2_{\tilde{d}^j_L}, M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^i_R}, |M_g|^2\Big)\big(\delta^d_{ij}\big)_{RR} \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_u\big)_{ij} (\mathbf{Y}_u)_{jj} \mu^\star I_3\Big(M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big) \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_u\big)_{ii} (\mathbf{Y}_u)_{ij} \mu^\star I_3\Big(M^2_{\tilde{u}^i_R}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big) \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_u\big)_{jj} (\mathbf{Y}_u)_{jj} \mu^\star M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big)\big(\delta^u_{ij}\big)_{LL} \\
&+ \frac{(\mathbf{Y}_d)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_u\big)_{ii} (\mathbf{Y}_u)_{jj} \mu^\star M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_R}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big)\big(\delta^u_{ij}\big)_{RR},
\end{aligned} \tag{6}$$

for the off-diagonal elements. These expressions (5) and (6), with $i, j = 1, 2, 3$, complete the radiative corrections to down quark interactions with Higgs fields. Repeating a similar analysis for the up quark sector, one finds

$$\begin{aligned}
\gamma^u_{ii} &= \frac{2\alpha_s}{3\pi}\big(\mathbf{Y}^A_u\big)_{ii} M^\star_g I_3\Big(M^2_{\tilde{u}^i_L}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big) \\
&+ \frac{2\alpha_s}{3\pi}\sum_{j=1}^{3}\big(\mathbf{Y}^A_u\big)_{jj} M^\star_g M^4_{\tilde{U}} I_5\Big(M^2_{\tilde{u}^i_L}, M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big)\big(\delta^u_{ij}\big)_{RR}\big(\delta^u_{ji}\big)_{LL} \\
&+ \frac{(\mathbf{Y}_u)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_d\big)_{ii} (\mathbf{Y}_d)_{ii} |\mu|^2 I_3\Big(M^2_{\tilde{d}^i_R}, M^2_{\tilde{u}^i_L}, |\mu|^2\Big), \\
\Gamma^u_{ii} &= \frac{2\alpha_s}{3\pi}(\mathbf{Y}_u)_{ii} \mu^\star M^\star_g I_3\Big(M^2_{\tilde{u}^i_L}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big) \\
&+ \frac{2\alpha_s}{3\pi}\sum_{j=1}^{3}(\mathbf{Y}_u)_{jj} \mu^\star M^\star_g M^4_{\tilde{U}} I_5\Big(M^2_{\tilde{u}^i_L}, M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big)\big(\delta^u_{ij}\big)_{RR}\big(\delta^u_{ji}\big)_{LL} \\
&+ \frac{(\mathbf{Y}_u)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_d\big)_{ii} (\mathbf{Y}_d)_{ii} \mu^\star I_3\Big(M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^i_L}, |\mu|^2\Big),
\end{aligned} \tag{7}$$

for the entries at the diagonal, and

$$
\begin{aligned}
\gamma^u_{ij} = {} & \frac{2\alpha_s}{3\pi}\big(\mathbf{Y}^A_u\big)_{ij} M_g^\star I_3\Big(M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big) \\
& + \frac{2\alpha_s}{3\pi}\big(\mathbf{Y}^A_u\big)_{ii} M_g^\star M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^i_L}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big)\big(\delta^u_{ij}\big)_{LL} \\
& + \frac{2\alpha_s}{3\pi}\big(\mathbf{Y}^A_u\big)_{jj} M_g^\star M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big)\big(\delta^u_{ij}\big)_{RR} \\
& + \frac{(\mathbf{Y}_u)_{ij}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_d\big)_{jj}(\mathbf{Y}_d)_{jj}|\mu|^2 I_3\Big(M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^j_L}, |\mu|^2\Big) \\
& + \frac{(\mathbf{Y}_u)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_d\big)_{jj}(\mathbf{Y}_d)_{jj}|\mu|^2 M^2_{\tilde{D}} I_4\Big(M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^j_L}, M^2_{\tilde{d}^i_L}, |\mu|^2\Big)\big(\delta^d_{ij}\big)_{LL} \\
& + \frac{(\mathbf{Y}_u)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^\dagger_d\big)_{ii}(\mathbf{Y}_d)_{jj}|\mu|^2 M^2_{\tilde{D}} I_4\Big(M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^i_R}, M^2_{\tilde{d}^i_L}, |\mu|^2\Big)\big(\delta^d_{ij}\big)_{RR}, \\
\Gamma^u_{ij} = {} & \frac{2\alpha_s}{3\pi}(\mathbf{Y}_u)_{ij}\mu^\star M_g^\star I_3\Big(M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big) \\
& + \frac{2\alpha_s}{3\pi}(\mathbf{Y}_u)_{ii}\mu^\star M_g^\star M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^i_L}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big)\big(\delta^u_{ij}\big)_{LL} \\
& + \frac{2\alpha_s}{3\pi}(\mathbf{Y}_u)_{jj}\mu^\star M_g^\star M^2_{\tilde{U}} I_4\Big(M^2_{\tilde{u}^j_L}, M^2_{\tilde{u}^j_R}, M^2_{\tilde{u}^i_R}, |M_g|^2\Big)\big(\delta^u_{ij}\big)_{RR} \\
& + \frac{(\mathbf{Y}_u)_{ij}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_d\big)_{jj}(\mathbf{Y}_d)_{jj}\mu^\star I_3\Big(M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^i_L}, |\mu|^2\Big) \\
& + \frac{(\mathbf{Y}_u)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_d\big)_{jj}(\mathbf{Y}_d)_{jj}\mu^\star M^2_{\tilde{D}} I_4\Big(M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^j_L}, M^2_{\tilde{d}^i_L}, |\mu|^2\Big)\big(\delta^d_{ij}\big)_{LL} \\
& + \frac{(\mathbf{Y}_u)_{ii}}{(4\pi)^2}\big(\mathbf{Y}^{A\,\dagger}_d\big)_{ii}(\mathbf{Y}_d)_{jj}\mu^\star M^2_{\tilde{D}} I_4\Big(M^2_{\tilde{d}^j_R}, M^2_{\tilde{d}^i_R}, M^2_{\tilde{d}^i_L}, |\mu|^2\Big)\big(\delta^d_{ij}\big)_{RR},
\end{aligned}
\tag{8}
$$

for the intergenerational ones. In these expressions $M_{\tilde{U},\tilde{D}}$ stand for the average up and down squark masses, and

$$
\big(\delta^{u,d}_{ij}\big)_{LL,RR} \equiv \big(M^2_{U,D}\big)^{ij}_{LL,RR}\big/M^2_{\tilde{U},\tilde{D}} \tag{9}
$$

are the mass insertions (MI) whose phases and sizes parametrize, respectively, the CP and flavor violations from the intergenerational entries of $(M^2_{U,D})_{LL,RR}$. Note that all entries of $\gamma^{u,d}$ and $\Gamma^{u,d}$ are computed at one loop approximation, and SUSY flavor violation effects are treated at single MI level everywhere except the diagonal entries which include dominant SUSY QCD contributions with two MIs in addition to the leading zero MI diagrams. The radiative corrections depend on the loop functions $I_{3,4,5}$, where

$$
\begin{aligned}
I_n\big(m_1^2, m_2^2, \ldots, m_n^2\big) = (-1)^{n+1}\Gamma(n-2)\int_0^1 dx_1 \int_0^{1-x_1} dx_2 \cdots \int_0^{1-x_1-\cdots-x_{n-2}} dx_{n-1} \\
\times \big(x_1 m_1^2 + x_2 m_2^2 + \cdots + (1 - x_1 - \cdots - x_{n-1})m_n^2\big)^{-n},
\end{aligned}
\tag{10}
$$

which approach, respectively, to $1/2m^2$, $-1/6m^4$ and $1/12m^6$ for $n = 3$, 4 and 5 when their arguments are equal.

