Higgs Boson Flavor-Changing Neutral Decays into Top Quark in a General Two-Higgs-Doublet Model

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Abstract

Higgs boson decays mediated by flavor changing neutral currents (FCNC) are very much suppressed in the Standard Model, at the level of $10^{-15}$ for Higgs boson masses of a few hundred GeV. Therefore, any experimental vestige of them would immediately call for new physics. In this paper we consider the FCNC decays of Higgs bosons into a top quark in a general two-Higgs-doublet model (2HDM). The isolated top quark signature, unbalanced by any other heavy particle, should help to identify the potential FCNC events much more than any other final state. We compute the maximum branching ratios and the number of FCNC Higgs boson decay events at the LHC collider at CERN. The most favorable mode for production and subsequent FCNC decay is the lightest CP-even state in the Type II 2HDM, followed by the other CP-even state, if it is not very heavy, whereas the CP-odd mode can never be sufficiently enhanced. Our calculation shows that the branching ratios of the CP-even states may reach $10^{-5}$, and that several hundred events could be collected in the highest luminosity runs of the LHC. We also point out some strategies to use these FCNC decays as a handle to discriminate between 2HDM and supersymmetric Higgs bosons.
1 Introduction

At the tree-level there are no Flavor Changing Neutral Currents (FCNC) processes in the Standard Model (SM), and at one-loop they are induced by charged-current interactions, which are GIM-suppressed \[1\]. Letting aside the meson-meson oscillations, such as $K^0 - \bar{K}^0$ and $B^0 - \bar{B}^0$, the decay processes mediated by FCNC are also of high interest and are strongly suppressed too. For instance, we have the radiative B-meson decays, with a typical branching ratio $BR(b \rightarrow s \gamma) \sim 10^{-4}$. But we also have the FCNC decays with the participation of the top quark as a physical field, which are by far the most suppressed decay modes \[2,3\]. Indeed, the top quark decays into gauge bosons ($t \rightarrow c V; V \equiv \gamma, Z, g$) are well known to be extremely rare events in the SM. The branching ratios are, according to Ref. \[3\]: $\sim 5 \times 10^{-13}$ for the photon, slightly above $1 \times 10^{-13}$ for the $Z$-boson, and $\sim 4 \times 10^{-11}$ for the gluon channel, or even smaller according to other estimates \[4\]. Similarly, the top quark decay into the SM Higgs boson, $H^{SM}$, is a very unusual decay, typically $BR(t \rightarrow c H^{SM}) \sim 10^{-14}$ \[5\]. However, when considering physics beyond the SM, new horizons of possibilities open up which may radically change the pessimistic prospects for FCNC decays involving a Higgs boson and the top quark. For example, in Ref. \[6\] it was shown that the vector boson modes can be highly enhanced within the context of the Minimal Supersymmetric Standard Model (MSSM) \[7\]. This fact was also dealt with in great detail in Ref. \[8\] where in addition a dedicated study was presented of the FCNC top quark decays into the various Higgs bosons of the MSSM (see also \[9\]), showing that these can be the most favored FCNC top quark decays – above the expectations on the gluon mode $t \rightarrow c g$. A similar study was performed for the FCNC top quark decays into Higgs bosons in a general two-Higgs-doublet model (2HDM) \[10\].

The low observed rates of FCNC processes among the SM fields strongly suggest that their FCNC couplings must be suppressed at the tree-level in any extension of the SM, in particular in the 2HDM models. As is well known, there are two canonical strategies to get rid of these tree-level couplings in that kind of models \[11\]. In Type I 2HDM (also denoted 2HDM I) one Higgs doublet, $\Phi_1$, does not couple to fermions at all and the other Higgs doublet, $\Phi_2$, couples to fermions in the same manner as in the SM. In contrast, in Type II 2HDM (also denoted 2HDM II) one Higgs doublet, $\Phi_1$, couples to down quarks (but not to up quarks) while $\Phi_2$ does the other way around. Such a coupling pattern is automatically realized in the framework of supersymmetry (SUSY), in particular in the MSSM, but it can also be arranged in non-supersymmetric extensions if we impose a discrete symmetry, e.g. $\Phi_1 \rightarrow -\Phi_1$ and $\Phi_2 \rightarrow +\Phi_2$ (or vice versa) plus a suitable transformation for the right-handed quark fields. In an analysis of FCNC top quark decays in 2HDM extensions of the SM \[10,12\] it was proven that while the maximum rates for $t \rightarrow c g$ were one order of magnitude more favorable in the MSSM than in the 2HDM, the corresponding rates for $t \rightarrow c H^0$ were comparable both for the MSSM and the general 2HDM, namely up to the $10^{-4}$ level and should therefore be visible both at the LHC and the LC \[10,12,13\].

Similarly, one may wonder whether the FCNC decays of the Higgs bosons themselves can be of some relevance. Obviously the situation with the SM Higgs is essentially hopeless, so again we have to move to physics beyond the SM. Some work on these decays,
performed in various contexts including the MSSM, shows that these effects can be important \[14,15,16,17,18\]. This could be expected, at least for heavy quarks in the MSSM, from the general SUSY study (including both strong and electroweak supersymmetric effects) of the FCNC vertices $htc$ \((h = h^0, H^0, A^0)\) made in Ref. \[8\]. However, other frameworks could perhaps be equally advantageous. Here we are particularly interested in the FCNC Higgs decay modes into top quark within a general 2HDM, which have not been studied anywhere in the literature to our knowledge. It means we restrict to Higgs bosons heavier than \(m_t\). From the above considerations, and most particularly on the basis of the detailed results obtained in \[10,12\] one may expect that some of the decays of the Higgs bosons

\[
h \rightarrow t \bar{c}, \quad h \rightarrow \bar{t}c \quad (h = h^0, H^0, A^0)
\]

in a general 2HDM can be substantially enhanced and perhaps can be pushed up to the visible level, particularly for \(h^0\) which is the lightest CP-even spinless state in these models \[11\]. This possibility can be of great relevance on several grounds. On the one hand the severe degree of suppression of the FCNC Higgs decay in the SM obviously implies that any experimental sign of Higgs-like FCNC decay \((1)\) would be instant evidence of physics beyond the SM. On the other hand, the presence of an isolated top quark in the final state, unbalanced by any other heavy particle, is an unmistakable carrier of the FCNC signature. Finally, the study of the maximum FCNC rates for the top quark modes \((1)\) within the 2HDM, which is the simplest non-trivial extension of the SM, should serve as a fiducial result from which more complicated extensions of the SM can be referred to. Therefore, we believe there are founded reasons to perform a thorough study of the FCNC Higgs decays in minimal extensions of the Higgs sector of the SM and see whether they can be of any help to discover new physics.

