Higgs Precision Measurements and Flavor Physics: 
A Supersymmetric Example

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Rare decays in flavor physics often suffer from Helicity suppress and Loop suppress. Helicity flip is a direct consequence of chiral $U(3)$ symmetry breaking and electroweak symmetry breaking. The identical feature is also shared by the mass generation of SM fermions. In this review, we use MSSM as an example to illustrate an explicit connection between bottom Yukawa coupling and rare decay process of $b \rightarrow s\gamma$. We take a symmetry approach to study the common symmetry breaking in supersymmetric correction to bottom quark mass generation and $b \rightarrow s\gamma$. We show that Large Peccei-Quinn symmetry breaking effect and $R$-symmetry breaking effect required by $b \rightarrow s\gamma$ inevitably lead to significant reduction of bottom Yukawa $y_b$. To compromise the reduction in $b\bar{b}$, a new decay is also needed to keep the Higgs total width as the SM value.

I. INTRODUCTION: CHIRAL AND ELECTROWEAK SYMMETRY BREAKING

A SM-like Higgs boson has been discovered at both ATLAS and CMS detectors at the CERN Large Hadron Collider first via the two cleanest channels, di-photon and four-lepton, with reconstructed invariant mass of 125 GeV[1]. The di-lepton mode was also seen with mass range consistent with the four-lepton measurement [2]. Updates from the two collaborations [3, 4] also prefer a CP-even spin-zero state $J^{PC} = 0^{++}$. The over-5 $\sigma$ evidence of $gg \rightarrow \phi \rightarrow ZZ^* \rightarrow \ell_i^+ \ell_i^- \ell_j^+ \ell_j^-$ clearly indicates that the boson $\phi$ is responsible for electroweak symmetry breaking and should be identified as the Higgs boson. In addition, both collaborations have also reported the boson decaying into tau pairs, $\phi \rightarrow \tau^+\tau^-$ which is the first evidence at the LHC that the Higgs-like boson actually couples to SM fermions. However, the final confirmation of whether the Higgs boson is the standard model (SM) Higgs boson still requires precision measurement of the Higgs couplings.
For instance, in the so-called “decoupling limit”, many new physics models beyond the SM also predict a light Higgs boson with couplings only differ from the SM ones by 10% or less. There is also example where the other couplings of this Higgs boson except the bottom Yukawa are very similar to the SM Higgs couplings while the bottom Yukawa measurement itself still suffer from large uncertainty. Higgs physics has entered an era of precision measurement and various $e^+e^-$ colliders as Higgs factory have been proposed as one intensity frontier with controlled background to improve the measurement of Higgs couplings.

On the other hand, the other type of intensity frontier as flavor factories have been playing important role in searching physics beyond the SM for many years. In this review, we try to illustrate the direct correlation between physics at two types of intensity frontier, the Higgs precision measurement and flavor physics.

Fermion mass is a consequence of chiral symmetry breaking. The Lagrangian of the kinetic energy and gauge interactions of the SM fermion fields is

$$i\bar{Q}_L^i \not{\partial} Q_L^i + i\bar{u}_R^i \not{\partial} u_R^i + i\bar{d}_R^i \not{\partial} d_R^i + i\bar{\ell}_L^i \not{\partial} \ell_L^i + i\bar{e}_R^i \not{\partial} e_R^i$$  \hspace{1cm} (1)

where index $i$ stands for generation. In Eq (1) all the fields carry unbroken gauge symmetry $U(1)_{EM}$ and some are fundamental representation of $SU(3)_C$, Majorana masses are strictly forbidden for the above fields. Lagrangian in Eq (1) is also invariant under global unitary transformations

$$f^i \rightarrow U_{ij}^f f^j, f^i \in \{Q_L^i, u_R^i, d_R^i, \ell_L^i, e_R^i\}$$  \hspace{1cm} (2)

which correspond to accidental chiral symmetries $U(3)_Q \times U(3)_u \times U(3)_d \times U(3)_\ell \times U(3)_e$ for three generations. Yukawa couplings $y_{u,d,e}$ of the SM fermions to the Higgs boson,

$$- y_{ij}^u \bar{Q}_L^i u_R^j \bar{H} - y_{ij}^d \bar{Q}_L^i d_R^j \bar{H} - y_{ij}^e \bar{\ell}_L^i e_R^j \bar{H} + h.c., \bar{H} = \epsilon H^*$$  \hspace{1cm} (3)

explicitly break the above chiral symmetries and SM fermion masses arise after $H$ developing the vacuum expectation value. Therefore, SM fermion mass generation is a consequence of both chiral symmetry breaking and electroweak symmetry breaking.