An important aspect of the nonholomorphic Yukawa structures $\gamma^{u,d}$ and $\Gamma^{u,d}$ is that they depend only on the ratio of the soft masses not on their absolute scale. This property guarantees that these radiative corrections remain nonvanishing even if $M_{\mathrm{SUSY}} \gg m_t$. The simplest case corresponds to an approximate universality of the soft masses, $|\mu| \sim |M_g| \sim M_{\tilde{u}^j_{L,R}} \sim M_{\tilde{d}^j_{L,R}} \sim |(A_{u,d})_{ii}| \sim M_{\tilde{U}} \sim M_{\tilde{D}} \equiv M_{\mathrm{SUSY}}$, in which case $\gamma^{u,d}$ and $\Gamma^{u,d}$

depend only on the gauge and Yukawa couplings in addition to CP and flavor violation textures from the soft masses. Although such a universality is not likely to occur at low energies even if it holds at the scale of local SUSY breaking, it proves useful in illustrating the salient features of the Higgs interactions with quarks. Using the limiting forms of the loop functions (10) one obtains

$$
\begin{aligned}
\gamma_{ii}^{d} \implies & (\mathbf{Y}_d)_{ii}\left[\frac{\alpha_s}{3\pi}e^{i(\theta_{ii}^{d}-\theta_g)}+\frac{1}{32\pi^2}\sum_{j=1}^{3}(\mathbf{Y}_u^{\dagger})_{ij}(\mathbf{Y}_u)_{ji}\right]+\frac{\alpha_s}{18\pi}\sum_{j=1}^{3}(\mathbf{Y}_d)_{jj}(\delta_{ij}^{d})_{RR}(\delta_{ji}^{d})_{LL}e^{i(\theta_{jj}^{d}-\theta_g)},\\
\Gamma_{ii}^{d} \implies & (\mathbf{Y}_d)_{ii}\left[\frac{\alpha_s}{3\pi}e^{-i(\theta_\mu+\theta_g)}+\frac{1}{32\pi^2}\sum_{j=1}^{3}(\mathbf{Y}_u^{\dagger})_{ij}(\mathbf{Y}_u)_{ji}(\delta_{ji}^{u})_A^{\star}e^{-i\theta_\mu}\right]\\
& +\frac{\alpha_s}{18\pi}\sum_{j=1}^{3}(\mathbf{Y}_d)_{jj}(\delta_{ij}^{d})_{RR}(\delta_{ji}^{d})_{LL}e^{-i(\theta_\mu+\theta_g)},\\
\gamma_{ij}^{d} \implies & (\mathbf{Y}_d)_{ii}\left[-\frac{\alpha_s}{9\pi}(\delta_{ij}^{d})_{LL}e^{i(\theta_{ii}^{d}-\theta_g)}\right.\\
& \left.+\frac{1}{96\pi^2}\left\{3(\mathbf{Y}_u^{\dagger})_{ij}(\mathbf{Y}_u)_{jj}+3(\mathbf{Y}_u^{\dagger})_{ii}(\mathbf{Y}_u)_{ij}-(\mathbf{Y}_u^{\dagger})_{jj}(\mathbf{Y}_u)_{jj}(\delta_{ij}^{u})_{LL}-(\mathbf{Y}_u^{\dagger})_{ii}(\mathbf{Y}_u)_{jj}(\delta_{ij}^{u})_{RR}\right\}\right]\\
& -(\mathbf{Y}_d)_{jj}\left[\frac{\alpha_s}{9\pi}(\delta_{ij}^{d})_{RR}e^{i(\theta_{jj}^{d}-\theta_g)}\right],\\
\Gamma_{ij}^{d} \implies & (\mathbf{Y}_d)_{ii}\left[-\frac{\alpha_s}{9\pi}(\delta_{ij}^{d})_{LL}e^{-i(\theta_\mu+\theta_g)}+\frac{1}{96\pi^2}\left\{3(\mathbf{Y}_u^{\dagger})_{ij}(\mathbf{Y}_u)_{jj}(\delta_{ji}^{u})_A^{\star}e^{-i\theta_\mu}+3(\mathbf{Y}_u^{\dagger})_{ii}(\mathbf{Y}_u)_{ij}e^{-i(\theta_{ii}^{u}+\theta_\mu)}\right.\right.\\
& \left.\left.-(\mathbf{Y}_u^{\dagger})_{jj}(\mathbf{Y}_u)_{jj}(\delta_{ij}^{u})_{LL}e^{-i(\theta_{jj}^{u}+\theta_\mu)}-(\mathbf{Y}_u^{\dagger})_{ii}(\mathbf{Y}_u)_{jj}(\delta_{ij}^{u})_{RR}e^{-i(\theta_{ii}^{u}+\theta_\mu)}\right\}\right]\\
& -(\mathbf{Y}_d)_{jj}\left[\frac{\alpha_s}{9\pi}(\delta_{ij}^{d})_{RR}e^{-i(\theta_\mu+\theta_g)}\right],
\end{aligned}
\tag{11}
$$

for down quark sector, and

$$
\begin{aligned}
\gamma_{ii}^{u} \implies & (\mathbf{Y}_u)_{ii}\left[\frac{\alpha_s}{3\pi}e^{i(\theta_{ii}^{u}-\theta_g)}+\frac{1}{32\pi^2}(\mathbf{Y}_d^{\dagger})_{ii}(\mathbf{Y}_d)_{ii}\right]+\frac{\alpha_s}{18\pi}\sum_{j=1}^{3}(\mathbf{Y}_u)_{jj}(\delta_{ij}^{u})_{RR}(\delta_{ji}^{u})_{LL}e^{i(\theta_{jj}^{u}-\theta_g)},\\
\Gamma_{ii}^{u} \implies & (\mathbf{Y}_u)_{ii}\left[\frac{\alpha_s}{3\pi}e^{-i(\theta_\mu+\theta_g)}+\frac{1}{32\pi^2}(\mathbf{Y}_d^{\dagger})_{ii}(\mathbf{Y}_d)_{ii}e^{-i(\theta_{ii}^{d}+\theta_\mu)}\right]\\
& +\frac{\alpha_s}{18\pi}\sum_{j=1}^{3}(\mathbf{Y}_u)_{jj}(\delta_{ij}^{u})_{RR}(\delta_{ji}^{u})_{LL}e^{-i(\theta_\mu+\theta_g)},\\
\gamma_{ij}^{u} \implies & (\mathbf{Y}_u)_{ij}\left[\frac{\alpha_s}{3\pi}e^{i(\theta_{ij}^{u}-\theta_g)}+\frac{1}{32\pi^2}(\mathbf{Y}_d^{\dagger})_{jj}(\mathbf{Y}_d)_{jj}\right]\\
& -(\mathbf{Y}_u)_{ii}\left[\frac{\alpha_s}{9\pi}(\delta_{ij}^{u})_{LL}e^{i(\theta_{ii}^{u}-\theta_g)}+\frac{1}{96\pi^2}\left\{(\mathbf{Y}_d^{\dagger})_{jj}(\mathbf{Y}_d)_{jj}(\delta_{ij}^{d})_{LL}+(\mathbf{Y}_d^{\dagger})_{ii}(\mathbf{Y}_d)_{jj}(\delta_{ij}^{d})_{RR}\right\}\right]\\
& -(\mathbf{Y}_u)_{jj}\left[\frac{\alpha_s}{9\pi}(\delta_{ij}^{u})_{RR}e^{i(\theta_{jj}^{u}-\theta_g)}\right],
\end{aligned}
$$