The paper is organized as follows. In Section 2 we summarize the 2HDM interactions most relevant for our study and estimate the expected FCNC rates of the Higgs decays in the SM and the general 2HDM. In Section 3 a detailed numerical analysis of the one-loop calculations of the FCNC decay widths and production rates of FCNC Higgs events is presented. Finally, in Section 4 we discuss the reach of our results and its phenomenological implications, and deliver our conclusions.

## 2 Expected branching ratios in the SM and the 2HDM

Before presenting the detailed numerical results of our calculation, it may be instructive to estimate the typical expected widths and branching ratios \((BR)\) both for the SM decay \(H^{SM} \rightarrow t \bar{c}\) and the non-standard decays \((1)\) in a general 2HDM. This should be especially useful in this kind of rare processes, which in the strict context of the SM are many orders of magnitude out of the accessible range. Therefore, one expects to be able to grossly reproduce the order of magnitude from simple physical considerations based on dimensional analysis, power counting, CKM matrix elements and dynamical features. By the same token it should be possible to guess at the potential enhancement factors in
the 2HDM extension of the SM. In fact, guided by the previous criteria the FCNC decay width of the SM Higgs of mass $m_H$ into top quark is expected to be of order

$$
\Gamma(H^{SM} \rightarrow t \bar{c}) \sim \left(\frac{1}{16\pi^2}\right)^2 |V_{tb}V_{bc}|^2 m_H \ (\lambda_b^{SM})^4 \sim \left(\frac{|V_{bc}|}{16\pi^2}\right)^2 \alpha_W \ G_F^2 \ m_H \ m_b^4, \ (2)
$$

where $G_F$ is Fermi's constant and $\alpha_W = g^2/4\pi$, $g$ being the $SU(2)_L$ weak gauge coupling. We have approximated the loop form factor by just a constant prefactor. Notice the presence of $\lambda_b^{SM} \sim m_b/M_W$, which is the SM Yukawa coupling of the bottom quark in units of $g$. The fourth power of $\lambda_b^{SM}$ in (2) gives the non-trivial suppression factor reminiscent of the GIM mechanism after summing over flavors. Since we are maximizing our estimation, a missing function related to kinematics and polarization sums, $F(m_t/m_H) \sim (1-m_t^2/m_H^2)^2$, has been approximated to one. To obtain the (maximized!) branching ratio it suffices to divide the previous result by $\Gamma(H^{SM} \rightarrow bb) \sim \alpha_W \ (\lambda_b^{SM})^2 \ m_H \sim G_F m_H m_b^2$

to obtain

$$
BR(H^{SM} \rightarrow t \bar{c}) \sim \left(\frac{|V_{bc}|}{16\pi^2}\right)^2 \alpha_W \ G_F m_b^2 \sim 10^{-13}, \ (3)
$$

with $V_{bc} = 0.04$, $m_b = 5 \text{GeV}$. In general this $BR$ will be even smaller, specially for higher Higgs boson masses ($m_H > 2 M_W$) for which the vector boson Higgs decay modes $H^{SM} \rightarrow W^+ W^- (Z Z)$ can be kinematically available and become dominant. In this case it is easy to see that $BR(H^{SM} \rightarrow t \bar{c})$ will be suppressed by an additional factor of $m_b^2/m_H^2$, which amounts at the very least to two additional orders of magnitude suppression, bringing it to a level of less than $10^{-15}$. Already the optimized branching ratio (3) will remain invisible to all foreseeable accelerators in the future! To obtain the corresponding maximized estimation for the 2HDM we remind the following basic interaction Lagrangians within the general 2HDM. The charged Higgs interactions with fermions read

$$
\mathcal{L}^{(j)}_{Htb} = \frac{g V_{tb}}{\sqrt{2} \ M_W} \ H \ H \ [m_t \ cot \beta \ P_R + m_b a_j \ P_L] \ t + h.c. \ (4)
$$

We use third-quark-family notation as a generic one; $P_{L,R} = (1/2) (1 \mp \gamma_5)$ are the chiral projection operators on left- and right-handed fermions; $j = I, II$ runs over Type I and Type II 2HDM’s. A most relevant parameter here is $\tan \beta$, the ratio between the vacuum expectation values of the two Higgs doublet models [11]. In the Lagrangian above, the parameter $a_j$ is such that $a_I = -\cot \beta$ and $a_{II} = +\tan \beta$. For the neutral Higgs interactions, the necessary pieces of the Lagrangian are the following:

$$
\mathcal{L}^{(j)}_{hqq} = \ -\frac{g m_b}{2 \ M_W} \ \bar{b} \ \begin{pmatrix} \cos \alpha \\ \sin \beta \cos \alpha \end{pmatrix} + \begin{pmatrix} \cos \alpha \\ -\sin \alpha \end{pmatrix} H^0 + i \frac{g m_b a_j}{2 \ M_W \ \gamma_5} \ b \ A^0
$$

$$
\quad + \ -\frac{g m_t}{2 \ M_W \ \sin \beta} \ \bar{t} \ \begin{pmatrix} \cos \alpha \cos H^0 \sin \alpha \end{pmatrix} t + i \frac{g m_t}{2 \ M_W \ \tan \beta} \ \gamma_5 \ t \ A^0, \ (5)
$$