Rare decays in flavor physics often suffer from Helicity suppress and Loop suppress. For instance, pseduo-scalar leptonic decay $\pi^- \rightarrow e^- \bar{\nu}_e$ is suppressed by the electron mass. The SM contribution to $B_s \rightarrow \mu^+\mu^-$ is suppressed by the muon mass insertion. Dipole operator

$$\bar{b}\sigma_{\mu\nu} s F^{\mu\nu}$$  \hspace{1cm} (4)
which correspond to $b \to s\gamma$. Existence of on-shell spin-one photon implies that the helicity in involving quark states must be flipped in $b \to s\gamma$ process. Helicity flip also breaks chiral symmetries and electroweak gauge symmetry. In SM, the helicity flip here corresponds to a mass insertion of $m_b$. Therefore, there may exist a direct correlation between $b$-quark mass generation and $b \to s\gamma$. In this review, we use MSSM as an example to illustrate the feature as how contribution to $b \to s\gamma$ may modify the $b$ Yukawa coupling.

II. TYPE-II 2HDM AND PECCEI-QUINN SYMMETRY

We are interested in the deviation in bottom Yukawa coupling as

$$m_b = y_b v_d + \Delta m_b.$$  (5)

This is a typical feature Type-II Two-Higgs-Doublet-Model (2HDM) where quark mass generation arises from two different electroweak symmetry breaking sources.

Minimal Supersymmetric Standard Model (MSSM) is a natural Type-II 2HDM. The superpotential being holomorphic so the $\bar{H} = \epsilon H^*$ is forbidden in superpotential. The anomaly cancellation conditions for $[SU(2)_L]^2U(1)_Y$ and Witten anomaly also require the introduction of second Higgsino, the Fermionic partner of the Higgs boson, so that the Higgsino contributions to anomalies vanish. MSSM superpotential is

$$W = y_u Q u^c H_u + y_d Q d^c H_d + y_e \ell e^c H_d + \mu H_u H_d$$  (6)

The $\mu$ is a dimensional parameter which is constrained. $\mu$ cannot be zero to avoid massless Higgsino and $\mu$ cannot be too large either so that the Higgs boson do not decouple. Suppose $\mu$ arise from a dynamical field $S$ which is SM singlet

$$W \ni S H_u H_d.$$  (7)

In order to forbid the bare $\mu$-term in the superpotential, we assume there exists a non-$R U(1)_X$ symmetry under which $S$ transforms non-trivially $s \neq 0$.

$$Q u^c H_u : q + u + h_u = 0$$

$$Q d^c H_d : q + d + h_d = 0$$

$$S H_u H_d : s + h_u + h_d = 0$$  (8)
FIG. 1. Mixed QCD anomaly $A_{[SU(3)_C]^2U(1)_X}$.

If one compute the mixed QCD anomaly with $U(1)_X$ as in Fig. 1, one can obtain the anomaly coefficient

$$A_{[SU(3)_C]^2U(1)_X} = \frac{N_f}{2} (2q + u + d) = -\frac{N_f}{2} (h_u + h_d) = \frac{N_f}{2} s \neq 0$$

(9)

Therefore, a non-zero $s$ charge results in non-vanishing of mixed QCD-$U(1)_X$ anomaly $A_{[SU(3)_C]^2U(1)_X}$. The $U(1)_X$ can then be identified as Peccei-Quinn (PQ) symmetry which is a global $U(1)$ symmetry with mixed QCD anomaly. The $\mu$-term which corresponds to the Higgsino mass term explicitly breaks the PQ symmetry.

In a full MSSM, PQ symmetry breaking not only appears as Higgsino mixing $\tilde{H}_u\tilde{H}_d$ but also appear in scalar potential as

$$V \ni |F_{H_d}|^2 = y_d\mu^*\tilde{Q}\tilde{d} + y_e\mu^*\tilde{H}_u^*\tilde{\ell}\tilde{e}$$

(10)

where

$$F_{H_d} = \frac{\partial W}{\partial H_d} = y_dQd^c + y_e\ell e^c + \mu H_u.$$  

(11)

