SUSY induced top quark FCNC decay t→cht→cht → ch after Run I of LHC postprint

Authors: Cao,J, Han,C, Wu,L, Yang,JM, Zhang,M

Date: 2016-12-28T00:00:00+00:00

Abstract

In light of the Higgs discovery and the nonobservation of sparticles at the LHC, we revisit the supersymmetric theory (SUSY) induced top quark flavor-changing decay into the Higgs boson. We perform a scan over the relevant SUSY parameter space by consider

Full Text

Preamble

SUSY Induced Top Quark FCNC Decay t → ch After Run I of LHC

Junjie Cao[1],[2], Chengcheng Han[3], Lei Wu , Jin Min Yang[3], Mengchao Zhang[3]

[1]Institute of Theoretical Physics, School of Physics, Xiamen University, Xiamen 361005, China
[2]State Key Laboratory of Theoretical Physics, Institute of Theoretical Physics, Academia Sinica, Beijing 100190, China
[3]School of Physics, Xiamen University, Xiamen 361005, China
ARC Center of Excellence for Particle Physics at the Terascale, School of Physics, The University of Sydney, NSW 2006, Australia

Abstract

In light of the Higgs discovery and the nonobservation of sparticles at the LHC, we revisit the SUSY-induced top quark flavor-changing decay into the Higgs boson. We perform a scan over the relevant SUSY parameter space by considering the constraints from the Higgs mass measurement, the LHC search for SUSY, vacuum stability, precision electroweak observables, as well as $B \to X_s\gamma$. We have the following observations: (1) In the MSSM, the branching ratio of $t \to ch$

can only reach 3.0×10^{-6}, which is about one order of magnitude smaller than previous results obtained before the advent of the LHC. Among the considered constraints, the Higgs mass and the LHC search for sparticles are found to play an important role in limiting the prediction. (2) In the singlet extension of the MSSM, since the squark sector is less constrained by the Higgs mass, the branching ratio of $t \to ch$ can reach the order of 10^{-5} in the allowed parameter space. (3) The chiral-conserving mixings δ_{LL} and δ_{RR} may have remnant effects on $t \to ch$ in the heavy SUSY limit. In the MSSM with squarks and gluino above 3 TeV and meanwhile the CP-odd Higgs boson mass around 1 TeV, the branching ratio of $t \to ch$ can still reach the order of 10^{-8} under the constraints.

Introduction

A scalar with mass around 125 GeV has been discovered at the LHC [**?**, **?**]. According to the analysis of the ATLAS and CMS collaborations, the measured properties of this scalar, albeit with large experimental uncertainties, agree well with those of the Higgs boson in the Standard Model (SM), which means that it plays a role in electroweak (EW) symmetry breaking and also in mass generation for the fermions in the SM [**?**, **?**, **?**]. Even so, due to the deficiencies of the SM itself in describing the symmetry breaking, it is well motivated to interpret this scalar in various frameworks of new physics. Obviously, in order to unambiguously decipher the nature of the scalar, it is mandatory to scrutinize both experimentally and theoretically the couplings of the scalar, including its self-interactions.

In this direction, the couplings of the scalar with the yet-known heaviest particle, the top quark, are of fundamental importance since, as suggested by the LHC Higgs data, the $h\bar{t}t$ coupling is strong, and meanwhile it is widely conjectured to be sensitive to new physics. In fact, great efforts have been paid recently to investigate the top-Higgs associated production processes like $pp \to t\bar{t}h$ [**?**, **?**] and $pp \to qth$ [**?**] at the LHC to extract the size and sign of the $h\bar{t}t$ Yukawa coupling, and also the top quark flavor-changing decay $t \to ch$ to probe anomalous top-Higgs interaction [**?**, **?**].

Among the new physics models, supersymmetric theory (SUSY) is a promising one due to its capability to solve the hierarchy problem of the SM, unify the gauge couplings, as well as provide a viable Dark Matter candidate [**?**, **?**]. In SUSY, a SM-like Higgs boson h around 125 GeV usually implies third-generation squarks at or heavier than 1 TeV, and the preference for heavy squarks is further corroborated by the absence of any signal in the search for SUSY at the LHC. If the SUSY scale is really high, which was the focus of many recent theoretical works [**?**], the only way to detect SUSY is through its possibly large remnant effects in EW processes. Such effects may exist in the Higgs process because the dominant part of the Higgs couplings to squarks is proportional to soft SUSY breaking parameters [**?**], and consequently, the suppression induced by the squark propagators in SUSY radiative corrections to the process may be compensated under certain conditions.

This feature has been demonstrated in the SUSY correction to the $h\bar{b}b$ vertex [?], the Higgs pair production process at the LHC [?], and also the Higgs rare decay $h \to \tau\bar{\mu}$ [?]. Here we emphasize that the existence of the remnant effect in the asymptotic large SUSY mass limit does not contradict the Appelquist-Carazzone theorem [?], which is valid only for supersymmetric theories with an exact gauge symmetry. Previous studies on such remnant effects in SUSY can also be found in [?].

In this work, we focus on the top quark flavor-changing decay $t \to ch$ in SUSY. The reasons for our interest mainly come from three considerations. Firstly, the LHC as a top factory has great capability to scrutinize the properties of the top quark, including its rare decay modes. As far as the flavor-changing decay $t \to ch$ is concerned, its branching ratio in the SM is only at the order of 10^{-14} [?], while in SUSY it may be greatly enhanced to 10^{-4} according to previous studies [?, ?]. Since any observation of this decay in the future will be robust evidence of new physics, this decay should receive attention in the LHC era, especially noting the fact that the Higgs boson has been recently discovered. Secondly, as introduced before, the LHC experiment has measured the Higgs mass and pushed SUSY to a rather high scale. These results have great impacts on the SUSY prediction for $t \to ch$, so it is necessary to update previous studies on $t \to ch$ in light of the experimental progress.

Thirdly, unlike other top FCNC processes in SUSY [?], the decay $t \to ch$ may have remnant effects in the heavy SUSY limit. From a theoretical point of view, it is worthwhile to investigate such a feature in detail. Besides, we remind that, if flavor mixings between scharm and stop are present, which may push up the rate of $t \to ch$ greatly, the LHC constraint on stop masses can be relaxed. This in return can alleviate the fine-tuning problem of SUSY [?].

This paper is organized as follows. In Section II, we parameterize the flavor mixings in the squark sector and define our conventions. We also list various constraints on SUSY. In Section III, we study the decay $t \to ch$ in both low-energy SUSY and heavy SUSY, and present some benchmark points at which the predictions for $t \to ch$ are optimized. We also exhibit the features of the remnant effect on $t \to ch$. Finally, we present our conclusions in Section IV.

