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Loop Induced Flavor Changing Neutral Decays of the Top Quark in a General Two-Higgs-Doublet Model

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ABSTRACT

Decays of the top quark induced by flavor changing neutral currents (FCNC) are known to be extremely rare events within the Standard Model. This is so not only for the decay modes into gauge bosons, but most notably in the case of the Higgs channels, e.g. $t \to H_{SM} + c$, with a branching fraction of 10^{-13} at most. Therefore, detection of FCNC top quark decays in a future high-energy, and high-luminosity, machine like the LHC or the LC would be an indisputable signal of new physics. In this paper we show that within the simplest extension of the SM, namely the general two-Higgs-doublet model, the FCNC top quark decays into Higgs bosons, $t \to (h^0, H^0, A^0) + c$, can be the most favored FCNC modes – comparable or even more efficient than the gluon channel $t \to g + c$. In both cases the optimal results are obtained for Type II models. However, only the Higgs channels can have rates reaching the detectable level (10^{-5}), with a maximum of order 10^{-4} which is compatible with the charged Higgs bounds from radiative B-meson decays. We compare with the previous results obtained in the Higgs sector of the MSSM.

1 Introduction

In the near and middle future, with the upgrades of the Tevatron (Run II, TeV33), the advent of the LHC, and the construction of an e^+e^- linear collider (LC), new results on top quark physics, and possibly also on Higgs physics, will be obtained that may be extremely helpful to complement the precious information already collected at LEP 100 and 200 from Z and W physics. Both types of machines, the hadron colliders and the LC will work at high luminosities and will produce large amounts of top quarks. In the LHC, for example, the production of top quark pairs will be $\sigma(t\bar{t}) = 800\,pb$ – roughly two orders of magnitude larger than in the Tevatron Run II. In the so-called low-luminosity phase ($10^{33}\,cm^{-2}s^{-1}$) of the LHC one expects about three $t\,\bar{t}$-pair per second (ten million $t\,\bar{t}$-pairs per year!) [1]. And this number will be augmented by one order of magnitude in the high-luminosity phase ($10^{34}\,cm^{-2}s^{-1}$). As for a future LC running at e.g. $\sqrt{s} = 500\,GeV$, one has a smaller cross-section $\sigma(t\bar{t}) = 650\,fb$ but a higher luminosity factor ranging from $5 \times 10^{33}\,cm^{-2}s^{-1}$ to $5 \times 10^{34}\,cm^{-2}s^{-1}$ and of course a much cleaner environment [2]. With datasets from LHC and LC increasing to several $100\,fb^{-1}$/year in the high-luminosity phase, one should be able to pile up an enormous wealth of statistics on top quark decays. Therefore, not surprisingly, these machines should be very useful to analyze rare decays of the top quark, viz. decays whose branching fractions are so small ($\lesssim 10^{-5}$) that they could not be seen unless the number of collected top decays is very large.

The reason for the interest in these decays is at least twofold. First, the typical branching ratios for the rare top quark decays predicted within the Standard Model (SM) are so small that the observation of a single event of this kind should be "instant evidence", so to speak, of new physics; and second, due to its large mass ($m_t = 174.3 \pm 5.1\,GeV$ [3]), the top quark could play a momentous role in the search for Higgs physics beyond the SM. While this has been shown to be the case for the top quark decay modes into charged Higgs bosons, both in the Minimal Supersymmetric Standard Model (MSSM) and in a general two-Higgs-doublet model (2HDM) [4, 5][1], we expect that a similar situation could apply for top quark decays into non-SM neutral Higgs bosons. Notice that the latter are rare top quark events of a particularly important kind: they are decays of the top quark mediated by flavor changing neutral currents (FCNC).

At the tree-level there are no FCNC processes in the SM, and at one-loop they are induced by charged-current interactions, which are GIM-suppressed [7]. In particular, FCNC decays of the top quark into gauge bosons ($t \to c\,V$; $V \equiv \gamma, Z, g$) are very unlikely. For the present narrow range of masses for the top quark, they yield maximum branching ratios of $\sim 5 \times 10^{-13}$ for the photon, slightly above 1×10^{-13} for the Z-boson, and $\sim 4 \times 10^{-11}$ for the gluon channel [8]. These are much smaller than the FCNC rates of a typical low-energy meson decay, e.g. $B(b \to s\,\gamma) \sim 10^{-4}$. And the reason is simple: for FCNC top quark decays in the SM, the loop amplitudes are controlled by down-type quarks, mainly by the bottom quark. Therefore, the scale of the loop amplitudes is set by m_b^2 and the partial widths are of order

$$\Gamma(t \to V\,c) \sim |V_{tb}^* V_{bc}|^2 \alpha\, G_F^2\, m_t\, m_b^4\, F \sim |V_{bc}|^2 \alpha_{em}^2 \alpha\, m_t \left(\frac{m_b}{M_W}\right)^4 F, \tag{1}$$