When the global $U(1)$ PQ symmetry is broken by anomaly, a pseudo-Goldstone boson is generated with its mass generated by non-perturbative QCD effect. The term in Eq. 7 would lead to Weinberg-Wilczek axion which is excluded by $K \to \pi a$ constraint. Therefore, one can introduce an extra $S^3$ term to explicitly break the $U(1)_{PQ}$ into $Z_3$ known as NMSSM approach or exotic quarks to cancel the above anomaly known as gauged $U(1)'$ approach. Coincidently, QCD invisible axion requires the PQ symmetry breaking scale is the same as the intermediate scale in gravity mediation supersymmetric theory, $M_{PQ}^2/M_{Pl} \sim M_{EW}$. The $U(1)_X$ can actually provide a simultaneous solution to strong CP problem and the $\mu$-term problem as Kim-Nilles mechanism based on supersymmetric DFSZ $H_uH_dS^2/M_{Pl}$ axion model [5].
III. R-SYMMETRY AND SUPERSYMMETRIC CORRECTIONS TO MASS GENERATION

In analogy to the chiral symmetry breaking that is associated with fermion mass generation, the Majorana gaugino mass in supersymmetric theory is associated with \( R \)-symmetry breaking. A global \( U(1) \) \( R \)-transformation is defined as a rotation over the anti-commuting coordinates (Grassmann variables) \( \theta \) and \( \bar{\theta} \)

\[
R : \theta \rightarrow e^{i\alpha} \theta, \bar{\theta} \rightarrow e^{-i\alpha} \bar{\theta}
\]

(12)

Gauge vector superfields are real so they are neutral under \( R \)-transformation. The gaugino component is then of \( R \)-charge 1 as

\[
R : \lambda \rightarrow e^{i\alpha} \lambda
\]

(13)

and gaugino mass term \( \frac{1}{2} M \lambda \lambda \) always break the \( U(1) \) \( R \)-symmetry.

One can categorize the soft-supersymmetry breaking Lagrangian based on the PQ and \( R \)-symmetries in

\[
\frac{1}{2} M \lambda \lambda + \tilde{A}_u \tilde{Q} \tilde{u}^c H_u + \ldots : \mathcal{R}
\]

\[
B_{\mu} H_u H_d \tilde{D} \tilde{Q} : \mathcal{K}
\]

\[
M_f \tilde{f}^\ast \tilde{f} : \mathcal{K}
\]

(14)

The gaugino mass and \( \tilde{A} \)-terms break \( R \)-symmetry. \( B_{\mu} \)-term breaks both \( R \) and PQ symmetries. The scalar mass term \( \tilde{f}^\ast \tilde{f} \) is trivial under any unitary transformation.

The two global \( U(1) \) assignments in MSSM are not uniquely defined. In Table III we list one sets of assignment of PQ and \( R \) charges consistent of \( SU(5) \).

| Field   | \( Q \) | \( u^c \) | \( e^c \) | \( d^c \) | \( \ell \) | \( H_u \) | \( H_d \) | \( \theta \) |
|---------|--------|--------|--------|--------|--------|--------|--------|--------|
| \( R \)-charge | \( \frac{1}{5} \) | \( \frac{1}{5} \) | \( \frac{1}{5} \) | \( \frac{7}{5} \) | \( \frac{7}{5} \) | \( \frac{8}{5} \) | \( \frac{2}{5} \) | 1 |
| PQ      | 0      | 0      | 0      | -1     | -1     | 0      | 1      | 0      |

TABLE I. Charge assignment under \( R \)-symmetry and Peccei-Quinn symmetry.

In MSSM, \( b \)-quark mass arises at the tree level from \( y_d Q d^c H_d \). The supersymmetric correction to \( m_b \) is effectively

\[
Q u^c \bar{H}_u
\]

(15)
which is known to be Lorentz invariant and gauge invariant in SM. Using charge assignments in Table III, one can substitute them into calculation of effective coupling $Qd^c H_u$ as

$$R[Qd^c H_u] : \frac{1}{5} + \frac{7}{5} - \frac{8}{5} = 0$$

$$\text{PQ}[Qd^c H_u] : 0 + (-1) + 0 = -1.$$  

(16)

(17)

Taking two Fermionic component, the $R$-invariant condition is of $R$-charge 2 while the above term is 0 so it breaks $R$-symmetry as well as the PQ symmetry. A trivial realization is that the correction can come from the two Higgs mixing term which is known as $B\mu$-term. $B\mu$-term explicitly breaks the PQ and $R$ symmetries as discussed previously. We can conclude that the supersymmetric correction to SM fermion masses must break PQ and $R$ symmetries in addition to the chiral symmetry and electroweak symmetry. Chiral symmetry breaking is quantized by the Yukawa coupling. In the case of $m_b$, the correction is also proportional to $H_u vev v_u$ which is typically dominated the electroweak symmetry breaking since it is the dominant contribution to top quark mass. Besides $y_b$ and $v_u$, the size of correction $\Delta m_b$ is then proportional to scales that break the PQ and $R$ symmetries, including $B\mu$-term, product of $\mu$-term and gaugino masses or $A$-terms.