II. FCNC Interactions in SUSY

In supersymmetric theories such as the Minimal Supersymmetric Standard Model (MSSM) [?] and the Next-to-Minimal Supersymmetric Standard Model (NMSSM) [?], the squark sector consists of six up-type squarks $(\tilde{u}_L, \tilde{c}_L, \tilde{t}_L, \tilde{u}_R, \tilde{c}_R, \tilde{t}_R)$ and six down-type squarks $(\tilde{d}_L, \tilde{s}_L, \tilde{b}_L, \tilde{d}_R, \tilde{s}_R, \tilde{b}_R)$. In general, states with different chiral and flavor quantum numbers in each type of squark will mix to form mass eigenstates, and consequently potentially large flavor-changing interactions arise from the misalignment between the rotations that diagonalize the quark and squark sectors. In the super-CKM basis, the 6×6 squark mass matrix $M^2_{\tilde{q}}$ ($\tilde{q} = \tilde{u}, \tilde{d}$) takes the form [?]

$$M^2_{\tilde{q}} = \begin{pmatrix} M^2_{\tilde{q},LL} + C^{LL}_{\tilde{q}} & M^2_{\tilde{q},LR} - C^{LR}_{\tilde{q}} \\ (M^2_{\tilde{q},LR} - C^{LR}_{\tilde{q}})^\dagger & M^2_{\tilde{q},RR} + C^{RR}_{\tilde{q}} \end{pmatrix},$$

where $M^2_{\tilde{q},LL} = m^2_{\tilde{q}} + m^2_q + \cos 2\beta M^2_Z (T^3_q - Q_q s^2_W)\hat{1}$ and $C^{LL}_{\tilde{q}} = m_q \mu (\tan\beta)^{-2T_q}\hat{1}$. Here $m_{\tilde{q}}$ and m_q are 3×3 diagonal matrices with $\hat{1}$ standing for the unit matrix in flavor space, m_q being the diagonal quark mass matrix and $T^3_q = \frac{1}{2}, -\frac{1}{2}$ for $q = u, d$ respectively, and $\tan\beta = v_2/v_1$ is the ratio of the vacuum expectation values of the SU(2) doublet Higgs fields. Similarly, $M^2_{\tilde{q},RR} = m^2_{\tilde{q}} + m^2_q + \cos 2\beta M^2_Z Q_q s^2_W \hat{1}$ and $C^{RR}_{\tilde{q}} = m_q \mu (\tan\beta)^{-2T_q}\hat{1}$.

If one only considers the flavor mixings between the second and third generation squarks, the soft breaking squared masses $(M^2_{\tilde{q}})_{LL}$, $(M^2_{\tilde{q}})_{LR}$ and $(M^2_{\tilde{q}})_{RR}$ can be parameterized as

$$(M^2_{\tilde{u}})_{LL} = \begin{pmatrix} M^2_{Q1} & 0 & 0 \\ 0 & M^2_{Q2} & \delta_{LL} M_{Q2} M_{Q3} \\ 0 & \delta_{LL} M_{Q2} M_{Q3} & M^2_{Q3} \end{pmatrix},$$

$$(M^2_{\tilde{u}})_{LR} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & \delta_{LR} v_2 M^{RL}_U \\ 0 & \delta_{RL} v_2 M^{RL}_U & m_t A_t \end{pmatrix},$$

$$(M^2_{\tilde{u}})_{RR} = (M^2_{\tilde{u}})_{LL}|_{M_{Qi} \to M_{Ui}, \delta_{LL} \to \delta_{RR}},$$

where M_{Qi} and M_{Ui} ($i = 1, 2, 3$ denotes generation index) are soft breaking parameters with mass dimension, M^{RL}_U represents the SUSY scale defined as $M^{RL}_U = (M_{U3} + M_{Q2})/2$ and $M^{LR}_U = (M_{U2} + M_{Q3})/2$, and δ_{LL}, δ_{LR}, δ_{RL} and δ_{RR} reflect the extent of flavor violation. Similarly, for down-squarks we have

$$(M^2_{\tilde{d}})_{LR} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & \delta^d_{LR} v_1 M^{RL}_D \\ 0 & \delta^d_{RL} v_1 M^{RL}_D & m_b A_b \end{pmatrix},$$

$$(M^2_{\tilde{d}})_{RR} = (M^2_{\tilde{u}})_{LL}|_{M_{Qi} \to M_{Di}, \delta_{LL} \to \delta^d_{RR}},$$

and due to SU(2) gauge symmetry, $(M^2_{\tilde{d}})_{LL}$ is determined by [?]

$$(M^2_{\tilde{d}})_{LL} = V^\dagger_{\text{CKM}} (M^2_{\tilde{u}})_{LL} V_{\text{CKM}},$$

with V_{CKM} denoting the Cabibbo-Kobayashi-Maskawa matrix in the SM. Note that in these equations, we only keep the chiral-flipping terms for third-family squarks because these terms are usually assumed to be proportional to corresponding quark masses and can only be non-negligible for third-family squarks.

The squark mass eigenstates can be obtained by diagonalizing the mass matrix presented above with a unitary rotation $U_{\tilde{q}}$, which is performed numerically in our analysis. The interaction of the field X with a pair of squark mass eigenstates is then obtained by

$$V(X\tilde{q}^*_\alpha \tilde{q}'_\beta) = U^\dagger_{\tilde{q},\alpha,i} U_{\tilde{q}',j,\beta} V(X\tilde{q}^*_i \tilde{q}'_j),$$

where $V(X\tilde{q}^*_i \tilde{q}'_j)$ denotes a generic vertex in the interaction basis and $V(X\tilde{q}^*_\alpha \tilde{q}'_\beta)$ is the vertex in the mass-eigenstate basis. It is clear that both the squark masses and their interactions depend on the mixing parameters δ_i.

In some fundamental supersymmetric theories like mSUGRA and gauge-mediated SUSY-breaking models, the mixing parameters are functions of the soft breaking masses and usually exhibit certain hierarchical structure [**?**]. In this work, in order to make our discussion as general as possible, we treat all δ_i as free parameters and limit them by physical observables. The constraints we consider include:

(I) The recently measured SM-like Higgs boson mass m_h. In the MSSM, this mass is determined by the renormalized self-energies of the doublet CP-even Higgs fields, h_u and h_d, and the transition between them. Squarks contribute to these quantities through the $\tilde{q}^*\tilde{q}S$ and $\tilde{q}^*\tilde{q}SS$ interactions with S denoting either h_u or h_d [**?**, **?**]. In the presence of flavor mixings, both the interactions and the squark masses may be quite different from those in the case of $\delta_i = 0$, and so is the SM-like Higgs boson mass. Among the flavor mixing parameters δ_i, the Higgs mass is more sensitive to the chiral-flipping ones δ_{LR} and δ_{RL}. In this work, we obtain m_h in the MSSM using the code FeynHiggs [**?**]. In our scan over the parameter space of the low-energy MSSM, we require the mass to be within about 2 GeV of its measured central value, i.e., $123\,\text{GeV} \leq m_h \leq 127\,\text{GeV}$. While for the MSSM in the heavy SUSY case (see below), noting that the mass obtained by FeynHiggs suffers from potentially large theoretical uncertainties, we require a moderately wider range, i.e., $121\,\text{GeV} \leq m_h \leq 129\,\text{GeV}$.