$$\begin{aligned}\Gamma^u_{ij} \implies\ & (\mathbf{Y}_u)_{ij}\left[\frac{\alpha_s}{3\pi}e^{-i(\theta_\mu+\theta_g)} + \frac{1}{32\pi^2}\left(\mathbf{Y}_d^\dagger\right)_{jj}(\mathbf{Y}_d)_{jj}e^{-i(\theta^d_{jj}+\theta_\mu)}\right] \\ & - (\mathbf{Y}_u)_{ii}\left[\frac{\alpha_s}{9\pi}\left(\delta^u_{ij}\right)_{LL}e^{-i(\theta_\mu+\theta_g)}\right. \\ & \left. + \frac{1}{96\pi^2}\left\{\left(\mathbf{Y}_d^\dagger\right)_{jj}(\mathbf{Y}_d)_{jj}\left(\delta^d_{ij}\right)_{LL}e^{-i(\theta^d_{jj}+\theta_\mu)} + \left(\mathbf{Y}_d^\dagger\right)_{ii}(\mathbf{Y}_d)_{jj}\left(\delta^d_{ij}\right)_{RR}e^{-i(\theta^d_{ii}+\theta_\mu)}\right\}\right] \\ & - (\mathbf{Y}_u)_{jj}\left[\frac{\alpha_s}{9\pi}\left(\delta^u_{ij}\right)_{RR}e^{-i(\theta_\mu+\theta_g)}\right],\end{aligned} \tag{12}$$

for up quark sector after introducing the CP-odd phases $\theta_\mu \equiv \mathrm{Arg}[\mu]$, $\theta^{u,d}_{ii} \equiv \mathrm{Arg}[(A_{u,d})_{ii}]$, and $\theta_g \equiv \mathrm{Arg}[M_g]$. That $\mathbf{Y}_u$ is not diagonal causes all entries of A_u to contribute Γ^d, and this is parametrized via the insertions $(\delta^u_{ij})_A = (A_u)_{ij}/M_{\mathrm{SUSY}}$, similar to (9). Note that in (11) and (12) there is no explicit dependence on the soft masses except for the fact that all Yukawa and gauge couplings are to be evaluated at the scale M_{SUSY}.

The nonholomorphic Yukawa structures $\gamma^{u,d}$ and $\Gamma^{u,d}$ contribute to the quark masses as well as the Higgs boson couplings to quarks after the electroweak breaking. When determining the vacuum configuration and the physical Higgs bosons it is essential to include the radiative corrections to the Higgs potential from sparticle loops. In particular, the CP-odd phases contained in trilinear couplings and the μ parameter generate sizeable scalar–pseudoscalar mixings which prevent the Higgs bosons to have definite CP parities [13] (these results can be further refined using the most recent complete two loop calculation [14]). Defining the Higgs VEVs as $v_{u,d} = \sqrt{2}\,\langle H^0_{u,d}\rangle$, with $\tan\beta \equiv \langle H^0_u\rangle/\langle H^0_d\rangle$, the radiatively-corrected Yukawa coupling matrices take the form

$$\Upsilon_d = \mathbf{Y}_d - \gamma^d + \tan\beta\,\Gamma^d, \qquad \Upsilon_u = e^{i\delta}\left[\mathbf{Y}_u + \gamma^u - \cot\beta\,\Gamma^u\right], \tag{13}$$

where δ, the relative phase between the two doublets, is generated by the SUSY CP phases via the radiative corrections [15]. These Yukawa matrices can be diagonalized via the rotations $\mathbf{d}_L \to V^d_L\mathbf{d}_L$, $\mathbf{u}_L \to V^{0\dagger}V^u_L\mathbf{u}_L$, $\mathbf{d}_R \to V^d_R\mathbf{d}_R$, and $\mathbf{u}_R \to V^u_R\mathbf{u}_R$. Then the misalignment between the left-handed quarks in up and down sectors generates the physical CKM matrix

$$V = \left(V^u_L\right)^\dagger V^0 V^d_L, \tag{14}$$

which would be identical to V^0 in the absence of radiative corrections. The defining relations for the unitary matrices $V^{u,d}_{L,R}$ are

$$\begin{aligned} &\left(V^d_L\right)^\dagger(\Upsilon_d)^\dagger\Upsilon_d V^d_L = \overline{\mathbf{Y}}^2_d, \qquad V\left(V^d_L\right)^\dagger(\Upsilon_u)^\dagger\Upsilon_u V^d_L V^\dagger = \overline{\mathbf{Y}}^2_u, \\ &\left(V^d_R\right)^\dagger\Upsilon_d(\Upsilon_d)^\dagger V^d_R = \overline{\mathbf{Y}}^2_d, \qquad \left(V^u_R\right)^\dagger\Upsilon_u(\Upsilon_u)^\dagger\ V^u_R = \overline{\mathbf{Y}}^2_u,\end{aligned} \tag{15}$$

where $\overline{\mathbf{Y}}_d = \mathrm{diag}\{\bar{h}_d, \bar{h}_s, \bar{h}_b\}$ and $\overline{\mathbf{Y}}_u = \mathrm{diag}\{\bar{h}_u, \bar{h}_c, \bar{h}_t\}$ with $\bar{h}_i$ being the running (physical) Yukawa coupling of the ith generation, e.g., $\bar{h}_s = g_2\bar{m}_s/\sqrt{2}\,M_W\cos\beta$. Note that in the above $\tan\beta$, V as well as $\bar{h}_i$ are all evaluated at the scale M_{SUSY} via the RGE running of their experimental values at M_Z using the two-Higgs doublet model as the effective theory below M_{SUSY} [16,17]. The mass-eigenstate quark fields above interact with the Higgs bosons via

$$\begin{aligned}-\mathcal{L} = {} & \bar{\mathbf{d}}_R\overline{\mathbf{Y}}_d\mathbf{d}_L H^0_d + \bar{\mathbf{d}}_R\left(V^d_R\right)^\dagger\Gamma^d V^d_L\mathbf{d}_L\left(H^{0\star}_u - \tan\beta\,H^0_d\right) \\ & + \bar{\mathbf{u}}_R\overline{\mathbf{Y}}_u\mathbf{u}_L H^0_u - \bar{\mathbf{u}}_R\left(V^u_R\right)^\dagger\Gamma^u V^d_L V^\dagger\mathbf{u}_L\left(H^{0\star}_d - \cot\beta\,H^0_u\right),\end{aligned} \tag{16}$$

where it is clear that the flavor structures of $\Gamma^{u,d}$ are crucial for Higgs bosons to develop FCNC couplings. In particular, when $\Gamma^d \propto \mathbf{Y}_d$ and/or $\Gamma^u \propto \mathbf{Y}_u V^0$ there is no flavor-changing couplings of the neutral Higgs bosons to down and/or up quarks. The tree level CKM matrix V^0 is some unitary matrix which does not need to confront the experimental data unless the radiative effects contained in $\gamma^{u,d}$ and $\Gamma^{u,d}$ are vanishingly small. The allowed ranges for individual entries of V^0 depend on the size of the SUSY flavor and CP violation effects since the physical CKM matrix (14) must saturate at least the bounds from tree level FCNC processes [16]. The mixing matrices $V^{u,d}_{L,R}$ can be computed via perturbation theory for small flavor mixings. On the other hand, if the mixings are sizeable it is useful to employ direct diagonalization by first transforming $\Upsilon_{u,d}$ into the nearest-neighbour-interaction basis [18] and then solving for Yukawa couplings and tree level CKM elements using the techniques given in [19].