An inapparent electroweak symmetry breaking source is the Wino-Higgsino mixing as in the neutralino mass matrix

$$N = \begin{pmatrix} M_1 & 0 & -g_1 v_d/\sqrt{2} & g_1 v_u/\sqrt{2} \\ 0 & M_2 & g_2 v_d/\sqrt{2} & -g_2 v_u/\sqrt{2} \\ -g_1 v_d/\sqrt{2} & g_2 v_d/\sqrt{2} & 0 & -\mu \\ g_1 v_u/\sqrt{2} & -g_2 v_u/\sqrt{2} & -\mu & 0 \end{pmatrix}. $$

(18)

IV. NON-DECOUPLING MSSM AN EXAMPLE

To illustrate the feature, we take a non-decoupling limit [6, 7] where supersymmetric correction to $b \rightarrow s\gamma$ is maximized to cancel the contribution from light charged Higgs. In this limit, the Higgs boson of 125 GeV is identified as the heavy neutral Higgs $H$ by taking $M_A$ around $m_Z$ scale and the charged Higgs $H^\pm$ is also around $O(100 \text{ GeV})$ as the tree level contribution to charged Higgs mass $M_{H^\pm}^2 = M_A^2 + m_{W^\pm}^2$. Such a light charged Higgs may significantly enhance the flavor violation $b \rightarrow s\gamma$. The 2HDM constraint on $b \rightarrow s\gamma$ has pushed the charged Higgs mass to be over 300 GeV. Significant cancellation to the light charged Higgs of $O(100 \text{ GeV})$ is then needed from
supersymmetric particles. As we argued, large supersymmetric correction is a consequence of large PQ symmetry breaking and large $R$ symmetry breaking so qualitatively it is easy to see where the allowed parameter region lies. At the same time, the inevitable supersymmetric correction significantly modifies the bottom Yukawa coupling. Large reduction in $H \to b\bar{b}$ then results in large reduction of Higgs total width. Since all other Higgs decay channels are at similar level as the SM predicts, new decay channel is then needed to compromise the reduction of Higgs total width. This feature has been discussed in details by our previous work [7] and we give a brief summary in this section.

In [7], all the numerical analysis are performed with *FeynHiggs 2.9.2* [10] with *HiggsBounds 3.8.0* [11] and *SUSY_Flavor 2.01* [12]. We implement the requirements as

- $M_H : 125 \pm 2$ GeV;
- $R_{\gamma\gamma} = \sigma_{\text{obs}}^{\gamma\gamma}/\sigma_{\text{SM}}^{\gamma\gamma} : 1 \sim 2$;
- Combined direct search bounds from HiggsBound3.8.0;
- $\text{BR}(B \to X_s\gamma) < 5.5 \times 10^{-4}$;
- $\text{BR}(B_s \to \mu^+\mu^-) < 6 \times 10^{-9}$.

We also calculate the constraint on $B^+ \to \tau^+\nu$ and find the destructive interference between the SM $W$ and the charged Higgs make the MSSM prediction about $20\% \sim 30\%$ smaller than the SM result of $(0.95 \pm 0.27) \times 10^{-4}$. While the experimental world average is $(1.65 \pm 0.34) \times 10^{-4}$ before 2012 [8], Belle updated their measurement at ICHEP2012 with much smaller value $0.72^{+0.29}_{-0.27} \times 10^{-4}$ for hadronic tag of $\tau$ [9]. So in the non-decoupling limit, a light charged Higgs with $\tan\beta \sim 10$ is well consistent with the new Belle measurement. In addition, the charged Higgs contribution to $B \to D^{*(\pm)}\tau\nu_\tau$ decays are not very significant in the interesting region of $M_{H^\pm}$ and $\tan\beta$. In *FeynHiggs*, Higgs boson masses are calculated to full two-loop. To illustrate the qualitative feature here, we use the leading one-loop expression with only contributions of top Yukawa couplings. Radiative corrections to the Higgs boson mass matrix elements and Higgs decay are [13-15]. Figure 2 give the allowed parameter region for the fixed choice of top squark mass as 500 GeV to enhance the correction. The points in blue region pass in addition the constraint of $\text{BR}(B \to X_s\gamma)$, while the points in black region pass all the constraints, including further the restriction of $\text{BR}(B_s \to \mu^+\mu^-)$.