(II) The LHC search for SUSY. By now both the ATLAS and CMS collaborations have made great efforts in searching for signals of gluinos, squarks, as well as charginos and neutralinos, and based on certain assumptions, they exclude some SUSY particles up to about 1 TeV [**?**]. These obtained results, however, cannot be applied directly to a general SUSY case, and in order to implement the LHC constraints, one has to perform detailed Monte Carlo simulation for each SUSY parameter point with the same strategies as those of the collaborations, then compare the simulated results with the LHC data [**?**]. In practice, such a process is rather time-consuming and cannot be applied to an extensive scan over the SUSY parameter space, where a large number of samples are involved. In order to simplify our analysis, we note that by now the gluino is preferred to be at the TeV scale without considering special cases such as compressed SUSY spectra [**?**], while the second and third generation squarks may still be as light as several hundred GeV [**?**, **?**], especially in the presence of flavor mixing when the limitation on the squark spectrum can be further relaxed [?]. So we make the following assumption in our discussion:

$$m_{\tilde{q}_\alpha}, m_{U3} \geq 200\,\text{GeV}, \quad m_{\tilde{g}} \geq 1\,\text{TeV}, \quad m_{Q2}, m_{U2}, m_{Q3} \geq 500\,\text{GeV}.$$

As will be shown below, our conclusions are not sensitive to such assumptions.

(III) The metastability of the vacuum state. This constraint reflects the fact that squarks as scalar fields contribute to the SUSY potential, and consequently their soft breaking parameters should be limited by the stability (or more general metastability) of the vacuum state [?, ?]. Assuming only δ_{LR} or δ_{RL} contributes to the potential, the metastability requires [?, ?]

$$|\delta_{LR}| \lesssim 1.2\times\sqrt{\frac{M_A^2\cos^2\beta + m_{\tilde{q}}^2}{m_{\tilde{q}}^2}}, \quad |\delta_{RL}| \lesssim 1.2\times\sqrt{\frac{M_A^2\cos^2\beta + m_{\tilde{q}}^2}{m_{\tilde{q}}^2}}, \quad |A_t| \lesssim 2.67\sqrt{M_A^2\cos^2\beta + m_{\tilde{q}}^2}.$$

Note that in previous studies on top quark flavor-changing neutral current processes, this constraint was usually missed.

(IV) EW precision observables M_W and $\sin^2\theta_{\text{eff}}$. In SUSY, the corrections to M_W and $\sin^2\theta_{\text{eff}}$ are dominated by squark loops, and their sizes reflect the mass disparity of left-handed SU(2) doublet squarks. To a good approximation, these two quantities are related with the $\delta\rho$ parameter by

$$\delta M_W \approx \frac{M_W}{2}\frac{c_W^2}{c_W^2 - s_W^2}\delta\rho, \quad \delta\sin^2\theta_{\text{eff}} \approx -\frac{c_W^2 s_W^2}{c_W^2 - s_W^2}\delta\rho,$$

where $\delta\rho = \frac{\Sigma_Z(0)}{M_Z^2} - \frac{\Sigma_W(0)}{M_W^2}$. In this work, we repeat our previous calculation of δM_W and $\delta\sin^2\theta_{\text{eff}}$ in [?] where three-generation squarks are considered to implement the SU(2) relation between $(M_{\tilde{u}}^2)_{LL}$ and $(M_{\tilde{d}}^2)_{LL}$, and require $\delta M_W \leq 21\,\text{MeV}$, $\delta\sin^2\theta_{\text{eff}} \leq 19.6\times10^{-5}$, which are their allowed ranges at 2σ level after considering experimental and theoretical uncertainties [?].

(V) Constraint from $B \to X_s\gamma$. In the MSSM, the SUSY contributions to $B \to X_s\gamma$ come from four kinds of loops mediated by charged Higgs bosons, charginos, neutralinos, and gluinos respectively. We calculate these contributions using the code FeynHiggs [?], and require $3.04\times10^{-4} \leq \text{Br}(B \to X_s\gamma) \leq 4.02\times10^{-4}$ (corresponding to its 2σ allowed range by experiments [?]) in our parameter scan. Since the neutralino contribution is usually small, the expression for $B \to X_s\gamma$ in the NMSSM is roughly identical to that of the MSSM.

Regarding the above constraints, it should be noted that constraints (I), (III), and (IV) do not diminish as the SUSY scale becomes higher; that is, they are not decoupled in the heavy SUSY limit, while the process $B \to X_s\gamma$ does not possess such a property.

[Figure 1: see original paper]: Feynman diagrams of the SUSY-QCD contribution to $t \to ch$. If the charm quark mass is neglected, the contribution from diagram (c) vanishes.

III. SUSY Prediction on the Rate of $t \to ch$

A. $t \to ch$ in Low-Energy SUSY

In the MSSM, the dominant contribution to the process $t \to ch$ arises from the SUSY-QCD diagrams shown in Fig.~1. The relevant Lagrangian is given by [?]

$$\mathcal{L} \supset \sqrt{2} g_s [\bar{\tilde{g}}^a (-U^*_{2\alpha} P_L + U^*_{5\alpha} P_R) \tilde{q}^*_{\alpha i} T^a_{ij} c_j + \bar{\tilde{g}}^a (-U^*_{3\alpha} P_L + U^*_{6\alpha} P_R) \tilde{q}^*_{\alpha i} T^a_{ij} t_j] + \text{h.c.} + \sum_{\alpha,\beta=1}^{6} C_{\alpha\beta} \tilde{q}^*_\alpha \tilde{q}_\beta h + Y_t \bar{t} t h,$$

where $\tilde{g}$ and $\tilde{q}_\alpha$ denote gluino and squark in the mass eigenstate respectively, T^a is the Gell-Mann matrix with i and j representing color indices, $U_{\tilde{q}}$ is the 6×6 rotation matrix to diagonalize the mass matrix for up-type squarks, $C_{\alpha\beta}$ parameterizes the coupling of the Higgs boson with squark mass eigenstates $\tilde{q}_\alpha$ and $\tilde{q}_\beta$, and $Y_t = m_t \cos\alpha / \sin\beta$ with α being the rotation angle to diagonalize the CP-even Higgs mass matrix.