where the upper row is for Type I models and the down row is for Type II. Here $\alpha$ is the mixing angle in the CP-even Higgs sector. Apart from the interactions with fermions, there is a set of potentially very relevant trilinear self-interactions of Higgs bosons in a
\[
\begin{array}{|c|c|}
\hline
H^\pm H^\mp H^0 & -\frac{g}{M_W \sin 2\beta} \left[ (m_{H^\pm}^2 - m_{A^0}^2 + \frac{1}{2}m_{H^0}^2) \sin 2\beta \cos(\beta - \alpha) + (m_{A^0}^2 - m_{H^0}^2) \cos 2\beta \sin(\beta - \alpha) \right] \\
H^\pm H^\mp h^0 & -\frac{g}{M_W \sin 2\beta} \left[ (m_{H^\pm}^2 - m_{A^0}^2 + \frac{1}{2}m_{h^0}^2) \sin 2\beta \cos(\beta - \alpha) + (m_{h^0}^2 - m_{A^0}^2) \cos 2\beta \sin(\beta - \alpha) \right] \\
h^0 h^0 H^0 & -\frac{g \cos(\beta - \alpha)}{2 M_W \sin 2\beta} \left[ (2m_{h^0}^2 + m_{H^0}^2) \sin 2\alpha - m_{A^0}^2 (3 \sin 2\alpha - \sin 2\beta) \right] \\
A^0 A^0 H^0 & -\frac{g}{2 M_W \sin 2\beta} \left[ m_{h^0}^2 \sin 2\beta \cos(\beta - \alpha) + 2(m_{h^0}^2 - m_{A^0}^2) \cos 2\beta \sin(\beta - \alpha) \right] \\
A^0 A^0 h^0 & -\frac{g}{2 M_W \sin 2\beta} \left[ m_{h^0}^2 \sin 2\beta \sin(\beta - \alpha) + 2(m_{h^0}^2 - m_{A^0}^2) \cos 2\beta \cos(\beta - \alpha) \right] \\
\hline
\end{array}
\]

Table 1: The needed trilinear self-couplings of the Higgs bosons in the 2HDM within the framework of Ref. [10].

general 2HDM. These are summarized in Table II*. Notice that the trilinear couplings in Table II are common to the 2HDM I and 2HDM II, and they are quite different from the ones obtained in the MSSM, whose pure gauge structure is enforced by the underlying supersymmetry.

Let us now first assume large tan \( \beta \) and restrict to Type II models. From the interaction Lagrangians above it is clear that we may replace \( \lambda_b^{SM} \to \lambda_b^{SM} \tan \beta \) in the previous formulae for the partial width. Moreover, the leading diagrams in the 2HDM contain the trilinear Higgs couplings \( \lambda_{H^+H^-h} \). Therefore, the maximum \( BR \) associated to the FCNC decays \( \{1\} \) in a general 2HDM II should be of order\(^{\dagger}\)

\[
BR^{II}(h \to t\bar{c}) \sim \left( \frac{|V_{bc}|}{16\pi^2} \right)^2 \alpha_W G_F m_b^2 \tan^2 \beta \lambda_{H^+H^-h}^2,
\]

where \( \lambda_{H^+H^-h} \) is defined here in units of \( g \) and dimensionless as compared to Table II. Clearly a big enhancement factor \( \tan^2 \beta \) appears, but this does not suffice. Fortunately, the trilinear couplings \( \lambda_{H^+H^-h} \) for \( h = h^0, H^0 \) (but not for \( h = A^0 \)) carry two additional sources of potential enhancement (Cf. Table II) which are absent in the MSSM case. Take e.g. \( h^0 \), then we see that under appropriate conditions (for example, large \( \tan \alpha \) and large \( \tan \beta \)) the trilinear coupling behaves as \( \lambda_{H^+H^-h^0} \sim (m_{h^0}^2 - m_{A^0}^2) \tan \beta / (M_W m_{H^\pm}) \), and in this case

\[
BR^{II}(h^0 \to t\bar{c}) \sim \left( \frac{|V_{bc}|}{16\pi^2} \right)^2 \alpha_W G_F m_b^2 \tan^4 \beta \left( \frac{m_{A^0}^2 - m_{h^0}^2}{M_W m_{H^\pm}} \right)^2.
\]

So finally \( BR^{II}(h^0 \to t\bar{c}) \), and of course \( BR^{II}(h^0 \to \bar{t}c) \), can be augmented by a huge factor \( \tan^4 \beta \) times the square of the relative splitting among the CP-even Higgs decaying

\(^\dagger\)Here we have normalized the \( BR \) with respect to the \( h \to b\bar{b} \) channel only, because the gauge boson modes will be suppressed in the relevant FCNC region, Cf. Section 3.

*Table II extends the one in Ref. [10], and therefore we also assume here \( \lambda_5 = \lambda_6 \) in the 2HDM Higgs potential [11]. We have omitted the couplings with the Goldstone bosons. Although the latter have been included in the calculation, their potential enhancement is far less significant than in the case of the Higgs trilinear couplings.
boson mass and the CP-odd Higgs mass. Since the neutral Higgs bosons do not participate in the loop form factors (Cf. Fig. 1 Ref. [10]), it is clear that various scenarios can be envisaged where these mass splittings can be relevant. In the next section this behaviour will be borne out by explicit calculations showing that $h^0 \rightarrow t \bar{c}$ can be raised to the visible level in the case of the Type II model. As for the Type I model the Higgs trilinear coupling enhancement is the same, but in the charged Higgs Yukawa coupling all quarks go with a factor $\cot \beta$; hence when considering the leading terms in the loops that contribute one sees that in the corresponding expression the term $m_b^2 \tan^2 \beta$ is traded for $m_t m_c \cot^2 \beta$, which is negligible at high $\tan \beta$. Both sources of enhancement are needed, and this feature is only tenable in the 2HDM II. Of course one could resort to the range $\tan \beta \ll 1$ for the Type I models, but this is not theoretically appealing. For example, for $\tan \beta \lesssim 0.1$ the top quark Yukawa coupling $g_t = g m_t / (\sqrt{2} M_W \sin \beta)$, which is present in the interaction Lagrangians above, is pushed into the non-perturbative region $g_t^2 / 16\pi^2 \gtrsim 1$ and then our calculation would not be justified. And what is worse: for the 2HDM I we would actually need $\tan \beta \lesssim \mathcal{O}(10^{-2})$ to get significant FCNC rates! In short, we consider that $BR^I(h \rightarrow t \bar{c} + \bar{t}c)$ is essentially small (for all $h$), and that these decays remain always invisible to speak of. Hereafter we abandon the study of the decays $\dagger$ for the 2HDM I and restrict ourselves to the general 2HDM II.

### 3 Numerical analysis

Let us now substantiate the previous claims and provide the precise numerical results of the full one-loop calculation of $BR^{II}(h \rightarrow t \bar{c} + \bar{t}c)$ as well as of the LHC production rates of these FCNC events. We shall closely follow the notation and methods of Refs. [8, 10, 12]. We refer the reader to these references for more details. See also [11] for basic definitions in the general 2HDM framework and [19, 20] for calculational techniques and further illustration of the $\tan \beta$ enhancement in other relevant Higgs processes both in the MSSM and the 2HDM. In what follows we shall limit ourselves to present the final results of our numerical analysis together with a detailed discussion, interpretation and phenomenological application. We have performed the calculations with the help of the numeric and algebraic programs FeynArts, FormCalc and LoopTools [21]. The calculation must obviously be finite without renormalization, and indeed the cancellation of UV divergences in the total amplitudes was verified explicitly.