[1]For a recent review of the main features of loop-induced supersymmetric effects on top quark production an decay, see e.g. Ref. [6].

where α is α_{em} for $V = \gamma, Z$ and α_s for $V = g$. The factor $F \sim (1 - m_V^2/m_t^2)^2$ results, upon neglecting m_c, from phase space and polarization sums. Notice that the dimensionless fourth power mass ratio, in parenthesis in eq. (1), stems from the GIM mechanism and is responsible for the ultralarge suppression beyond naive expectations based on pure dimensional analysis, power counting and CKM matrix elements. From that simple formula, the approximate orders of magnitude mentioned above ensue immediately.

Even more dramatic is the situation with the top quark decay into the SM Higgs boson, $t \to c\, H_{SM}$, which has recently been recognized to be much more disfavored than originally claimed [8]: $BR(t \to c\, H_{SM}) \sim 10^{-13} - 10^{-15}$ $(m_t = 175\, GeV;\ M_Z \leq M_H \leq 2\ M_W)$ [9]. This extremely tiny rate is far out of the range to be covered by any presently conceivable high luminosity machine. On the other hand, the highest FCNC top quark rate in the SM, namely that of the gluon channel $t \to c\, g$, is still 6 orders of magnitude below the feasible experimental possibilities at the LHC. All in all the detection of FCNC decays of the top quark at visible levels (viz. $BR(t \to c\, X) > 10^{-5}$) by the high luminosity colliders round the corner (especially LHC and LC) seems doomed to failure in the absence of new physics. Thus the possibility of large enhancements of some FCNC channels up to the visible threshold, particularly within the context of the general 2HDM and the MSSM, should be very welcome. Unfortunately, although the FCNC decay modes into electroweak gauge bosons $V_{ew} = W, Z$ may be enhanced a few orders of magnitude, it proves to be insufficient to raise the meager SM rates mentioned before up to detectable limits, and this is true both in the 2HDM – where $BR(t \to V_{ew}\, c) < 10^{-6}$ [8] – and in the MSSM – where $BR(t \to V_{ew}\, c) < 10^{-7}$ except in highly unlikely regions of the MSSM parameter space [10][2]. In this respect it is a lucky fact that these bad news need not to apply to the gluon channel, which could be barely visible $(BR(t \to g\, c) \lesssim 10^{-5})$ both in the MSSM [12, 13] and in the general 2HDM [8]. But, most significant of all, they may not apply to the non-SM Higgs boson channels $t \to (h^0, H^0, A^0) + c$ either. As we shall show in the sequel, these Higgs decay channels of the top quark could lie above the visible threshold for a parameter choice made in perfectly sound regions of parameter space!

While a systematic discussion of these "gifted" Higgs channels was made in Ref. [13] for the MSSM case and in other models[3], to the best of our knowledge there is no similar study in the general 2HDM. And we believe that this study is necessary, not only to assess what are the chances to see traces of new (renormalizable) physics in the new colliders round the corner but also to clear up the nature of the virtual effects; in particular to disentangle whether the origin of the hypothetically detected FCNC decays of the top quark is ultimately triggered by SUSY or by some alternative, more generic, renormalizable extension of the SM such as the 2HDM or generalizations thereof. Of course the alleged signs of new physics could be searched for directly through particle tagging, if the new particles were not too heavy. However, even if accessible, the corresponding signatures could be far from transparent. In contrast, the indirect approach based on the FCNC processes has the advantage that one deals all the time with the dynamics of the top

[2]Namely, regions in which there are wave-function renormalization thresholds due to (extremely fortuitous!) sharp coincidences between the sum of the sparticle masses involved in the self-energy loops and the top quark mass. See e.g. Ref. [11] for similar situations already in the conventional $t \to W\, b$ decay within the MSSM. In our opinion these narrow regions should not be taken too seriously.