The amplitude of $t \to ch$ can then be expressed by

$$i\mathcal{M}(t \to ch) = \bar{u}(p_c)[(F_{1L} + F_{2L}) P_L + (F_{1R} + F_{2R}) P_R] u(p_t),$$

where, after neglecting the charm quark mass, F_i are given by

$$F_{1L} = \frac{\alpha_s}{6\pi^2} \sum_{\alpha,\beta=1}^{6} C_{\alpha\beta} \left[U^*_{3\alpha} U_{5\beta} m_{\tilde{g}} C_0(p_c^2, p_h^2, p_t^2, m^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}, m^2_{\tilde{q}_\beta}) - U^*_{6\alpha} U_{5\beta} m_t C_{12}(p_c^2, p_h^2, p_t^2, m^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}, m^2_{\tilde{q}_\beta}) \right],$$

$$F_{1R} = \frac{\alpha_s}{6\pi^2} \sum_{\alpha,\beta=1}^{6} C_{\alpha\beta} \left[U^*_{6\alpha} U_{2\beta} m_{\tilde{g}} C_0(p_c^2, p_h^2, p_t^2, m^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}, m^2_{\tilde{q}_\beta}) - U^*_{3\alpha} U_{2\beta} m_t C_{12}(p_c^2, p_h^2, p_t^2, m^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}, m^2_{\tilde{q}_\beta}) \right],$$

$$F_{2L} = -\frac{Y_t}{6\pi^2 m_t} \sum_{\alpha=1}^{6} U^*_{3\alpha} U_{5\alpha} B_0(p_c^2, m^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}), \quad F_{2R} = -\frac{Y_t}{6\pi^2 m_t} \sum_{\alpha=1}^{6} U^*_{6\alpha} U_{2\alpha} B_0(p_c^2, M^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}).$$

In the above expressions, p_t, p_c and p_h denote the momenta of the top quark, charm quark, and Higgs boson respectively, $m_{\tilde{g}}$ and $m_{\tilde{q}_\alpha}$ represent the masses of the gluino and squark respectively, and B_0, C_0 and C_{12} are the standard two-point and three-point loop functions respectively [?]. In the heavy SUSY case discussed below, since the involved sparticle masses are much larger than m_t, the contribution from C_{12} can be safely ignored, and B_0, C_0 can be approximated by

$$C_0(p_c^2, p_h^2, p_t^2, m^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}, m^2_{\tilde{q}_\beta}) = \frac{1}{m^2_{\tilde{g}}} \left[1 + \frac{1-\delta_\beta}{\delta_\alpha - \delta_\beta} \ln \delta_\alpha + \frac{1-\delta_\alpha}{\delta_\beta - \delta_\alpha} \ln \delta_\beta \right] + \mathcal{O}\left(\frac{m_t^2}{m^2_{\tilde{g}}} \right),$$

$$B_0(p_c^2, m^2_{\tilde{g}}, m^2_{\tilde{q}_\alpha}) = \frac{1}{m^2_{\tilde{g}}} [1 - \ln \delta_\alpha] + \mathcal{O}\left(\frac{m_t^2}{m^2_{\tilde{g}}} \right),$$

where δ_α is defined as $\delta_\alpha = m^2_{\tilde{q}_\alpha} / m^2_{\tilde{g}}$.

As shown in [?], the SUSY-QCD contribution to $t \to ch$ in low-energy MSSM has the following features: - In the case that only one flavor mixing parameter δ_i is non-zero, the rate of $t \to ch$ increases monotonically as δ_i becomes larger, while if several non-vanishing δ_i coexist, their effects may cancel each other out. - Since the effective $h\bar{c}t$ vertex involves both chiral-flipping and flavor-changing, the chiral-conserving parameter δ_{LL}/δ_{RR} must be accompanied by chiral-flipping $h\tilde{t}^*\tilde{t}_L$ interaction in contributing to the vertex, while the chiral-flipping parameter δ_{LR}/δ_{RL} alone is able to lead to the $h\bar{c}t$ vertex. As a result, the effective vertex is usually more sensitive to δ_{LR} and δ_{RL} if we do not consider the constraints on the mixing parameters. - Unlike other top quark FCNC processes, the Super-GIM mechanism does not apply to the decay $t \to ch$. But since there exists a strong cancellation between diagrams (a) and (b) in Fig.~1 (see discussion below), the rate of $t \to ch$ depends on the soft mass parameters in a complex way.

In the following, we do not intend to exhibit these features, but instead, noting that the constraints (I-V) introduced above were not considered before, we try to figure out the order of Br($t \to ch$) that SUSY can predict after considering these constraints. For this purpose, we perform two independent scans over relevant SUSY parameters by imposing the constraints. Details of our scans are as follows:

- **Scan-I**: We restrict our discussion to the MSSM and calculate the Higgs mass with the code FeynHiggs. The parameter region we explore is given by

$$
\begin{aligned}
&500\,\text{GeV} \le m_{Q2}, m_{Q3}, m_{U2}, m_{D2}, m_{D3} \le 2\,\text{TeV}, \\
&200\,\text{GeV} \le m_{U3} \le 2\,\text{TeV}, \quad |A_t| \le 6\sqrt{m_{Q3} m_{U3}}, \\
&1\,\text{TeV} \le m_{\tilde{g}} \le 2\,\text{TeV}, \quad 1 \le \tan\beta \le 40, \\
&-1 \le \delta_{LL}, \delta_{RR} \le 1, \quad -2.0 \le \delta_{LR}, \delta_{RL} \le 2.0, \\
&-0.5 \le \delta^d_{LR}, \delta^d_{RL} \le 0.5, \quad 400\,\text{GeV} \le m_A \le 2\,\text{TeV}, \\
&-2\,\text{TeV} \le \mu \le 2\,\text{TeV}.
\end{aligned}
$$

  In drawing up the strategy of this scan, we note that although the down-type squark parameters like δ^d_{LR} and M_D do not affect the rate of $t \to ch$, they are needed in the $\delta\rho$ and $B \to X_s\gamma$ calculations. So to make our conclusions as general as possible we vary them in reasonable regions. We also note that since too many parameters are involved in the scan, the traditional random scan method is not efficient in searching for the maximal value of Br($t \to ch$). So we adopt the Markov chain method for this task. During the scan we adjust the optimal value by the results obtained from previous samplings until it reaches stable values.