The input set for our numerical analysis is given by the data row

$$\left( m_{h^0}, m_{H^0}, m_{A^0}, m_{H^\pm}, \tan \alpha, \tan \beta \right)$$

made out of six independent parameters in the general 2HDM. Remaining inputs as in [22]. In practice there are some phenomenological restrictions on the data [8] which were already described in [10] and references therein, particularly [23]. It suffices to say here that the Higgs boson masses in [8] are much less restricted than in the MSSM. Actually, the direct experimental searches put a soft bound of $\sim 70 \text{ GeV}$ for the lightest

\[ \text{Here and throughout we use the notation } BR^{II}(h^0 \rightarrow t \bar{c} + \bar{t}c) = BR^{II}(h^0 \rightarrow t \bar{c}) + BR^{II}(h^0 \rightarrow \bar{t}c). \]
neutron Higgs boson and around 80 GeV for the charged one \(H^\pm\). These are much more relaxed than the bound of 114.1 GeV placed on the SM Higgs boson [24]. In the MSSM, direct LEP 200 searches set a limit of \(\gtrsim 90\) GeV for \(m_{h^0}\) [24] and a similar limit for \(m_{A^0}\) which translates (approximately) to \(m_{H^\pm} \simeq 120\) GeV due to the supersymmetric Higgs mass relations. Moreover, in contrast to the general 2HDM, in the MSSM case there is a theoretical upper bound \(m_{h^0} \lesssim 122, 135\) GeV (for minimal and maximal top-squark mixing, respectively) [25]. Concerning tan \(\beta\), in a general 2HDM it is restricted only, by perturbativity arguments, to lie within the approximate range

\[
0.1 \lesssim \tan \beta \lesssim 60. \quad (9)
\]

Strictly speaking the upper bound could be larger, but as in [10] we shall stick to the above moderate range. In practice, since Type I models are not considered, the effective range for our calculation will be the high tan \(\beta\) end of (9). For comparison, in the MSSM case there are additional phenomenological restrictions that bring the lower bound on tan \(\beta\) to roughly 2.4 for the so-called maximal \(m_h^0\) scenario, and 10.5 for the so-called no mixing scenario [24], the upper bound being the same. Furthermore, in contrast to the MSSM, the parameter tan \(\alpha\) is free in the 2HDM.
Apart from the direct limit \( m_{H^\pm} \gtrsim 80 \text{ GeV} \) from LEP 200, there are more stringent restrictions on the charged Higgs boson mass. Although they are indirect and in this sense not so indisputable, it is wise to assess their impact in our analysis. These restrictions emerge from the virtual contributions to \( b \to s \gamma \) that have been computed at the NLO in QCD [26]. From these calculations and from the experimental limits on the radiative B-meson decays, these authors can place a lower bound on the 2HDM II charged Higgs mass of \( m_{H^\pm} \gtrsim 350 \text{ GeV} \) [26]. We will, in general, apply this lower bound, but since the direct and indirect searches complement each other, we may disregard it in some parts of our numerical analysis. The parts where we deviate will be clear in the text. We recall that at the moment the MSSM may escape this indirect bound because the positive charged Higgs virtual contributions to \( b \to s \gamma \) can be compensated for by negative stop and chargino loops, if they are not too heavy. Therefore, in the MSSM the charged Higgs can stay relatively light, \( m_{H^\pm} \gtrsim 120 \text{ GeV} \), just to comply with the aforementioned LEP 200 bounds on \( m_{A^0} \) [24].

Finally, we will restrict the input data row \([8]\) with the important constraint from \( \delta \rho \), extensively used in \([8, 10]\) for the FCNC top quark decays in the MSSM and the 2HDM. Namely, the one-loop corrections to the \( \rho \)-parameter (i.e. the ratio between the charged and neutral current Fermi constants) are bound within \( |\delta \rho^{\text{2HDM}}| \lesssim 0.1\% \). To be precise, the latter is the extra effect that \( \delta \rho \) can accommodate at one standard deviation \((1 \sigma)\) from the 2HDM fields beyond the SM contribution \([8]\). This is a stringent restriction that affects the possible mass splittings among the Higgs fields of the 2HDM, and its implementation in our codes does severely prevent the possibility from playing with the Higgs boson masses to artificially enhance the FCNC contributions.

With these restrictions in mind we have computed the number of FCNC Higgs decay events into top quark at the LHC:

\[
pp \to h + X \to t \bar{c} (\bar{t} c) + X \quad (h = h^0, H^0, A^0).
\]

Figure 2: Contour lines in the \((m_{H^\pm}, m_{h^0})\)-plane for the branching ratios (2HDM II case) (a) \( BR^{II}(h^0 \to t \bar{c} + \bar{t} c) \) and (b) \( BR^{II}(H^0 \to t \bar{c} + \bar{t} c) \) assuming \( \delta \rho^{\text{2HDM}} \) at 1 \( \sigma \).
Figure 3: (a) Contour lines in the $(m_{H^\pm}, m_{h^0})$-plane for the maximum number of light CP-even Higgs FCNC events $h^0 \to t \bar{c} + \bar{t}c$ produced at the LHC for 100 fb$^{-1}$ of integrated luminosity; (b) Contour lines showing the value of $m_{A^0}$ that maximizes the number of events; (c) As in (a) but within $\delta \rho^{2\text{HDM}}$ at 3$\sigma$; (d) As in (b) but with $\delta \rho^{2\text{HDM}}$ at 3$\sigma$.