For determining the Higgs interactions with quarks (16) it is necessary to express the tree level parameters $(\mathbf{Y}_{u,d}, V^0)$ in terms of the physical ones $(\overline{\mathbf{Y}}_{u,d}, V)$ via (14) and (15). Since a given entry of $\Upsilon_{u,d}$ depends on the Yukawa couplings of other generations, a direct solution of (15) will eventually need a scanning of the parameter space by taking into account all the available constraints. However, for the purpose of illustrating the essential features of SUSY flavor and CP violation effects on Higgs–quark interactions it suffices to have an approximate solution for Yukawa couplings, i.e., one can neglect flavor mixings from V^0, and discard the SUSY electroweak corrections all together which induces $\sim 20\%$ error in estimating the bottom Yukawa. Furthermore, for compactness it is useful to use the limiting forms (11), (12) keeping in mind that size of the radiative corrections can be significantly altered if the universality assumption is relaxed. Within these approximations, the Yukawa couplings admit the solutions

$$\begin{aligned}
h_d &= \frac{g_2 \bar{m}_d}{\sqrt{2} M_W} \frac{\tan\beta}{1+\epsilon\tan\beta}\left[1 - \frac{\epsilon\tan\beta}{1+\epsilon\tan\beta}\left\{\frac{\bar{m}_s}{\bar{m}_d}\left(\delta^d_{12}\right)_{LR} + \frac{\bar{m}_b}{\bar{m}_d}\left(\delta^d_{13}\right)_{LR}\right\}\right],\\
h_s &= \frac{g_2 \bar{m}_s}{\sqrt{2} M_W} \frac{\tan\beta}{1+\epsilon\tan\beta}\left[1 - \frac{\epsilon\tan\beta}{1+\epsilon\tan\beta}\frac{\bar{m}_b}{\bar{m}_s}\left(\delta^d_{23}\right)_{LR}\right],\\
h_b &= \frac{g_2 \bar{m}_b}{\sqrt{2} M_W} \frac{\tan\beta}{1+\epsilon\tan\beta}, \qquad h_u = \frac{g_2 \bar{m}_u}{\sqrt{2} M_W}\left[1 - \epsilon_1\left\{\frac{\bar{m}_c}{\bar{m}_u}\left(\delta^u_{12}\right)_{LR} + \frac{\bar{m}_t}{\bar{m}_u}\left(\delta^u_{13}\right)_{LR}\right\}\right],\\
h_c &= \frac{g_2 \bar{m}_c}{\sqrt{2} M_W}\left[1 - \epsilon_2\frac{\bar{m}_t}{\bar{m}_c}\left(\delta^u_{23}\right)_{LR}\right], \qquad h_t = \frac{g_2 \bar{m}_t}{\sqrt{2} M_W},
\end{aligned} \tag{17}$$

where the SUSY flavor violation contributions are separated from the ones which already exist in the minimal flavor violation (MFV) scheme [8,9]. These expressions for Yukawas have been obtained by keeping only those terms not suppressed by $\tan\beta$ and linear in $(\delta^d_{ij})_{LR}$. The new parameters in (17) are defined as $\epsilon = (\alpha_s/3\pi)e^{-i(\theta_\mu+\theta_g)}$, $\epsilon_i = (\alpha_s/3\pi)e^{i(\theta^u_{ii}-\theta_g)}$, and

$$\left(\delta^d_{ij}\right)_{LR} = \frac{1}{6}\left(\delta^d_{ij}\right)_{RR}\left(\delta^d_{ji}\right)_{LL}, \qquad \left(\delta^u_{ij}\right)_{LR} = \frac{1}{6}e^{i(\theta^u_{jj}-\theta^u_{ii})}\left(\delta^u_{ij}\right)_{RR}\left(\delta^u_{ji}\right)_{LL} \tag{18}$$

which generate the effective LR transitions needed for correcting the diagonal Yukawa elements. In contrast to the MFV scheme, the Yukawa couplings acquire sizeable corrections from the those of the heavier generations as suggested by (17). Indeed, the radiative corrections to $h_d/\bar{h}_d$, $h_s/\bar{h}_s$, $h_u/\bar{h}_u$ and $h_c/\bar{h}_c$ involve, respectively, the large factors $\bar{m}_b/\bar{m}_d \sim (\tan\beta)^2_{\max}$, $\bar{m}_b/\bar{m}_s \sim (\tan\beta)_{\max}$, $\bar{m}_t/\bar{m}_u \sim (\tan\beta)^3_{\max}$, and $\bar{m}_t/\bar{m}_c \sim (\tan\beta)^2_{\max}$ with $(\tan\beta)_{\max} \lesssim \bar{m}_t/\bar{m}_b$. Unlike the light quarks, the top and bottom Yukawas remain stuck to their MFV values to a good approximation. Therefore, the SUSY flavor violation sources mainly influence the light sector whereby modifying several processes they participate. The modifications in the Yukawa couplings are important even at low $\tan\beta$ values. As an example, consider $(\delta^d_{13})_{LR} \sim 10^{-2}$ for which $h_d/h_d^{\mathrm{MFV}} \simeq 0.02(2.11), -2.3(6.6), -4.6(17.7)$ for $\tan\beta = 5, 20, 40$ at $\theta_\mu + \theta_g \rightsquigarrow 0(\pi)$. Note that the Yukawas are enhanced especially for $\theta_\mu + \theta_g \rightsquigarrow \pi$ which is the point preferred by Yukawa-unified models such as SO(10).

In general, as $\tan\beta \to (\tan\beta)_{\max}$ the Yukawa couplings of down and strange quarks become approximately degenerate with the bottom Yukawa for $(\delta^d_{13,23})_{LR} \sim 0.1$ and $\theta_\mu + \theta_g \rightsquigarrow \pi$. There is no $\tan\beta$ enhancement for up quark sector but still the large ratio $\bar{m}_t/\bar{m}_u$ sizeably folds h_u compared to its SM value: $h_u \simeq 0.6 e^{i(\theta^u_{11}-\theta_g)}\bar{h}_c$ with $(\delta^u_{13})_{LR} \sim 0.1$. These spectacular enhancements in light quark Yukawas, though possible in a small corner of the parameter space, imply that the SUSY flavor violation effects can induce strong modifications in light quark couplings to Higgs bosons—an important aspect for both Higgs boson searches and FCNC processes to be detailed below.