- **Scan-II**: Same as Scan-I, but in order to relax the Higgs mass bound, we go beyond the MSSM by considering extra contributions to the mass. To be more specific, now we write the Higgs mass as $m_h^2 = m^2_{h,\text{MSSM}} + \lambda^2 v^2 \sin^2 2\beta$ with λ being a free parameter in the range from 0 to 0.7.

This treatment of the Higgs mass is motivated by singlet extensions of the MSSM such as the NMSSM, where the interaction between the singlet Higgs field and the doublet Higgs fields in the MSSM provides an additional contribution to the mass at tree level [?]. In order to get a large contribution from this singlet extension, we set $\tan\beta = 1.5$.

In Table I, we show two benchmark points for each scan. These points correspond to very optimal cases in predicting a large rate of $t \to ch$. After analyzing our scan results, we have the following observations:

- A small $m_{\tilde{g}}$ (around its experimental lower bound) seems to be favored to maximize the rate of $t \to ch$ after considering the constraints. Meanwhile, it is interesting to learn from Table I that, although we allow m_{U3} to be as low as 200 GeV, the optimal points do not correspond to low m_{U3} values.
- In low-energy SUSY, the size of $\mathrm{Br}(t \to ch)$ for the optimal points in Table I is not sensitive to the parameters m_A and μ. For example, our results indicate that shifting m_A from its value in Table I by 100 GeV only results in a change of $\mathrm{Br}(t \to ch)$ by less than 1%. While as will be shown below, these two parameters play an important role in determining the rate of the rare decay in the heavy SUSY limit.
- Among the considered constraints, the most stringent ones come from the Higgs mass and the LHC search for SUSY. As a result, the branching ratio of $t \to ch$ can only reach 10^{-6} in low-energy MSSM, which is about one order of magnitude smaller than previous predictions obtained before the advent of the LHC [?]. In Scan-II, however, since the Higgs mass constraint on squark masses is comparatively relaxed, $\mathrm{Br}(t \to ch)$ can reach 10^{-5}.

Regarding our results on $t \to ch$, we remind that we do not include the SUSY-EW contribution to $t \to ch$. The reason is that the amplitude of the SUSY-EW contribution is roughly determined by $\alpha m_{\tilde{\chi}^\pm}/\mathrm{Max}(m^2_{\tilde{q}}, m^2_{\tilde{\chi}^\pm})$ from naive estimation [?], while the SUSY-QCD contribution is determined by $\alpha_s m_{\tilde{g}}/\mathrm{Max}(m^2_{\tilde{q}}, m^2_{\tilde{g}})$. Noting that $m_{\tilde{q}_D}$ is not much smaller than $m_{\tilde{g}}$ as suggested by the LHC search for the second and third generation squarks [?, ?], and also that the flavor mixings in the down-type squark sector are more tightly constrained by B-physics than those in the up-type squark sector, we conclude that the SUSY-EW contribution should not be comparable with the SUSY-QCD contribution. So our estimates on the magnitude of $t \to ch$ will not change after including the SUSY-EW contribution. Another reason to neglect the EW contribution is that once it is considered, too many parameters will be involved, but meanwhile this does not change our conclusion.

Before we end this subsection, we have two comments. One is that in extensions of the MSSM, the decay chain of a certain sparticle may be quite different from its MSSM prediction, and the analysis of the ATLAS and CMS collaborations in searching for SUSY may become irrelevant [?]. As a result, the constraint on the sparticle mass may be relaxed. This in return may push up the SUSY prediction on the rate of $t \to ch$. Moreover, in extensions of the MSSM the couplings of the Higgs boson with quarks and squarks may be slightly changed. The influence of such changes on the rate of $t \to ch$ is usually not as significant as the relaxation of the Higgs mass constraint. The other comment is that, given $\mathrm{Br}(t \to ch) \sim 10^{-5}$, it is difficult to detect such a top quark rare decay at the 14-TeV LHC with an integrated luminosity of 100 fb^{-1} [**?**], but at future linear colliders such as TLEP [**?**], detection of the decay is still possible.

B. Remnant Effect of $t \to ch$ in Heavy SUSY Limit

As mentioned in Section I, because the $h\tilde{q}^*\tilde{q}$ coupling strength is mainly determined by soft SUSY breaking parameters, the SUSY-QCD contribution to the effective $h\bar{c}t$ interaction may exhibit remnant effects of SUSY in the heavy SUSY limit. In order to investigate this issue, we assume a common SUSY mass scale $m_{Q2} = m_{Q3} = m_{U2} = m_{U3} = m_{\tilde{g}} = M_{\mathrm{SUSY}}$, and study the dependence of $\mathrm{Br}(t \to ch)$ on M_{SUSY} for different choices of A_t and μ in Fig.~2. In obtaining Fig.~2, we consider the case that only δ_{LL} is non-vanishing and fix $\delta_{LL} = 0.7$, $\tan\beta = 4, 10$ and $m_A = 400, 800$ GeV. We set $A_t = M_{\mathrm{SUSY}}$ and $\mu = 0$ for the left panel, and $A_t = 0$ and $\mu = M_{\mathrm{SUSY}}$ for the right panel. These settings are only for exhibiting the decoupling behavior of $\mathrm{Br}(t \to ch)$ and we do not consider the constraint from the LEP search for charginos, which requires $\mu \gtrsim 103$ GeV. Fig.~2 then indicates that SUSY has remnant effects on the rare decay rate only when μ is at the SUSY scale and meanwhile m_A is at the weak scale, and the size of the effect depends strongly on m_A and $\tan\beta$, e.g., small values of m_A and $\tan\beta$ tend to enhance the effect. We also investigate the case that only δ_{LR} and δ_{RL} are non-vanishing, and we do not find such remnant effects in the heavy SUSY limit for any choices of μ and m_A.