[3]Preliminary SUSY analysis of the Higgs channels are given in [14], but they assume the MSSM Higgs mass relations at the tree-level. Therefore these are particular cases of the general MSSM approach given in [13]. Studies beyond the MSSM (e.g. including R-parity violation) and also in quite different contexts from the present one (for instance in effective field theories) are available in the literature, see e.g. [15].

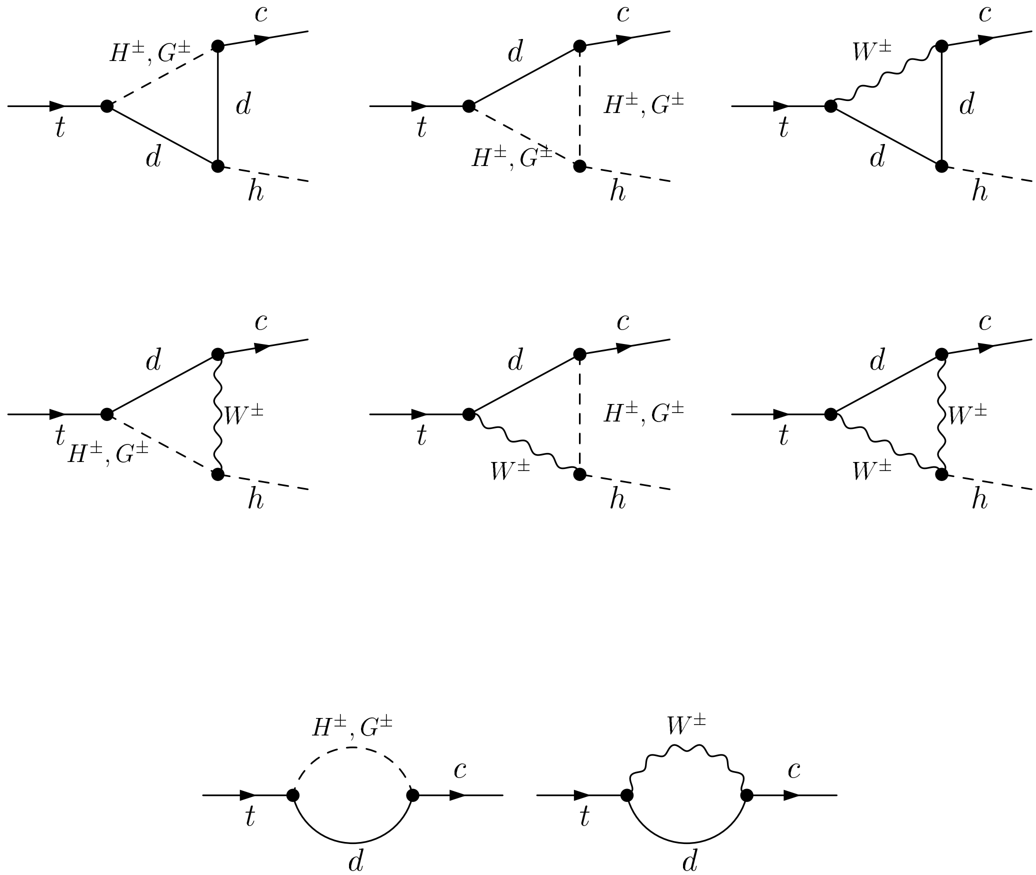

Figure 1: One-loop vertex diagrams contributing to the FCNC top quark decays (2). Shown are the vertices and mixed self-energies with all possible contributions from the SM fields and the Higgs bosons from the general 2HDM. The Goldstone boson contributions are computed in the Feynman gauge.

quark. Thus by studying potentially new features beyond the well-known SM properties of this quark one can hopefully uncover the existence of the underlying new interactions.

2 Relevant fields and interactions in the 2HDM

We will mainly focus our interest on the loop induced FCNC decays

$$t \to c\, h \quad (h = h^0, H^0, A^0)\,, \tag{2}$$

in which any of the three possible neutral Higgs bosons from a general 2HDM can be in the final state. However, as a reference we shall compare throughout our analysis the Higgs channels with the more conventional gluon channel

$$t \to c\, g\,. \tag{3}$$

Although other quarks could participate in the final state of these processes, their contribution is obviously negligible – because it is further CKM suppressed. The lowest order diagrams entering these decays are one-loop diagrams in which Higgs, quarks, gauge and