At this point one may wonder if the leptonic Yukawas can also be enhanced. By replacing the gluino–squark loops with bino–slepton loops, one finds that the radiative corrections are actually down by two orders of magnitude compared to the quark sector even when $\tan\beta \sim (\tan\beta)_{\max}$ and bino is nearly degenerate with sleptons. Moreover, if bino is the dark matter candidate these threshold corrections are further suppressed. In summary, as follows from (17), the SUSY flavor and CP violations modify the hierarchy of the Yukawa couplings strongly even for small or moderate $\tan\beta$ values. In fact, when $\tan\beta$ assumes its maximal value and the MIs are $\mathcal{O}(1)$ one finds that (i) the down quark Yukawas acquires an approximate universality, (ii) the up quark Yukawa becomes degenerate with the charm, and (iii) the Yukawa couplings of the third generation quarks, of the charm quark, and of all leptons remain stuck to their MFV values to a good approximation.

The couplings of Higgs bosons to quarks are fully determined by (16). The off-diagonal entries of $V^d_{R,L}$ in (15) are approximately given by $-(1/3)\epsilon\tan\beta\,(\delta^d_{ij})_{RR,LL}$. The corresponding entries of $V^u_{L,R}$ are down by a factor or $1/\tan\beta$. Clearly, for any regime of the parameter values, the texture of the tree level CKM matrix V^0 plays an important role in generating the physical CKM matrix V. As for an approximate analysis, one may take $V^{u,d}_{L,R}$ diagonal and neglect scalar–pseudoscalar mixings in the Higgs sector [13]. Note that errors made in these approximations are sensitive to $\tan\beta$; therefore, they must be avoided in an accurate treatment of the problem. Modulo these approximations, assuming for simplicity a universal phase for $(A_d)_{ii}$ and $(A_u)_{ii}$ each, the couplings of the Higgs bosons to quarks take the form

$$\begin{aligned}
-\mathcal{L} = {} & \frac{g_2 \bar{m}_{d^i}(M_{\rm SUSY})}{2M_W}\left[\frac{h^i_d}{\bar{h}^i_d}\tan\beta\, C^d_a + \left(\frac{h^i_d}{\bar{h}^i_d} - 1\right)\left(e^{i(\theta^d_{ii}+\theta_\mu)} C^d_a - C^{u\star}_a\right)\right]\bar{d}^i_R d^i_L H_a \\
& + \frac{g_2 \bar{m}_{d^i}(M_{\rm SUSY})}{6M_W}\epsilon\tan\beta\left[\frac{h^i_d}{\bar{h}^i_d}\left(\delta^d_{ij}\right)_{LL} + \frac{h^j_d}{\bar{h}^i_d}\left(\delta^d_{ij}\right)_{RR}\right]\left(\tan\beta\, C^d_a - C^{u\star}_a\right)\bar{d}^i_R d^j_L H_a \\
& + \frac{g_2 \bar{m}_{u^i}(M_{\rm SUSY})}{2M_W}\left[C^u_a + e^{-i(\theta^u_{ii}+\theta_\mu)}\left(\frac{h^i_u}{\bar{h}^i_u} - 1\right)\left(C^{d\star}_a - \cot\beta\, C^u_a\right)\right]\bar{u}^i_R u^i_L H_a \\
& + \frac{g_2 \bar{m}_{u^i}(M_{\rm SUSY})}{6M_W}\epsilon\left[\frac{h^i_u}{\bar{h}^i_u}\left(\delta^u_{ij}\right)_{LL} + \frac{h^j_u}{\bar{h}^i_u}\left(\delta^u_{ij}\right)_{RR}\right]\left(C^{d\star}_a - \cot\beta\, C^u_a\right)\bar{u}^i_R u^j_L H_a,
\end{aligned} \tag{19}$$

where $C^d_a \equiv \{-\sin\alpha, \cos\alpha, i\sin\beta, -i\cos\beta\}$ and $C^u_a \equiv \{\cos\alpha, \sin\alpha, i\cos\beta, i\sin\beta\}$ in the basis $H_a \equiv \{h, H, A, G\}$. In deriving (19) $\tan\beta$ is assumed to be large though not necessarily close to $(\tan\beta)_{\max}$. That the MFV contributions as well as $\mathcal{O}[(\tan\alpha_{ij})^2]$ terms are absent in the flavor-violating couplings, that the CKM matrix does not have a direct contribution, that the Higgs bosons assume well-defined CP parities, ... are just the artefacts of the simplifying assumptions made above. These missing pieces can be incorporated into the effective Lagrangian by a more accurate analysis using the exact formulae derived before.

The effective Lagrangian (19) for Higgs–quark interactions has a multitude of phenomenological implications covering hadronic, atomic as well as Higgs systems. Quite generically, all the SUSY effects contained in (19) scale as $1/M_A^2$ whereas the analogous effective Lagrangian for gauge boson interactions with quarks as well as four-fermion operators generated by box diagrams do so as $1/M^2_{\rm SUSY}$ [3,4]. Consequently, when the SUSY Higgs sector lies at the weak scale the Higgs–quark interactions probe superpartners at all scales, as dictated by (11), (12) and (19), via their persistent effects on low energy observables. On the other hand, the non-Higgs SUSY contributions can be important only when $M_{\rm SUSY}$ lies around the weak scale. When discussing the Higgs boson effects on various observables it is convenient to separate the atomic and hadronic observables from those in the Higgs system:

1. Implications for hadronic and atomic systems

For such observables the Higgs boson effects filter through Higgs mediation which may (FCNC observables) or may not (atomic EDMs) require flavor violation. Starting with the FCNC processes, one notes that $\Delta F = 2$ transitions proceed via double Higgs penguins. The dominant contributions come from the scalar operators of the form $\bar{h}_R l_L \bar{h}_L l_R$, where $(h,l) = (s,d), (c,u), (b,d)$ and (b,s) for K^0–$\overline{K}^0$, D^0–$\overline{D}^0$, B_d^0–$\overline{B}_d^0$ and B_s^0–$\overline{B}_s^0$ mixings, respectively. The Wilson coefficients of these operators follow from the flavor-changing parts of (19), and at large $\tan\beta$ read as

$$C_2^{LR}\left(K^0\text{–}\overline{K}^0\right) = -\frac{1}{9}|\epsilon|^2 \tan^2\beta \left(h_s\left(\delta_{21}^d\right)_{LL} + h_d\left(\delta_{21}^d\right)_{RR}\right)\left(h_d^\star\left(\delta_{21}^d\right)_{LL} + h_s^\star\left(\delta_{21}^d\right)_{RR}\right) \times \left(\frac{\sin^2(\alpha-\beta)}{M_H^2} + \frac{\cos^2(\alpha-\beta)}{M_h^2} + \frac{1}{M_A^2}\right), \tag{20}$$

with $C_2^{LR}(D^0\text{–}\overline{D}^0) = C_2^{LR}(K^0\text{–}\overline{K}^0)[\tan\beta \to 1, h_d \to h_u, h_s \to h_c]$. The expressions for $C_2^{LR}(B_{d,s}^0\text{–}\overline{B}_{d,s}^0)$ follow from (20) with obvious replacements:

$$C_2^{LR}\left(B_d^0\text{–}\overline{B}_d^0\right) = -\frac{1}{9}|\epsilon|^2 \tan^2\beta \left(h_b\left(\delta_{31}^d\right)_{LL} + h_d\left(\delta_{31}^d\right)_{RR}\right)\left(h_d^\star\left(\delta_{31}^d\right)_{LL} + h_b^\star\left(\delta_{31}^d\right)_{RR}\right) \times \left(\frac{\sin^2(\alpha-\beta)}{M_H^2} + \frac{\cos^2(\alpha-\beta)}{M_h^2} + \frac{1}{M_A^2}\right), \tag{21}$$

where it is clear from both (20) and (21) that the Higgs double penguins contribute to the CP violation in mixing—a property not present in the MFV scheme [8,9]. For all four distinct meson systems C_2^{LR} is quadratic in the Yukawa coupling of the heaviest quark, and requires six MIs when the radiative corrections in (17) dominate. It is clear that strength of C_2^{LR} depends on the absolute sizes of MIs as well as relative phases between the LL and RR sector contributions.