The behaviors shown in Fig.~2 can be understood by the effective Lagrangian that describes the Higgs and quark system. After including loop effects, one can write down the Lagrangian as follows:

$$\mathcal{L}_{\mathrm{eff}} = \sum_{i,j=1}^{3} \left\{ \bar{q}'_i (m'_{ij} + \delta m'_{ij}) P_L q'_j + h \bar{q}'_i (Y'_{ij} + \delta Y'_{ij}) P_L q'_j \right\} + \mathrm{h.c.}$$

$$= \sum_{i,j=1}^{3} \left\{ (\bar{q} V_R^\dagger)_i (m'_{ij} + \delta m'_{ij}) P_L (V_L q)_j + h (\bar{q} V_R^\dagger)_i (Y'_{ij} + \delta Y'_{ij}) P_L (V_L q)_j \right\} + \mathrm{h.c.}$$

$$= \sum_{i=1}^{3} \bar{q}_i m_i P_L q_i + \sum_{i,j=1}^{3} h \bar{q}_i \left[V_R^\dagger \left(\delta Y' - \frac{\delta m'}{v} \right) V_L \right]_{ij} P_L q_j + \mathrm{h.c.},$$

where q'_i and m'_{ij} are the quark field and its mass matrix at tree level with i, j denoting flavor indices, and $\delta m'$ and $\delta Y'$ represent loop corrections to the mass matrix and the Yukawa coupling respectively. The second equation reflects the definition of quark mass eigenstates with V_L and V_R denoting the rotation matrices for left-handed quarks and right-handed quarks respectively. After such a definition, the loop-corrected mass matrix $m' + \delta m'$ is diagonal with its diagonal element m_i representing the physical quark mass determined by experiments. At this stage, the correction to the $h\bar{q}_i q_j$ interaction is given by $h\bar{q}_i[V_R^\dagger(\delta Y' - \delta m'/v)V_L]_{ij}P_L q_j + \text{h.c.}$. Obviously, if $\delta Y' = \delta m'/v$, the new physics contribution to the $h\bar{q}_i q_j$ interaction vanishes. In actual calculation, $\delta m'$ is obtained from $q_i - q_j$ transition diagrams like diagram (b) of Fig.~1 without the emission of the Higgs particle, and $\delta Y'$ comes from the vertex correction like diagram (a) of Fig.~1. The effective Lagrangian then indicates that the two contributions should cancel each other in contributing to the $h\bar{q}_i q_j$ interaction.

As far as the SUSY-QCD correction to the $h\bar{q}_i q_j$ interaction is concerned, its behavior in the heavy SUSY limit can be analyzed with the mass insertion approximation. In this method, squark masses are taken to be the diagonal elements of the squark mass matrix, and the non-diagonal elements are treated as interactions. In order to illustrate this method in explaining the remnant effect, we first consider the well-studied SUSY-QCD correction to the $h\bar{b}b$ vertex [?]. In this example, to get δm_b one needs to insert the $\tilde{b}_L - \tilde{b}_R$ transition an odd number of times into the sbottom propagator entering bottom quark self-energy diagrams, and sum the corresponding contribution to infinite orders of the insertion. One can check that, with one more insertion, the corresponding contribution is suppressed by a factor $m_b(A_b - \mu\tan\beta)/M_{\text{SUSY}}^2$ compared with that without the insertion, and only for the first-order insertion is δm_b not suppressed by M_{SUSY}, which means that the corresponding contribution is non-decoupled in the heavy SUSY limit. In a similar way, one can check that an even number of insertions are needed to get the expression of δY_b in calculating the $h\bar{b}b$ vertex correction, and only for the zero-th insertion is the contribution non-decoupled. Putting these two contributions together, one can learn that the non-decoupled terms proportional to A_b are exactly canceled out, and the remaining contribution is proportional to the well-known form $\mu m_{\tilde{g}}/M_{\text{SUSY}}^2 \tan\beta$ [?]. This effect, albeit scaling as $1/m_{\tilde{g}}^2$, does not diminish as $\mu \sim m_{\tilde{g}} \sim M_{\text{SUSY}} \to \infty$ for m_A at the weak scale, and is therefore dubbed as the remnant effect of SUSY.

Next we turn to analyze the SUSY-QCD contribution to the $h\bar{c}t$ vertex. Since such an interaction involves both chiral-flipping and flavor-changing, appropriate insertions are needed to accomplish both tasks. For the chiral-flipping mixings δ_{LR} and δ_{RL}, their role is quite similar to A_b in the SUSY-QCD correction to the $h\bar{b}b$ coupling, and their non-decoupling contribution is completely canceled out. While for the chiral-conserving mixings δ_{LL} and δ_{RR}, their contribution to the $h\bar{b}b$ vertex can be split into two parts with one part proportional to A_t and the other part proportional to μ. Fig.~2 then reflects that the non-decoupling contribution of the former part is exactly canceled out, while that of the latter part is maintained if m_A is not at the same order as M_{SUSY}. This situation is actually similar to the SUSY-QCD correction to the $h\bar{b}b$ vertex with the only difference coming from the fact that such remnant effect is not enhanced by $\tan\beta$. In fact, if both M_{SUSY} and m_A approach infinity simultaneously, the genuine

SUSY contribution to $t \to ch$ should vanish since now the Higgs sector of the MSSM is identical to that of the SM, while if only $M_{\rm SUSY}$ approaches infinity, the Higgs sector is described by a Two-Higgs-Doublet model, and SUSY may leave its imprint in the Higgs sector [?]. Our results in Fig.~2 actually reflect such a possibility.

We finally discuss how large SUSY can predict the rate of $t \to ch$ in the heavy SUSY case. For this purpose, we require all squarks to be heavier than 3 TeV and perform two independent scans over relevant SUSY parameter space by considering the constraints listed in Section II. These scans are:

- **Scan-III**: Similar to Scan-I except that we fix $m_A = 800$ GeV and consider the following parameter space:

$$
\begin{aligned}
&4\,\text{TeV} \le m_{Q2}, m_{Q3}, m_{U2}, m_{U3}, m_{D2}, m_{D3} \le 7\,\text{TeV}, \\
&|A_t| \le 6\sqrt{m_{Q3} m_{U3}}, \quad 4\,\text{TeV} \le m_{\tilde{g}} \le 10\,\text{TeV}, \\
&|\mu| \le 20\,\text{TeV}, \quad 1 \le \tan\beta \le 40, \\
&-1 \le \delta_{LL}, \delta_{RR} \le 1, \quad -2.0 \le \delta_{LR}, \delta_{RL} \le 2.0, \\
&-0.5 \le \delta^d_{LR}, \delta^d_{RL} \le 0.5.
\end{aligned}
$$

  In our scan, we do not consider very large squark soft breaking parameters because in such a case, the Higgs mass calculated by FeynHiggs suffers from large theoretical uncertainties.

- **Scan-IV**: Similar to Scan-II except that the scan regions are now given by the parameter space above.

For both scans, we find the rate of $t \to ch$ may reach 10^{-8} in the optimal case, and a smaller value of m_A can lead to a larger branching ratio. In Table II, we provide two benchmark points for future study with Point 5 obtained from Scan-III and Point 6 from Scan-IV. One can learn that, compared with the prediction in heavy SUSY with that in low-energy SUSY, although the optimal values of $\text{Br}(t \to ch)$ are suppressed by at least two orders of magnitude, they are still 10^6 times larger than the corresponding SM prediction.