The necessary cross-sections to compute the production of neutral Higgs bosons at this collider, including all known QCD corrections, have been computed by adapting the codes HIGLU 1.0 and HQQ 1.0\cite{27} – originally written for the MSSM case\cite{28} – to the general case of the 2HDM$^8$. Folding the cross-sections with the one-loop branching ratios of the processes $^{11}$ we have obtained the number of FCNC Higgs decay events at the LHC. Let us first consider the branching ratios themselves. In Fig. $^{11}$a,b we show $BR^{II}(h^0 \to t \bar{c} + \bar{t}c)$ for the lightest CP-even state. In particular, Fig. $^{11}$a shows $BR^{II}(h^0 \to t \bar{c} + \bar{t}c)$ versus the charged Higgs mass $m_{H^\pm}$. In this figure we fix the values of the parameters in $^8$

$^8$We have used the default parton distribution functions and renormalization/factorization scales used in these programs, namely GRV94 with $\mu_R = \mu_F = m_h$ for HIGLU, and CTEQ4L with $\mu_R = \mu_F = \sqrt{\hat{s}} \equiv \sqrt{(p_h + p_Q + p_{\bar{Q}})^2}$ for HQQ.
Figure 4: Contour lines $\alpha/\pi = \text{const.}$ ($\alpha$ is the mixing angle in the CP-even sector) corresponding to Fig. 3 for $\delta \rho^{2\text{HDM}}$ at (a) $1\sigma$ and (b) $3\sigma$.

which are not varying as follows:

\[
(m_{h^0} = 350 \text{ GeV}, m_{H^0} = 600 \text{ GeV}, m_{A^0} = 550 \text{ GeV}, m_{H^\pm} = 375 \text{ GeV}, \\
tan \alpha = 30, \tan \beta = 60)
\]  

(11)

After crossing a local maximum (associated to a pseudo-threshold of the one-loop vertex function involving the $h^0 H^+ H^- \text{ coupling}$) the subsequently falling behavior of the $BR$ with $m_{H^\pm}$ clearly shows that the previously discussed bounds on $m_{H^\pm}$ are quite relevant. The branching ratio, however, stays within $10^{-6} - 10^{-5}$ for a wide range of heavy charged Higgs masses extending up to $m_{H^\pm} \leq 600 \text{ GeV}$ in Fig. 1a. Hence, for $m_{H^\pm}$ heavy enough to satisfy the indirect bounds from radiative $B$-meson decays [26], the maximum $BR$ is still sizeable. In Fig. 1a the production rate of $h^0$ bosons at the LHC is shown as a function of $m_{h^0}$, for fixed parameters [33]. The production cross-sections for the subprocesses

\[
gg \rightarrow h^0 + X, \quad gg, qq \rightarrow h^0 + Q\bar{Q},
\]  

(12)

contributing to (11) in the case of the light CP-even Higgs $h^0$ are explicitly separated in Fig. 1a. The gluon-gluon fusion process proceeds at one-loop and the $h^0 Q\bar{Q}$ associated production proceeds at tree-level [29]. Similar subprocesses and results apply for $H^0$ and $A^0$ production. At large $\tan \beta$ and the larger the Higgs boson masses the particular associated production mechanism with the bottom quark, $Q = b$, i.e. $h^0 b\bar{b}$, becomes dominant by far. All other mechanisms for Higgs boson production in Type II models [28,29,30], like vector-boson fusion (which contributes also to $h^0 Q\bar{Q}$ when $Q$ are light quarks), vector-boson bremsstrahlung ($q\bar{q} \rightarrow h V$) and associated $t\bar{t}$ production, are subdominant at large $\tan \beta$ and can be neglected for our purposes. Admittedly, some of these mechanisms can be relevant for Higgs boson production in the case of the Type I 2HDM at low $\tan \beta$, but we have already warned that the corresponding FCNC branching ratios are never sufficiently high.
The control over $\delta \rho^{2HDM}$ is displayed in Fig. 1d. Recall that $\delta \rho$ is not sensitive to the mass splitting between $m_{h^0}$ and $m_{H^0}$, because of $CP$-conservation in the gauge boson sector, but it does feel all the other mass splittings among Higgs bosons, charged and neutral. A more systematic search of $BR$ values in the parameter space is presented in Figs. 2a,b corresponding to $BR^{II}(h^0 \rightarrow t \bar{c} + \bar{t}c)$ and $BR^{II}(H^0 \rightarrow t \bar{c} + \bar{t}c)$ respectively. Here we have scanned independently on the parameters $\delta \rho^{2HDM}$ bound at $1 \sigma$. The contour lines in these figures represent the locus of points in the $(m_{H^\pm}, m_{h^0})$-plane giving maximized values of the $BR$ in the 2HDM II. Let us remark that the highest value of $\tan \beta$ is always preferred, and therefore all these contour lines correspond to $\tan \beta = 60$.

In practice, to better assess the possibility of detection at the LHC, one has to study the
production rates of the FCNC events. These are determined by combining the production cross-sections of neutral 2HDM II Higgs bosons at the LHC and the FCNC branching ratios. If we just adopt the mild LEP bound $m_{H^\pm} \gtrsim 80 \text{GeV}$ and let $m_{H^\pm}$ approach the maximum in Fig. 1a then the $BR$ can be as large as $10^{-3}$ and the number of FCNC events can be huge, at the level of ten thousand per 100 fb$^{-1}$ of integrated luminosity. But of course the region near the maximum is too special. Moreover, if we switch on the above mentioned indirect bound from $b \to s \gamma$ [26], then the typical $BR$ is much smaller (of order $10^{-5}$) and the number of events is reduced dramatically, at a level of hundred or less for the same integrated luminosity. On the other hand it may well happen that there are regions of parameter space where $BR \sim 10^{-5}$ (see Fig. 2) but the production cross-section is too small because the decaying Higgs boson is too heavy. Therefore, it is the product of the two quantities that matters.

The systematic search of the regions of parameter space with the maximum number of FCNC events for the light CP-even Higgs is presented in the form of contour lines in the multiple Fig. 3. For instance, each isoline in Figs. 3a,c corresponds to a fixed number of produced FCNC events at the LHC while keeping the value of $\delta \rho^{\text{2HDM}}$ within 1 $\sigma$ or 3 $\sigma$ respectively of its central experimental value. When scanning over the parameter space we have found again that $\tan \beta$ is preferred at the highest allowed value ($\tan \beta = 60$) – for Type II models. We have also determined (see Figs. 3b,d) the corresponding contour lines for $m_{A^0}$ associated to these events. The $m_{A^0}$-lines are important because the FCNC processes under consideration are sensitive to the mass splittings between $m_{A^0}$ and the corresponding decaying Higgs boson, see e.g. eq. (7) and Table 1. The combined figures 3a-d are very useful because they give a panoramic view of the origin of our results in the parameter space. To complete the map of the numerical analysis we provide Fig. 4 in which we have projected the contour lines of the CP-even mixing angle $\alpha$ associated to the previous plots. For a given contour line $\alpha/\pi = \text{const.}$, the set of inner points have a value of $\alpha/\pi$ smaller than the one defined by the line itself. In particular, the large
domains in Figs. 4a,b without contour lines correspond to $\alpha/\pi > 0.4$ and so to relatively large (and positive) $\tan \alpha$. There are a few and small neighborhoods where the FCNC rates for $h^0$ can be sizeable also for small $\tan \alpha$.