The Higgs exchange diagrams with a single flavor flip generate $\Delta F = 1$ transitions of which $K_L \to \pi e^+ e^-$, $B_d \to \phi K_s$, $B_s \to \mu^+\mu^-$, $D \to \pi\pi\pi$, $B_d \to (\pi, K)\ell^+\ell^-$ form a few examples. For instance, at the matching scale the Higgs penguins generate the scalar operator $\bar{u}_R c_L \bar{d}_R d_L$ with the coefficient

$$C_1^{LR}(D \to \pi\pi\pi) = -\frac{1}{3}\epsilon h_d \left(h_u\left(\delta_{12}^u\right)_{LL} + h_c\left(\delta_{12}^u\right)_{RR}\right)\left(\frac{\sin^2(\alpha-\beta)}{M_H^2} + \frac{\cos^2(\alpha-\beta)}{M_h^2} + \frac{1}{M_A^2}\right), \tag{22}$$

which is responsible for D meson decays into three pions. Similarly, semileptonic operator $\bar{s}_R b_L \bar{\ell}_R \ell_L$, generated by Higgs exchange, contributes to pure leptonic decay of B_d meson via

$$C_1^{LR}(B_d \to \bar{\mu}\mu) = -\frac{1}{3}\epsilon \tan\beta\, h_\mu \left(h_d\left(\delta_{13}^d\right)_{LL} + h_b\left(\delta_{13}^d\right)_{RR}\right)\left(\frac{\sin^2(\alpha-\beta)}{M_H^2} + \frac{\cos^2(\alpha-\beta)}{M_h^2} + \frac{1}{M_A^2}\right), \tag{23}$$

with a similar expression for B_s mode. It is clear that the strength of C_1^{LR} is directly correlated with the associated C_2^{LR} coefficient [8,9,20].

The EDMs of heavy atoms are sensitive to CP-violating semileptonic four-fermion operators [10,11,21] in addition to the electron EDM contribution [22]. Especially operators of the form $\bar{q}q\bar{e}i\gamma_5 e$ couple the spin of the electron cloud to the nuclear density, and the resulting contribution to the EDM of the atom grows with its atomic number. For example, the EDM of ^{205}Tl is given by $d_{\text{Tl}} = -585 d_e - 8.5 \times 10^{-17} C_S$ e cm, where d_e is the electron EDM and

$$C_S = 5.5 \times 10^{-10}\left(\frac{100\ \text{GeV}}{M_A}\right)^2 \text{Im}\left[(1 - 0.25\kappa)h_b^\star h_e + 3.3\kappa h_s^\star h_e + 0.5 h_b^\star h_e\right], \tag{24}$$

with $\kappa \sim 1$ parametrizing the uncertainty in $\langle N|m_s \bar{s}s|N\rangle$.

The hadronic and atomic system observables examplified by (20)–(24) are partly under experimental investigation and are partly constrained by the existing data. In general, the bounds on these Higgs-exchange amplitudes depend on the size of non-Higgs contributions to a given observable. Conversely, the existing bounds on various SUSY parameters [4,23] derived by considering only the non-Higgs effects can be significantly modified once the Higgs mediation effects are taken into account. This has already been shown to happen for the atomic EDMs [11]: the Higgs-exchange amplitude largely cancels with the two-loop electron EDM contribution in certain regions of the parameter space. Therefore, it is after a combined analysis of the Higgs and non-Higgs contributions that one can arrive at conclusions about the size and phase contents of various MIs. Indeed, the present bounds on various flavor violation sources [4,23], following from meson mixings by taking into account only the gluino box contributions, can be relaxed or strengthened depending on the parameter space.

Among various FCNC observables, the pure leptonic decays of B mesons put stringent constraints on Higgs-mediated FCNC since the SM predictions for $\mathrm{BR}(B_{d,s} \to \bar{\mu}\mu) \sim (1.5, 35) \times 10^{-10}$ which are roughly three orders of magnitude below the current experimental bounds $(6.8, 26) \times 10^{-7}$ [24] whereas the Higgs-exchange contributions well exceed the bounds even in minimal flavor violation scheme for $\tan\beta \gtrsim 50$ [8]. Due to the smallness of the SM background the Higgs effects on these decays are important irrespective of if M_{SUSY} is close to or far above the weak scale. In the present framework, to agree with the bounds the Wilson coefficient (23) must be suppressed in other words the quantity $h_d^i(\delta_{ij}^d)_{LL} + h_d^j(\delta_{ij}^d)_{RR}$ (with $i = 3,\ j = 1, 2$ and vice versa) must be sufficiently small depending on $\tan\beta$ and M_A. This constraint is important since if it forces $(\delta_{13,23}^d)_{LL,RR}$ to take small values the light quark Yukawas cannot assume sizeable enhancements noted before. Since (19) is far from being precise enough (neglect of flavor violation from V^0 and $V_{L,R}^d$ as well as the SUSY electroweak corrections) to perform a scanning of the parameter space, it is useful check if (23) can be suppressed in parameter regions with low M_A, large $\tan\beta$ and $\mathcal{O}(1)$ MIs. This indeed happens. To see this one incorporates terms higher order in MIs into the Yukawa couplings listed in (17). For instance, the down quark Yukawa takes the form

$$h_d = h_d^{\mathrm{MFV}} \frac{1 - a^2(\delta_{23}^d)_{LR}(\delta_{32}^d)_{LR} - aA_{12}\frac{\bar{m}_s}{\bar{m}_d} - aA_{13}\frac{\bar{m}_b}{\bar{m}_d}}{1 - a^2 A_2 - a^3 A_3}, \tag{25}$$

where

$$a = \frac{\epsilon \tan\beta}{1 + \epsilon \tan\beta}, \qquad A_{12} = (\delta_{12}^d)_{LR} - a(\delta_{13}^d)_{LR}(\delta_{32}^d)_{LR}, \qquad A_{13} = A_{12}(2 \leftrightarrow 3),$$
$$A_2 = |(\delta_{12}^d)_{LR}|^2 + |(\delta_{13}^d)_{LR}|^2 + |(\delta_{23}^d)_{LR}|^2, \qquad A_3 = (\delta_{12}^d)_{LR}(\delta_{23}^d)_{LR}(\delta_{31}^d)_{LR} + \mathrm{h.c.}$$