IV. Conclusion

In this work, we studied the top quark FCNC decay $t \to ch$ in the MSSM under the constraints from the Higgs mass measurement, the LHC searches for sparticles, vacuum stability, precision electroweak observables, and $B \to X_s\gamma$. From a scan over the relevant parameter space, we found:

- Due to the strong constraints from the measured Higgs mass and the results of SUSY searches at the LHC, the branching ratio of $t \to ch$ can only reach $\mathcal{O}(10^{-6})$ in the MSSM, which is about one order of magnitude smaller than the old results.

- In the singlet extension of the MSSM, which can lift the Higgs mass at tree level, $\mathrm{Br}(t \to ch)$ can reach $\mathcal{O}(10^{-5})$ in the allowed parameter space.
- The chiral-conserving mixings δ_{LL} and δ_{RR} can induce SUSY remnant effects on the rate of $t \to ch$. For heavy squarks and gluino above 3 TeV, $\mathrm{Br}(t \to ch)$ can still reach 10^{-8}.

Acknowledgement

We thank Nima Arkani-Hamed, M.J. Herrero, Zhaoxia Heng, and Yosef Nir for helpful discussions. This work was supported by the National Natural Science Foundation of China (NNSFC) under grant Nos. 10821504, 11222548, 11305049 and 11135003, by the Program for New Century Excellent Talents in University, and also by the ARC Center of Excellence for Particle Physics at the Tera-scale.

References

[1] [ATLAS Collaboration], Phys. Lett. B 716 (2012) 1 [arXiv:1207.7214 [hep-ex]].
[2] [CMS Collaboration], Phys. Lett. B 716 (2012) 30 [arXiv:1207.7235 [hep-ex]].
[3] Bruno Mansoulié, talk at the Rencontres de Moriond EW 2013, on behalf of the ATLAS collaboration.
[4] [CMS Collaboration], CMS-PAS-HIG-13-005.
[5] [ATLAS Collaboration], ATLAS-CONF-2013-034.
[6] W. Beenakker, et al., Phys. Rev. Lett. 87, 201805 (2001) [hep-ph/0107081]; L. Reina, S. Dawson and D. Wackeroth, Phys. Rev. D 65, 053017 (2002) [hep-ph/0109066]; W. Beenakker, et al., Nucl. Phys. B 653, 151 (2003) [hep-ph/0211352]; S. Dawson, L. H. Orr, L. Reina and D. Wackeroth, Phys. Rev. D 67, 071503 (2003) [hep-ph/0211438].
[7] The ATLAS collaboration, ATLAS-CONF-2013-080; N. Craig, M. Park and J. Shelton, arXiv:1308.0845 [hep-ph].
[8] T. M. P. Tait and C. -P. Yuan, Phys. Rev. D 63, 014018 (2000) [hep-ph/0007298]; F. Maltoni, K. Paul, T. Stelzer and S. Willenbrock, Phys. Rev. D 64 (2001) 094023 [hep-ph/0106293]; S. Biswas, E. Gabrielli and B. Mele, JHEP 1301 (2013) 088 [arXiv:1211.0499 [hep-ph]]; M. Farina, et al., JHEP 1305 (2013) 022 [arXiv:1211.3736 [hep-ph]]; P. Agrawal, S. Mitra and A. Shivaji, JHEP 1312 (2013) 077 [arXiv:1211.4362 [hep-ph]]; S. Biswas, E. Gabrielli, F. Margaroli and B. Mele, JHEP 07 (2013) 073 [arXiv:1304.1822 [hep-ph]].
[9] N. Craig, et al., Phys. Rev. D 86, 075002 (2012) [arXiv:1207.6794 [hep-ph]]; K. -F. Chen, W. -S. Hou, C. Kao and M. Kohda, Phys. Lett. B 725, 378 (2013) [arXiv:1304.8037 [hep-ph]]; D. Atwood, S. K. Gupta and A. Soni, arXiv:1305.2427 [hep-ph]; C. Zhang and F. Maltoni, Phys. Rev. D 88, 054005

(2013) [arXiv:1305.7386 [hep-ph]]; T. Han and R. Ruiz, arXiv:1312.3324 [hep-ph]; B. Yang, N. Liu and J. Han, arXiv:1308.4852 [hep-ph].
[10] G. Aad et al. [ATLAS Collaboration], arXiv:1403.6293 [hep-ex].
[11] H. E. Haber and G. L. Kane, Phys. Rept. 117, 75 (1985); J. F. Gunion and H. E. Haber, Nucl. Phys. B 272, 1 (1986) [Erratum-ibid. B 402, 567 (1993)].
[12] U. Ellwanger, C. Hugonie and A. M. Teixeira, Phys. Rept. 496, 1 (2010) [arXiv:0910.1785 [hep-ph]].
[13] see, e.g., L. J. Hall, Y. Nomura and S. Shirai, JHEP 1301 (2013) 036 [arXiv:1210.2395 [hep-ph]]; J. Unwin, Phys. Rev. D 86 (2012) 095002 [arXiv:1210.4936 [hep-ph]]; E. Arganda, J. L. Diaz-Cruz and A. Szynkman, Eur. Phys. J. C 73 (2013) 2384 [arXiv:1211.0163 [hep-ph]]; N. Arkani-Hamed, et al., arXiv:1212.6971 [hep-ph]; E. Arganda, J. L. Diaz-Cruz and A. Szynkman, Phys. Lett. B 722 (2013) 100 [arXiv:1301.0708 [hep-ph]]; D. McKeen, M. Pospelov and A. Ritz, arXiv:1303.1172 [hep-ph]; N. Liu, L. Wu, P. W. Wu and J. M. Yang, JHEP 1301, 161 (2013) [arXiv:1208.3413 [hep-ph]]; A. Djouadi and J. Quevillon, arXiv:1304.1787 [hep-ph]; C. Han, et al., arXiv:1307.3790 [hep-ph].
[14] H. E. Haber, et al., Phys. Rev. D 63, 055004 (2001) [hep-ph/0007006].
[15] J. Cao, et al., JHEP 1304, 134 (2013) [arXiv:1301.6437 [hep-ph]].
[16] M. Arana-Catania, E. Arganda and M. J. Herrero, JHEP 1309, 160 (2013) [arXiv:1304.3371 [hep-ph]].
[17] T. Appelquist and J. Carazzone, Phys. Rev. D 11 (1975) 2856.
[18] J. A. Coarasa Perez, R. A. Jimenez and J. Sola, Phys. Lett. B 389, 312 (1996) [hep-ph/9511402]; J. A. Coarasa Perez, D. Garcia, J. Guasch, R. A. Jimenez and J. Sola, Eur. Phys. J. C 2, 373 (1998) [hep-ph/9607485]; K. S. Babu and C. F. Kolda, Phys. Lett. B 451, 77 (1999) [hep-ph/9811308]; J. Guasch, W. Hollik and S. Penaranda, Phys. Lett. B 515, 367 (2001) [hep-ph/0106027]; A. Crivellin, Phys. Rev. D 83 (2011) 056001 [arXiv:1012.4840 [hep-ph]]; A. Crivellin, L. Hofer and J. Rosiek, JHEP 1107 (2011) 017 [arXiv:1103.4272 [hep-ph]]; A. Crivellin and C. Greub, Phys. Rev. D 87 (2013) 015013 [arXiv:1210.7453 [hep-ph]].
[19] G. Eilam, J. L. Hewett and A. Soni, Phys. Rev. D 44, 1473 (1991) [Erratum-ibid. D 59, 039901 (1999)]; B. Mele, S. Petrarca and A. Soddu, Phys. Lett. B 435, 401 (1998) [hep-ph/9805498]; J. L. Diaz-Cruz, et al., "Flavor-changing radiative decay of the t quark," Phys. Rev. D41 (1990).
[20] J. Cao, G. Eilam, K. -i. Hikasa and J. M. Yang, Phys. Rev. D 74, 031701 (2006) [hep-ph/0604163].
[21] J. Cao, et al., Phys. Rev. D 75, 075021 (2007) [hep-ph/0702264].
[22] Y. Wang, F. P. Huang, C. S. Li, B. H. Li, D. Y. Shao and J. Wang, Phys. Rev. D 86 (2012) 094014 [arXiv:1208.2902 [hep-ph]]; J. M. Yang, C. S. Li, Physical Review D, 1994, 49(7) 3412; J. Cao, L. Wang, L. Wu and J. M. Yang, Phys. Rev. D 84, 074001 (2011) [arXiv:1101.4456 [hep-ph]]; J. Cao, Z. Heng, L. Wu and J. M. Yang, Phys. Rev. D 81, 014016 (2010) [arXiv:0912.1447 [hep-ph]].
[23] M. Blanke, et al., JHEP 1306, 022 (2013) [arXiv:1302.7232 [hep-ph]].
[24] See, e.g., S. Dimopoulos, D. Sutter, Nucl. Phys. B 452, 496 (1996); F. Gabbiani, et al., Nucl. Phys. B 477, 321 (1996); M. Misiak, S. Pokorski, J.