Goldstone bosons – in the Feynman gauge – circulate around. While the diagrams for the decays (2) are depicted in Fig. 1, the ones for the decay (3) are not shown [8]. Here we follow the standard notation [16], namely h^0, H^0 are CP-even Higgs bosons (where by convention $m_{h^0} < m_{H^0}$) and A^0 is a CP-odd one (often called a "pseudoscalar"). As it is well-known, when the quark mass matrices are diagonalized in non-minimal extensions of the Higgs sector of the SM, the Yukawa couplings do not in general become simultaneously diagonalized, so that one would expect Higgs mediated FCNC's at the tree-level. These are of course unwanted, since they would lead to large FCNC processes in light quark phenomenology, which are stringently restricted by experiments. Apart from letting the Higgs masses to acquire very large values, one has two additional, more elegant, canonical choices to get rid of them. In Type I 2HDM (also denoted 2HDM I) one Higgs doublet, Φ_1, does not couple to fermions at all and the other Higgs doublet, Φ_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, Φ_1, couples to down quarks (but not to up quarks) while Φ_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 if we impose a discrete symmetry, e.g. $\Phi_1 \to -\Phi_1$ and $\Phi_2 \to +\Phi_2$ (or vice versa) plus a suitable transformation for the right-handed quark fields. We shall not worry here on the ultimate theoretical origin of this ad hoc symmetry, but we will accept it as a guiding principle. As mentioned above, the SUSY case has recently been investigated in Ref. [13], so we wish to concentrate here on Type I and Type II models of a sufficiently generic nature, to wit, those which are characterized by the following set of free parameters:

$$(m_{h^0}, m_{H^0}, m_{A^0}, m_{H^\pm}, \tan\alpha, \tan\beta)\,, \tag{4}$$

where $m_{H^\pm}$ is the mass of the charged Higgs companions $H^\pm$, $\tan\alpha$ defines the mixing angle α in the diagonalization of the CP-even sector, and $\tan\beta$ gives the mixing angle β in the CP-odd sector. The latter is a key parameter in our analysis. It is given by the quotient of the vacuum expectation values (VEV's) of the two Higgs doublets $\Phi_{2,1}$, viz. $\tan\beta = v_2/v_1$, where the parameter sum $v^2 \equiv v_1^2 + v_2^2$ is fixed by the W mass: $M_W^2 = (1/2)\, g^2\, v^2$ (g is the weak $SU(2)$ gauge coupling) or, equivalently, by the Fermi constant: $G_F = 1/(2\sqrt{2})\, v^2$. It is well-known [16] that the most general 2HDM Higgs potential subject to hermiticity, $SU(2) \times U(1)$ gauge invariance and a discrete symmetry of the sort mentioned above involves six scalar operators with six free (real) coefficients $\lambda_i\, (i = 1, 2, \ldots, 6)$ and the two VEV's[4]. We will furthermore assume that $\lambda_5 = \lambda_6$ in the general 2HDM Higgs potential [18]. This allows to remove the CP phase in the potential by redefining the phase of one of the doublets. In this way one can choose the VEV's of $\Phi_{1,2}$ real and positive. The alternative set (4) is just a (more physical) reformulation of this fact after diagonalization of the mass matrices and imposing the aforementioned set of constraints.

As stated above, the two canonical types of 2HDM's only differ in the couplings to fermions but they share Feynman rules generally different from the corresponding ones

[4]By imposing the discrete symmetry one is able, in principle, to get rid of two additional quartic and one bilinear operators in the Higgs potential. These quartics are not shown in [16], but the most general 2HDM potential could contain all of these additional terms [17]. The bilinear ones are eventually kept in most cases as one usually makes allowance for the discrete symmetry to be only (softly) violated by the dimension two operators. The resulting model is still a minimal setup compatible with the absence of Higgs-mediated dangerous FCNC.