Using these improved expressions for Yukawas in (19), one finds that the flavor-changing Higgs vertices bsH^a and bdH^a become vanishingly small for $\tan\beta \simeq 60$ when all MIs are $\mathcal{O}(1)$, for $\tan\beta \simeq 65$ when $(\delta_{12}^d)_{LL,RR} \simeq 0$, and finally for $\tan\beta \simeq 68$ when $(\delta_{12}^d)_{LL} \simeq -(\delta_{12}^d)_{RR}$ provided that $\phi_\mu + \phi_g \rightsquigarrow \pi$ in all three cases. The existence of such a parameter domain is important in that it allows one to overcome constraints from $B_{d,s} \to \bar{\mu}\mu$ without suppressing the MIs $(\delta_{13,23}^d)_{LL,RR}$ which are crucial for enhancing the light quark Yukawas. Of course, saturating the bounds implies by no means vanishing of bsH^a and bdH^a vertices instead what is needed is to suppress such flavor-changing entries without enforcing the MIs to unobservably small values. If $M_{\mathrm{SUSY}} \gg m_t$ vanishing of such vertices reduces $B_{d,s} \to \bar{\ell}\ell$ and $B_{d,s} \to (K, \pi)\ell^+\ell^-$ decays as well as $B_{d,s}^0$–$\overline{B}_{d,s}^0$ mixings to their SM predictions with $\mathcal{O}(1)$ flavor mixings between the third and first generations of quarks. If $M_{\mathrm{SUSY}} \sim m_t$, however, for flavor mixings to be still sizeable the Higgs exchange contributions to $B_{d,s}^0$–$\overline{B}_{d,s}^0$ should balance the gluino boxes [4,23] on top of suppressing $B_{d,s} \to \bar{\mu}\mu$ below the bounds. This can be decided only after a global analysis of all the existing FCNC data.

For flavor transitions between the first two generations, the relevant observables are K^0–$\overline{K}^0$ and D^0–$\overline{D}^0$ mixings as well as the rare K and D decays. The Higgs-exchange amplitudes (20) and (22) can be suppressed either by balancing them with the gluino boxes or by tuning the LL and RR pieces in (19) depending, respectively, on whether $M_{\mathrm{SUSY}} \sim m_t$ or $M_{\mathrm{SUSY}} \gg m_t$. One notices that, in the latter case, the equality $(\delta_{12}^{u,d})_{LL} \simeq -(\delta_{12}^{u,d})_{RR}$

automatically suppresses the FCNC in K and D systems without falsifying the aforementioned enhancements in light quark Yukawas which rest mainly on $(\delta^{u,d}_{13,23})_{LR}$.

The atomic EDMs, in particular, ^{205}Tl EDM, can be suppressed by balancing the contribution of (24) with d_e (especially its two-loop part) [11] if $M_{\rm SUSY}$ is close to the weak scale. In the opposite limit, $M_{\rm SUSY} \gg m_t$, (24) is the only contribution and its suppression is almost automatic in parameter regions where the Higgs-mediated FCNC in B system are suppressed.

2. Implications for Higgs boson searches

The sizes and phases of various MIs and other SUSY parameters that have survived the bounds from EDMs and FCNC observables can open new channels, strengthen or weaken the existing ones for Higgs boson production and decays. The relevant interactions can be read off from (19) for each quark and Higgs flavor. Concerning the collider searches for a light fundamental scalar, the main object is the lightest Higgs boson which possesses both flavor-changing and flavor-conserving couplings to quarks

$$\mathrm{Re}\big[g^d_{ii}\big]h\bar{d}^i d^i + \mathrm{Re}\big[g^u_{ii}\big]h\bar{u}^i u^i + \frac{1}{2}\Big\{\big(g^d_{ij} + g^{d\star}_{ji}\big)h\bar{d}^i d^j + \big(g^u_{ij} + g^{u\star}_{ji}\big)h\bar{u}^i u^j + \text{h.c.}\Big\}, \tag{26}$$

where $i \neq j$, and various couplings read as

$$\begin{aligned}
g^d_{ii} &= -\big(h^i_d\big)_{\rm SM}\frac{\sin\alpha}{\cos\beta}\left[1 + \left(\frac{h^i_d}{\bar{h}^i_d} - 1\right)\left(1 + \frac{1}{\tan\alpha\tan\beta}\right)\right],\\
g^d_{ij} &= -\big(h^i_d\big)_{\rm SM}\frac{\sin\alpha}{\cos\beta}\frac{\epsilon}{3}\left[\frac{h^i_d}{\bar{h}^i_d}\big(\delta^d_{ij}\big)_{LL} + \frac{h^j_d}{\bar{h}^i_d}\big(\delta^d_{ij}\big)_{RR}\right](\tan\beta + \cot\alpha),\\
g^u_{ii} &= \big(h^i_u\big)_{\rm SM}\frac{\cos\alpha}{\sin\beta}\left[1 + \left(1 - \frac{h^i_u}{\bar{h}^i_u}\right)e^{i(\theta^u_{ii} + \theta_\mu)}(\cot\beta + \tan\alpha)\right],\\
g^u_{ij} &= \big(h^i_u\big)_{\rm SM}\frac{\cos\alpha}{\sin\beta}\frac{\epsilon}{3}\left[\frac{h^i_u}{\bar{h}^i_u}\big(\delta^u_{ij}\big)_{LL} + \frac{h^j_u}{\bar{h}^i_u}\big(\delta^u_{ij}\big)_{RR}\right](\cot\beta + \tan\alpha),
\end{aligned} \tag{27}$$

where the Yukawa couplings are given in (17). Clearly, all flavor-diagonal couplings reduce to those in the SM, and all flavor-changing couplings vanish when the Higgs sector enters the decoupling regime, $M_A \gg M_Z$ [6]. In this limit the lightest Higgs becomes the standard Higgs boson, and all the aforementioned Higgs-mediated FCNC amplitudes, which are generated by the heavier Higgs bosons, are suppressed as $1/M_A^2$. Therefore, only with a light Higgs sector, i.e., $M_A \gtrsim M_Z$ or equivalently $|\cot\alpha| \ll \tan\beta$, that there exist observable SUSY effects in the decay channels of h.

The Higgs bosons possess both flavor-changing and flavor-conserving decay and production modes. For example, the lightest Higgs decays into down and strange quarks in both flavor-changing

$$\frac{\Gamma(h \to \bar{s}d + \bar{d}s)}{\Gamma(h \to \bar{b}b)} \simeq \left|\frac{\epsilon h_s + \epsilon^\star h^\star_d}{3h_b}\big(\delta^d_{21}\big)_{LL} + \frac{\epsilon h_d + \epsilon^\star h^\star_s}{3h_b}\big(\delta^d_{21}\big)_{RR}\right|^2 \tan^2\beta \tag{28}$$

and flavor-conserving fashion

$$\frac{\Gamma(h \to \bar{d}d)}{\Gamma(h \to \bar{b}b)} \simeq \left(\mathrm{Re}\left[\frac{h_d}{h_b}\right]\right)^2, \tag{29}$$

where the differences in phase spaces are neglected, and it is assumed that the Higgs sector is light, i.e., $|\cot\alpha| \ll \tan\beta$. Whether the ratios above achieve any observable significance depends on the sizes and phases of $h_{d,s}$ as well as $(\delta^d_{21})_{LL,RR}$, which eventually need a global analysis of all the relevant FCNC data.