Knowing that high $\tan \alpha$ is generally preferred by $h^0 \to t \bar{c} + \bar{t} c$, and noting from Fig. 3 that large mass splittings between $m_{h^0}$ and $m_{A^0}$ are allowed, we find that the trilinear coupling $\lambda_{H+H-H^0}^\beta$ can take the form $\lambda_{H+H-H^0}^\beta \sim (m_{h^0}^2 - m_{A^0}^2) \tan \beta / (M_W m_{H^+})$. Hence it provides a substantial additional enhancement beyond the $\tan \beta$ factor. One can check from the approximate formula (7) that the maximum FCNC branching ratios $\text{BR}^{II}(h^0 \to t \bar{c} + \bar{t} c)$ can eventually reach the $10^{-5}$ level even in regions where the charged Higgs boson mass preserve the stringent indirect bounds from radiative $B$-meson decays (20). These expectations are well in agreement with the exact numerical analysis presented in Fig. 2, thus showing that eq. (7) provides a reasonable estimate, and therefore a plausible explanation for the origin of the maximum contributions. As a matter of fact, we have checked that the single (finite) Feynman diagram giving rise to the estimation (7) – the one-loop vertex Feynman diagram with a couple of charged Higgs bosons and a bottom quark in the loop – reproduces the full result with an accuracy better than 10% for $\tan \beta \gtrsim 10 - 20$. At lower $\tan \beta$ values large deviations are possible but, as warned before, eq. (7) is expected to be valid only at large $\tan \beta$. Furthermore, for low values of $\tan \beta \lesssim 20$ the FCNC $\text{BR}s$ are too small to be of any phenomenological interest. The exact numerical analysis is of course based on the full expression for the branching ratio

$$\text{BR}^{II}(h^0 \to t \bar{c} + \bar{t} c) = \frac{\Gamma(h^0 \to t \bar{c} + \bar{t} c)}{\Gamma(h^0 \to b \bar{b}) + \Gamma(h^0 \to t \bar{t}) + \Gamma(h^0 \to V V) + \Gamma(h^0 \to H H)},$$

(13)

where all decay widths in the denominator of this formula have been computed at the tree-level in the 2HDM II, since this provides a consistent description of eq. (13) at leading order. Here we have defined

$$\Gamma(h^0 \to V V) \equiv \Gamma(h^0 \to W^+W^-) + \Gamma(H^0 \to ZZ),$$

(14)

$$\Gamma(h^0 \to H H) \equiv \Gamma(h^0 \to A^0 A^0) + \Gamma(h^0 \to H^+ H^-).$$

(15)

We disregard the loop induced decay channels, since they have branching ratios below the percent level all over the parameter space. The $\tau$-lepton decay channel is also neglected, since it is suppressed by a factor of $\mathcal{O}(10^{-2})$ with respect the $b\bar{b}$-channel in the whole 2HDM parameter space. In general the effect of the gauge boson channels $h^0 \to W^+W^-, ZZ$ in the $\text{BR}^{II}$ is not so important as in the SM, actually for $\beta = \alpha$ they vanish in the $h^0$ case because they are proportional to $\sin^2(\beta - \alpha)$. This is approximately the case for large $\tan \alpha$ and large $\tan \beta$, the dominant FCNC region for $h^0$ decay (Cf. Fig. 4a and 4b). In this region, the mode $h^0 \to t \bar{t}$ is, when kinematically allowed, suppressed: $\text{BR}(h^0 \to t \bar{t}) \propto \cos^2 \alpha / \sin^2 \beta \to 0$ (Cf. Eq. (20)). On the other hand there are domains in our plots where the decays $h^0 \to H^+ H^-$ and $h^0 \to A^0 A^0$ are kinematically possible and non-(dynamically) suppressed. Indeed, this can be checked from the explicit structure of the trilinear couplings $h^0 H^+ H^-$ and $h^0 A^0 A^0$ in Table 1, in the dominant region for
the decays $h^0 \to t\bar{c} + \bar{t}c$ both of these couplings are $\tan \beta$-enhanced. Nevertheless the decay $h^0 \to A^0 A^0$ is only possible for $m_{A^0} < m_{h^0}/2$, and since the optimal FCNC regions demand the largest possible values of $m_{A^0}$, this decay is kinematically blocked there. On the other hand the mode $h^0 \to H^+ H^-$ is of course allowed if we just take the aforementioned direct limits on the 2HDM Higgs boson masses. But it is never available if we apply the indirect bound from $b \to s\gamma$ on the charged Higgs mass mentioned above, unless $m_{h^0} > 2m_{H^\pm} > 700 GeV$, in which case $h^0$ is so heavy that its production cross-section is too small for FCNC studies to be further pursued.

The corresponding results for the heavy CP-even Higgs boson are displayed in Figs. 5 and 6. The exact formula for the $BR$ in this case reads

$$BR^{II}(H^0 \to t\bar{c} + \bar{t}c) = \frac{\Gamma(H^0 \to t\bar{c} + \bar{t}c)}{\Gamma(H^0 \to b\bar{b}) + \Gamma(H^0 \to t\bar{t}) + \Gamma(H^0 \to V V) + \Gamma(H^0 \to H H)},$$

where we have defined

$$\Gamma(H^0 \to V V) \equiv \Gamma(H^0 \to W^+W^-) + \Gamma(H^0 \to Z Z),$$

$$\Gamma(H^0 \to H H) \equiv \Gamma(H^0 \to h^0 h^0) + \Gamma(H^0 \to A^0 A^0) + \Gamma(H^0 \to H^+ H^-).$$

From the contour lines in Fig. 4a,c it is patent that the number of FCNC top quark events stemming from $H^0$ decays is comparable to the case of the lightest Higgs boson. However, Fig. 5a,b clearly reveals that these events are localized in regions of the parameter space generally different from the $h^0$ case, namely they prefer $\tan \alpha \simeq 0$. Even so, there are some “islands” of events at large $\tan \alpha$. This situation is complementary to the one observed for $h^0$ in Fig. 4. However, in both cases these isolated regions are mainly concentrated in the segment $m_{H^\pm} < 350 GeV$. Therefore, if the bound on $m_{H^\pm}$ from $b \to s\gamma$ is strictly preserved, it is difficult to find regions of parameter space where the two CP-even states of a general 2HDM II may both undergo a FCNC decay of the type $t\bar{t}$.