It is clear that the survival region corresponds to large PQ symmetry and $R$
FIG. 2. (a) Scan Results in \([A_t, \mu] \) plane. The heavy (light) stop scenario with \( \tilde{Q}_3 = \tilde{t} = 1 \) (0.5) TeV is shown in the left (right) plot. The red region pass the direct search bounds from HiggsBounds with a heavy CP-even Higgs \( M_H = 125 \pm 2 \) GeV and an enhanced diphoton rate \( 1 < R_{\gamma\gamma} < 2 \). The blue region pass in addition the constraint of \( \text{BR}(B \to X_s \gamma) \), while the black region pass all the constraints, including further the restriction of \( \text{BR}(B_s \to \mu^+ \mu^-) \). (b) \( M_{h, H, H^\pm} \) vary with respect to \( M_A \) for \( \tilde{t} = 500 \) GeV, \( A_t = -740 \) GeV, \( \tan \beta = 11, \mu = 2300 \) GeV.

Symmetry breaking where \( \mu \sim 2 - -3 \) TeV and \( A_t \sim -750 \) GeV. Figure 2-b plots the Higgs masses respect to \( M_A \) for one set of benchmark points in the allowed parameter region.

When the supersymmetric correction in flavor physics processes is significant, bottom Yukawa also receives significant corrections while at the same time \( \tau \) Yukawa is not modified significantly. Figure 3 shows the correlation between \( \text{BR}(H \to \tau^+ \tau^-) \) normalized by its SM value and \( \text{BR}(H \to b\bar{b}) \) normalized by the corresponding SM value as well as the corresponding \( \text{BR}(t \to bH^+) \) with respect to \( M_{H^\pm} \) by assuming \( \text{BR}(H^+ \to \tau^+ \nu_\tau) = 100\% \) for the survival points. In Fig 3-a, the region where \( R_{\tau^+ \tau^-} \sim 1, H \to b\bar{b} \) partial width is much less than its SM value but a new decay of \( H \to hh \) is opened and compromises the reduction of \( b\bar{b} \) partial width to make the Higgs total width remaining unchanged. Fig 3-b clearly shows that all the parameter points that pass our selections are below the search of light charged Higgs boson via top decay \( t \to bH^+ \) with \( H^+ \to \tau^+ \nu_\tau \).
FIG. 3. (a) $\text{BR}(H \rightarrow \tau^+\tau^-)$ in correlation with $\text{BR}(H \rightarrow b\bar{b})$ normalized by the SM values respectively. (b) $\text{BR}(t \rightarrow bH^+)$ vs $M_{H^\pm}$ by assuming $\text{BR}(H^+ \rightarrow \tau^+\nu_\tau) = 100\%$. Red dots are parameter points that pass all our selection and constraints.

V. CONCLUSIONS

Rare decays in flavor physics often suffer from Helicity suppress and Loop suppress. Helicity flip is a direct consequence of chiral $U(3)$ symmetry breaking and electroweak symmetry breaking. The identical feature is also shared by the mass generation of SM fermions so one would expect that there exists a general correlation between the helicity flip flavor physics process and SM fermion mass generation. To illustrate the feature, we take a non-decoupling limit [6, 7] where supersymmetric correction to $b \rightarrow s\gamma$ is maximized to cancel the contribution from light charged Higgs. In this limit, the Higgs boson of 125 GeV is identified as the heavy neutral Higgs $H$ by taking $M_A$ around $m_Z$ scale and the charged Higgs $H^\pm$ is also around $\mathcal{O}(100$ GeV) as the tree level contribution to charged Higgs mass $M_{H^\pm}^2 = M_A^2 + m_W^2$. Such a light charged Higgs may significantly enhance the flavor violation $b \rightarrow s\gamma$. The 2HDM constraint on $b \rightarrow s\gamma$ has pushed the charged Higgs mass to be over 300 GeV. Significant cancellation to the light charged Higgs of $\mathcal{O}(100$ GeV) is then needed from supersymmetric particles. At the same time, the inevitable supersymmetric correction significantly modifies the bottom Yukawa coupling. Large reduction in $H \rightarrow b\bar{b}$ then results in large reduction of Higgs total width. Since all other Higgs decay channels are at similar level as the SM predicts, new decay channel $H \rightarrow hh$ is then needed to compromise the reduction of Higgs total width.
A more general approach to study correlation between Yukawa couplings and helicity flipped flavor violation operators is to appear in [16].

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