It is useful to start analyzing (28) and (29) in the limit of enhanced Yukawas: $h_d \sim h_s \sim h_b$ and $h_u \sim h_c$. In the limit of heavy superpartners, $M_{\mathrm{SUSY}} \gg m_t$, bounds on B, D and K system FCNC can be satisfied with $\mathcal{O}(1)$ MIs, as discussed above. In this case, flavor-changing Higgs interactions can be sizeable only in the FCNC top decays [25] with $\Gamma(t \to ch) \simeq \Gamma(t \to uh)$, whose likelihood depends on future observations at Tevatron and LHC. Although all flavor-changing Higgs decay channels are shut-off by the FCNC data, decays into $\bar{q}q$ final states are maximally enhanced. Indeed, (29) implies that $\Gamma(h \to \bar{d}d) \simeq \Gamma(h \to \bar{s}s) \simeq \Gamma(h \to \bar{b}b)$ and $\Gamma(h \to \bar{u}u) \simeq \Gamma(h \to \bar{c}c)$. Therefore, $h \to \bar{b}b$ is no longer the dominant decay channel as expected in SM, instead all channels are equally possible. In fact, $\mathrm{BR}(h \to \bar{b}b)$ is typically $\sim 30\%$ which well below the SM expectation. One recalls that, in the parameter domains which lead to degenerate Yukawas $\Gamma(h \to \bar{b}b)$ is typically an order of magnitude larger than that in the SM [26], and thus, the main signature of enhanced down Yukawas is not the suppression of $\bar{b}b$ production rate instead it is the drop in $\mathrm{BR}(h \to \bar{b}b)$ due to the strengthening of the other channels. Note that, even the existing LEP data can accommodate a light Higgs boson with mass $\lesssim 100$ GeV provided that the Yukawa couplings are comparable in size [27] in contrast to SM-like hierarchical couplings [28].

When M_{SUSY} is close to the weak scale, the FCNC observables receive contributions from not only the Higgs mediation but also from SUSY box and penguin diagrams. In this case, the allowed sizes of the MIs depend on if Higgs and non-Higgs contributions sufficiently cancel. Note that even if this happens one still needs to suppress the pure leptonic decay modes $B_{s,d} \to \ell^+\ell^-$ by tuning the SUSY parameters. This necessarily suppresses the corresponding flavor-changing Higgs decays $h \to \bar{b}(s,d) + \overline{(s,d)}b$ [29]. In case all the MIs survive FCNC constraints without excessive suppression, then (28) and (29) can both be sizeable and therefore $\mathrm{BR}(h \to \bar{b}b)$ can fall to 15–20% level as an optimistic estimate.

Having discussed the implications of FCNC data for Higgs decays and production for enhanced light quark Yukawas, it is useful to discuss (28) and (29) in a different parameter domain, i.e., suppose, though not realistic at all, that the Higgs-exchange contributions to FCNC data are negligible so that the MIs remain stuck to their bounds obtained via non-Higgs amplitudes:

$$\begin{aligned} &\left(\delta^d_{12}\right)_{LL} \simeq \left(\delta^d_{12}\right)_{RR} \simeq 8.0 \times 10^{-2}, \qquad \left(\delta^d_{13}\right)_{LL} \simeq \left(\delta^d_{13}\right)_{RR} \simeq \left(\delta^u_{12}\right)_{LL} \simeq \left(\delta^u_{12}\right)_{RR} \simeq \left(\delta^u_{13}\right)_{LL} \simeq 2.0 \times 10^{-1}, \\ &\left(\delta^d_{23}\right)_{LL} \simeq \left(\delta^d_{23}\right)_{RR} \simeq \left(\delta^u_{23}\right)_{LL} \simeq \left(\delta^u_{23}\right)_{RR} \simeq \left(\delta^u_{13}\right)_{RR} \simeq 1, \end{aligned}$$

as follows from the analyses of [4,23] for $M_{\mathrm{SUSY}} = 1$ TeV. For definiteness, take $\theta_\mu + \theta_g \rightsquigarrow \pi$, and consider $\tan\beta = 20$ and $\tan\beta = \bar{m}_t/\bar{m}_b \simeq 60$ as two sample points for illustrating low and high $\tan\beta$ behaviours. Then the down and strange quark Yukawas are enhanced as $h_d = (2.4 \times 10^{-3}, 0.04)h_b$ and $h_s = (0.06, 0.91)h_b$ for $\tan\beta = (20, 60)$. Consequently, (29) gives $\Gamma(h \to \bar{d}d) = (6 \times 10^{-4}\%, 0.14\%)\Gamma(h \to \bar{b}b)$, $\Gamma(h \to \bar{s}s) = (0.35\%, 83\%)\Gamma(h \to \bar{b}b)$. Similarly, from (28) it follows that $\Gamma(h \to \bar{s}d + \bar{d}s) = (7.2 \times 10^{-5}\%, 0.1\%)\Gamma(h \to \bar{b}b)$, $\Gamma(h \to \bar{b}d + \bar{d}b) = (0.12\%, 1.1\%)\Gamma(h \to \bar{b}b)$, and $\Gamma(h \to \bar{b}s + \bar{s}b) = (16\%, 96\%)\Gamma(h \to \bar{b}b)$. It is clear that enhancements in Yukawas depend crucially on the allowed sizes of the MIs as well as $\tan\beta$: at low $\tan\beta$ $h \to \bar{b}b$ is the dominant decay channel with a large branching fraction as in the SM. On the other hand, as $\tan\beta$ grows to its maximal value, $h \to \bar{s}s$ as well as $h \to \bar{b}s + \bar{s}b$ become as strong as $h \to \bar{b}b$ since $(\delta^d_{23})_{LL,RR}$ are $\mathcal{O}(1)$. In particular, one notes that $h \to \bar{b}b$ branching fraction falls to $\sim 30\%$—a completely non-SM signal testable in present [27,28] as well as future colliders. Note that, though $\Gamma(h \to \bar{d}d, \bar{s}s) \ll \Gamma(h \to \bar{b}b)$ for $\tan\beta = 20$ they are still an order of magnitude larger than the SM prediction. The above analysis can also be repeated for the up quark sector. For instance, taking $\theta^u_{ii} - \theta_g \rightsquigarrow \pi$ one finds $h_u \simeq 2.3 \times 10^{-2}\bar{h}_c$ and $h_c \simeq 1.7\bar{h}_c$ so that the most important enhancement occurs for $h \to \bar{c}c$ decay whose rate is roughly four times larger than the SM expectation. Similarly, with $(\delta^u_{23})_{LL} \simeq (\delta^u_{23})_{RR} \simeq (\delta^u_{13})_{RR} \simeq 1$ one expects $\Gamma(t \to ch) \simeq 4\Gamma(t \to uh)$ whose absolute size depends on how large $\tan\alpha$ is as follows from (26). In spite of all these estimates for Higgs decay and production rates, one keeps in mind that the MIs used above have been determined [4,23] by discarding the Higgs-exchange amplitudes which grow with $\tan\beta$. Therefore, it is after a complete analysis of various FCNC and EDM observables by including all Higgs as well as non-Higgs contributions that one can achieve an accurate determination of Higgs boson decay and production rates.

Here it must be emphasized that the discussions above neglect the radiative corrections in the Higgs sector which are, however, important and must be taken into account in an accurate analysis of the aforementioned observables because (i) at large $\tan\beta$ the radiative corrections can suppress the H_u^0–H_d^0 mixing in the Higgs mass-squared matrix so that the lightest Higgs can become effectively blind to all down type fermions [13,30], and (ii) the SUSY CP violation in the Higgs sector modifies Higgs couplings to quarks and vector bosons thereby altering the Higgs production and decay processes [13,31].

The discussions presented in the text show that the SUSY CP and flavor violation effects can have important implications for atomic EDMs, rare processes as well as the collider searches for Higgs bosons. Several effects: sizeable modifications in light quark Yukawas, filtering of SUSY CP violation into the meson mixings, enhancements and certain regularities in Higgs boson decay and production rates, ... all induce observable effects at meson factories and colliders.

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