In the dominant regions of the FCNC mode $H^0 \to t\bar{c}$ (where $\tan \alpha$ is small and $\tan \beta$ is large), the decay of $H^0$ into the $t\bar{t}$ final state is suppressed: $BR(H \to t\bar{t}) \propto \sin^2\alpha/\sin^2\beta \rightarrow 0$. In the same regions the gauge boson channels in (16) are suppressed too because $\Gamma(H^0 \to W^+W^-, ZZ) \propto \cos^2(\beta - \alpha)$. In principle the heavy CP-even Higgs boson $H^0$ also could (as $h^0$) decay into $A^0 A^0$ and $H^+ H^-$. But there is a novelty here with respect to the $h^0$ decays, in that there could be regions where $H^0$ could decay into the final state $h^0 h^0$. This contingency has been included explicitly in eq. (18). However, in practice, neither one of these three last channels is relevant in the optimal FCNC domains of parameter space. First, the decay $H^0 \to h^0 h^0$, although it is kinematically possible, is dynamically suppressed in the main FCNC region for $H^0$. This can be seen from Table 1, where the trilinear coupling $H^0 h^0 h^0$ becomes vanishingly small at large $\tan \beta$ and small $\tan \alpha$. Second, the coupling $H^0 A^0 A^0$ in Table 1 is non-suppressed in the present region, but again the mode $H^0 \to A^0 A^0$ is kinematically forbidden in the optimal FCNC domains because the latter favor large values of the CP-odd mass (see Fig.5b,d). Third, although in these domains the decay $H^0 \to H^+ H^-$ is also non-dynamically suppressed
(see the corresponding trilinear coupling in Table 11), it becomes kinematically shifted to the high mass range $m_{H^0} > 700 \text{ GeV}$ if we switch on the indirect bound from $b \to s \gamma$. Obviously, in this latter case the $H^0$ production cross section becomes too small and the FCNC study has no interest. All in all the contributions from (14), (15), (17) and (18) are irrelevant for $m_{h^0}, m_{H^0} < 700 \text{ GeV}$ as their numerical impact on $B^{\ell\Pi}(h^0, H^0 \to t\bar{c} + \ell c)$ is negligible. Our formulae (13) and (16) do contain all the decay channels and we have verified explicitly these features.

As remarked before, in general the most favorable regions of parameter space for the FCNC decays of $h^0$ and $H^0$ do not overlap much. The trilinear Higgs boson self-couplings in Table 1 (also the fermionic ones) are interchanged when performing the simultaneous substitutions $\alpha \to \pi/2 - \alpha$ and $m_{h^0} \to m_{H^0}$ [10]. Furthermore, the LHC production rates of the neutral Higgs bosons fall quite fast with the masses of these particles, as seen e.g. in Fig. 1b for the $h^0$ state. As a consequence that exchange symmetry on the branching ratios does not go over to the final event rates, so in practice the number of FCNC events from $H^0$ decays are smaller (for the same values of the other parameters) as compared to those for $h^0$; thus $H^0$ requires e.g. lighter charged Higgs masses to achieve the same number of FCNC events as $h^0$. As for the CP-odd state $A^0$, we have seen that it plays an important indirect dynamical role on the other decays through the trilinear couplings in Table 1 but its own FCNC decay rates never get a sufficient degree of enhancement due to the absence of the relevant trilinear couplings, so we may discard it from our analysis.

We notice that this picture is consistent with the decoupling limit in the 2HDM: for $\alpha \to \beta$, the heaviest CP-even Higgs boson ($H^0$) behaves as the SM Higgs boson, whereas $h^0$ decouples from the electroweak gauge bosons and may develop enhanced couplings to up and down-like quarks, depending on whether $\tan \beta$ is small or large respectively; in the opposite limit ($\alpha \to \beta - \pi/2$), it is $h^0$ that behaves as $H^{SM}$, while $H^0$ decouples from gauge bosons and may develop the same enhanced couplings to quarks as $h^0$ did in the previous case. Indeed these are the situations that we find concerning the FCNC decay rates. We recall that the numerical results presented in our figures correspond to an integrated luminosity of $100 \text{ fb}^{-1}$. However, the combined ATLAS and CMS detectors might eventually accumulate a few hundred inverse femtobarn [31, 32]. Therefore, hopefully, a few hundred FCNC events (1) could eventually be collected in the most optimistic scenario. Actually, the extreme rareness of these events in the SM suggests that if only a few of them could be clearly disentangled, it should suffice to claim physics beyond the SM.

4 Discussion and conclusions

Detection strategies at the CERN-LHC collider for the search of the SM Higgs boson, and also for the three spinless fields of the MSSM Higgs sector, have been described in great detail in many places of the literature [11, 31, 32, 33, 34], but not so well for the corresponding charged and neutral Higgs bosons of the general 2HDM. The result is that the discovery of the SM Higgs boson is guaranteed at the LHC in the whole
presumed range $100 \text{GeV} \lesssim m_H \lesssim 1 \text{TeV}$. However, the discovery channels are different in each kinematical region and sometimes the most obvious ones are rendered useless. For example, due to the huge irreducible QCD background from $b\bar{b}$ dijets, the decay mode $H^{SM} \rightarrow b\bar{b}$ is difficult and one has to complement the search with many other channels, particularly $H^{SM} \rightarrow \gamma\gamma \, \, \, [31,32]$. We have shown in this work that there are scenarios in the 2HDM parameter space where alternative decays, like the FCNC modes $h^0 \rightarrow t\bar{c} + \bar{t}c$ and $H^0 \rightarrow t\bar{c} + \bar{t}c$, can also be useful. For instance, in the $h^0$ case, this situation occurs when $\tan\beta$ and $\tan\alpha$ are both large and the CP-odd state is much heavier than the CP-even ones. The potential enhancement is then spectacular and it may reach up to ten billion times the SM value $BR(H^{SM} \rightarrow t\bar{c}) \sim 10^{-15}$, thereby bringing the maximum value of the FCNC branching ratio $BR(h^0 \rightarrow t\bar{c})$ to the level of $\sim 10^{-5}$. As a matter of fact, the enhancement would be much larger were it not because we eventually apply the severe (indirect) lower bound on the charged Higgs mass from $b \rightarrow s\gamma \, \, \, [26]$. Although these decays have maximal ratios below $BR(h \rightarrow \gamma\gamma) \sim 10^{-3}$, they should be essentially free of QCD background.