Rosiek, hep-ph/9703442.
[25] K. Hikasa and M. Kobayashi, Phys. Rev. D 36, 724 (1987); S. Bertolini, F. Borzumati, A. Masiero and G. Ridolfi, Nucl. Phys. B 353, 591 (1991); K. I. Izawa, Y. Nomura, K. Tobe and T. Yanagida, Phys. Rev. D 56, 2886 (1997) [hep-ph/9705228].
[26] S. Heinemeyer, W. Hollik, F. Merz and S. Penaranda, Eur. Phys. J. C 37, 481 (2004) [hep-ph/0403228].
[27] M. Frank, et al., JHEP 0702, 047 (2007) [hep-ph/0611326]; G. Degrassi, et al., Eur. Phys. J. C 28, 133 (2003) [hep-ph/0212020]; S. Heinemeyer, W. Hollik and G. Weiglein, Eur. Phys. J. C 9, 343 (1999) [hep-ph/9812472]; S. Heinemeyer, W. Hollik and G. Weiglein, Comput. Phys. Commun. 124, 76 (2000) [hep-ph/9812320].
[28] G. Aad et al. [ATLAS Collaboration], JHEP 1310, 130 (2013) [arXiv:1308.1841 [hep-ex]]; G. Aad et al. [ATLAS Collaboration], arXiv:1402.7029 [hep-ex]; G. Aad et al. [ATLAS Collaboration], arXiv:1403.5294 [hep-ex]; G. Aad et al. [ATLAS Collaboration], arXiv:1403.4853 [hep-ex]; S. Chatrchyan et al. [CMS Collaboration], arXiv:1402.4770 [hep-ex]; S. Chatrchyan et al. [CMS Collaboration], arXiv:1311.4937 [hep-ex]; S. Chatrchyan et al. [CMS Collaboration], Eur. Phys. J. C 73, 2677 (2013) [arXiv:1308.1586 [hep-ex]]; CMS Collaboration [CMS Collaboration], CMS-PAS-SUS-13-006; CMS Collaboration [CMS Collaboration], CMS-PAS-SUS-12-022.
[29] See, e.g., J. Cao, et al., JHEP 1211, 039 (2012) [arXiv:1206.3865 [hep-ph]]; C. Han, et al., JHEP 1310, 216 (2013) [arXiv:1308.5307 [hep-ph]].
[30] R. Mahbubani, et al., Phys. Rev. Lett. 110, no. 15, 151804 (2013) [arXiv:1212.3328 [hep-ph]].
[31] J. -h. Park, Phys. Rev. D 83, 055015 (2011) [arXiv:1011.4939 [hep-ph]].
[32] D. Chowdhury, R. M. Godbole, K. A. Mohan and S. K. Vempati, arXiv:1310.1932 [hep-ph].
[33] M. Baak, et al., arXiv:1310.6708 [hep-ph].
[34] Y. Amhis et al. [Heavy Flavor Averaging Group Collaboration], arXiv:1207.1158 [hep-ex].
[35] J. Rosiek, hep-ph/9511250.
[36] G. Passarino and M. J. G. Veltman, Nucl. Phys. B 160, 151 (1979); G. ' t Hooft and M. J. G. Veltman, Nucl. Phys. B 153, 365 (1979).
[37] J. Guasch, In *Barcelona 1997, Quantum effects in the minimal supersymmetric standard model* 256-265 [hep-ph/9710267]; J. Guasch and J. Sola, Nucl. Phys. B 562, 3 (1999) [hep-ph/9906268]; J. Guasch and J. Sola, hep-ph/9909503.
[38] D. Das, U. Ellwanger and A. M. Teixeira, JHEP 1204, 067 (2012) [arXiv:1202.5244 [hep-ph]].
[39] K. Agashe et al. [Top Quark Working Group Collaboration], arXiv:1311.2028 [hep-ph].
[40] M. Bicer et al. [TLEP Design Study Working Group Collaboration], JHEP 1401, 164 (2014) [arXiv:1308.6176 [hep-ex]].
[41] M. S. Carena, D. Garcia, U. Nierste and C. E. M. Wagner, Nucl. Phys. B 577, 88 (2000) [hep-ph/9912516]; J. Guasch, P. Hafliger and M. Spira, Phys.

Rev. D 68, 115001 (2003) [hep-ph/0305101].

Note: Figure translations are in progress. See original paper for figures.