While in the MSSM almost the full $(m_{A^0}, \tan\beta)$-parameter space is covered, with better efficiency at high $\tan\beta$ though, we should insist that within the general 2HDM the tagging strategies are not so well studied and one would like to have further information to disentangle the MSSM scenarios from the 2HDM ones. Here again the study of the FCNC Higgs decays can play a role. Of course the statistics for the FCNC Higgs decays is poor due to the weakness of the couplings and the large masses of the Higgs bosons to be produced. However, in the favorable regions, which are generally characterized by large values of $\tan\beta$ and of $\tan\alpha$, one may collect a few hundred events of the type ($\Pi$) – mainly from $h^0$ – in the high luminosity phase of the LHC. As we have said, this is basically due to the enormous enhancement that may undergo the FCNC decay rates, but also because in the same regions of parameter space where the $BR$’s are enhanced, also the LHC production rates of the Higgs bosons can be significantly larger (one order of magnitude) in the 2HDM II as compared to the SM.

Interestingly enough, in many cases one can easily distinguish whether the enhanced FCNC events ($\Pi$) stem from the dynamics of a general, unrestricted, 2HDM model, or rather from some supersymmetric mechanisms within the MSSM. This is already obvious from the fact that the ranges of neutral and charged Higgs boson masses in the 2HDM case can be totally incompatible with the corresponding ones in the MSSM. But there are many other ways to discriminate these rare events. For instance, in the 2HDM case the CP-odd modes $A^0 \rightarrow t\bar{c} + \bar{t}c$ are completely hopeless whereas in the MSSM they can be enhanced $[8,15,16,18]$. Using this information in combination with the masses of potentially detected Higgs bosons could be extremely useful to pinpoint the supersymmetric or non-supersymmetric nature of them. We may describe a few specific strategies. As it was first shown in Ref. $[8]$, the leading SUSY-FCNC effects associated to the $h t c$ vertices ($h = h^0, H^0, A^0$) come from the FCNC gluino interactions which are induced.

\*\*\* Misidentification of $b$-quarks as $c$-quarks in $tb$ production might be a source of background to our FCNC events. However, to rate the actual impact of that misidentification one would need a dedicated simulation of the signal versus background, which is beyond the scope of this paper.\*\*
by potentially large misalignments of the quark and squark mass matrices. These effects are not particularly sensitive to $\tan \beta$ and they can be very sizeable for both high and moderately low values of this parameter. This sole fact can be another distinguishing feature between FCNC events of MSSM or 2HDM origin. If, for example, a few of these events were observed and at the same time the best MSSM fits to the electroweak precision data would favor moderate values of $\tan \beta$, say in the range $10 - 20$, then it is clear that those events could originate in the FCNC gluino interactions but in no way within the context of the general 2HDM. In this respect it should be mentioned that the FCNC gluino couplings recently became more restricted from the low-energy meson data, and will presumably become further restricted in the near future. The reason being that the same couplings are related, via $SU(2)$ gauge invariance and CKM rotation, to those affecting the down-like quark sector, which will most likely become constrained by the increasingly more precise low-energy meson physics. In that circumstance the only source of FCNC Higgs decays in the MSSM will stem purely from the electroweak interactions within the super-CKM basis. Then, in the absence of these SUSY-QCD FCNC effects, we could judiciously conclude from the work of Ref. $[8]$ – in which both the SUSY-QCD and the SUSY electroweak contributions were computed for the $h t c$ vertices – that the FCNC rates in the MSSM should diminish dramatically (two to three orders of magnitude). In such case we can imagine the following “provocative” scenario. Suppose that the LHC finds a light neutral Higgs boson of mass $\lesssim 140 GeV$ (sufficiently enough, in a mass range near the MSSM upper bound for $m_{h^0}$!) and subsequently, or about simultaneously, a charged Higgs boson and another neutral Higgs boson both with masses around $400 GeV$ or more. At this point one could naively suspect that a MSSM picture out these findings is getting somehow confirmed. If, however, later on a few FCNC events are reported and potentially ascribed to the previously discovered heavy neutral Higgs boson (presumably $H^0$), then the overall situation could not correspond at all to the MSSM, while it could be perfectly compatible with the 2HDM II. Alternatively, suppose that the FCNC gluino couplings were not yet sufficiently restricted, but (still following the remaining hypotheses of the previous example) a third neutral Higgs boson (presumably $A^0$) is found, also accompanied with a few FCNC events. Then this situation would be incompatible with the 2HDM II, and in actual fact it would put forward strong (indirect) evidence of the MSSM!!

We should also mention that there are other FCNC Higgs decay modes, as for example $h \rightarrow b \bar{s} + b \bar{s}$, which could be, in principle, competitive with the top quark modes. In some cases these bottom modes can be highly enhanced in the MSSM case. Actually, a more complete assessment of the FCNC bottom modes in the MSSM case – namely one which takes also into account the supersymmetric contributions to the highly restrictive radiative B-meson decays – shows that they are eventually rendered at a similar level of the top modes under study in most of the parameter space.

To summarize, the FCNC decays of the Higgs bosons into top quark final states can be a helpful complementary strategy to search for signals of physics beyond the SM in the LHC. Our comprehensive numerical analysis shows that the FCNC studies are feasible for CP-even Higgs masses up to about $500 GeV$. While the statistics of these FCNC decays
is of course poor, the advantage is that a few tagged and well discriminated events of this sort could not be attributed by any means to the SM, and therefore should call for various kinds of new physics. In this paper we have shown that a general 2HDM II is potentially competitive to be ultimately responsible for these FCNC decays, if they are ever found, and we have exemplified how to discriminate this possibility from the more restricted one associated to the MSSM.

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$\text{BR}^{\Pi}(H^0 \rightarrow t\bar{c}+t\bar{c})$

$|\delta \rho^{2\text{HDM}}| < 3\sigma$
(b)