in the MSSM [16]. The necessary Feynman rules to compute the diagrams in Fig. 1 are mostly well-known. For instance, the interactions between Higgs and Goldstone bosons and the interactions between Goldstone bosons and gauge bosons are the same as in the MSSM [16]. Similarly for the interactions among Goldstone bosons and gauge bosons, and Goldstone bosons and fermions. However, the 2HDM Feynman rules for the charged and neutral Higgs interactions with fermions and also for the trilinear self-couplings of Higgs bosons may drastically deviate from the MSSM. The charged Higgs interactions with fermions are encoded in the following interaction Lagrangian

$$\mathcal{L}_{Htb}^{(j)} = \frac{g V_{tb}}{\sqrt{2}\, M_W}\, H^- \bar{b}\, [m_t \cot\beta\, P_R + m_b\, a_j\, P_L]\; t + h.c. \tag{5}$$

where here, and hereafter, we use third-quark-family notation as a generic one; V_{tb} is the corresponding CKM matrix element, $P_{L,R} = (1/2)(1 \mp \gamma_5)$ are the chiral projection operators on left- and right-handed fermions, $j = I, II$ runs over Type I and Type II 2HDM's, and we have introduced a parameter a_j such that $a_I = -\cot\beta$ and $a_{II} = +\tan\beta$. For the neutral Higgs interactions, the necessary pieces of the Lagrangian (see Fig. 1) can be written in the compact form

$$\begin{aligned} \mathcal{L}_{hqq}^{(j)} &= \frac{-g\, m_b}{2\, M_W \left\{ \begin{array}{c} \sin\beta \\ \cos\beta \end{array} \right\}}\, \bar{b} \left[h^0 \left\{ \begin{array}{c} \cos\alpha \\ -\sin\alpha \end{array} \right\} + H^0 \left\{ \begin{array}{c} \sin\alpha \\ \cos\alpha \end{array} \right\} \right] b + \frac{i\, g\, m_b\, a_j}{2\, M_W}\, \bar{b}\, \gamma_5\, b\, A^0 \\ &+ \frac{-g\, m_t}{2\, M_W\, \sin\beta}\, \bar{t} \left[h^0\, \cos\alpha + H^0\, \sin\alpha \right] t + \frac{i\, g\, m_t}{2\, M_W \tan\beta}\, \bar{t}\, \gamma_5\, t\, A^0 \,, \end{aligned} \tag{6}$$

where the upper row is for $j = I$ and the down row is for $j = II$. As far as it goes to the 2HDM Feynman rules for the trilinear couplings among Higgses, and Higgses with Goldstone bosons, they are summarized in Table 1, and are valid for Type I and Type II models. Let us recall that Type II models are specially important in that the Higgs sector of the MSSM is precisely of this sort. Had we not imposed the restriction $\lambda_5 = \lambda_6$, then the trilinear rules would be explicitly dependent on the λ_5 parameter. However the numerical analysis that we perform in the next section does not depend in any essential way on this simplification. In essence we have just traded λ_5 for $m_{A^0}^2$ in these rules and so by varying with respect to m_{A^0} we do explore most of the quantitative potential of the general 2HDM. In the MSSM case the condition $\lambda_5 = \lambda_6$ is automatic by the underlying supersymmetry, and the values of these couplings are determined by the $SU(2) \times U(1)$ electroweak gauge couplings. But of course we treat here the parameters of Type II models in a way not restricted by SUSY prejudices. On the other hand there is no need to depart arbitrarily from the SUSY frame, as it can be useful for a better comparison. Be as it may, our Type II Higgs model is still sufficiently general that it cannot be considered as the limit of the MSSM when all the sparticle masses are decoupled. Both in the generic 2HDM II and in the MSSM, the Feynman rules for the lightest CP-even Higgs, h^0, go over to the SM Higgs boson ones in the limit $\sin(\beta - \alpha) \to 1$. In the particular case of the MSSM, but *not* in a general 2HDM II, this limit is equivalent to $m_{A^0} \to \infty$. Moreover, in the MSSM, $m_{h^0} \lesssim 135\, GeV$ [19] whereas in the general Type II model there is no upper bound on m_{h^0}, and by the same token the corresponding lower bound is considerably less stringent (see below).

Since we shall perform our calculation in the on-shell scheme, we understand that the physical inputs are given by the electromagnetic coupling and the physical masses of all

$H^\pm H^\pm H^0$	$(-ig)\left[(m^2_{H^\pm}-m^2_{A^0}+\frac{1}{2}m^2_{H^0})\sin(2\beta)\cos(\beta-\alpha)+ +(m^2_{A^0}-m^2_{H^0})\cos(2\beta)\sin(\beta-\alpha)\right]\frac{1}{M_W\sin(2\beta)}$
$H^\pm H^\pm h^0$	$(-ig)\left[(m^2_{H^\pm}-m^2_{A^0}+\frac{1}{2}m^2_{h^0})\sin(2\beta)\sin(\beta-\alpha)+ +(m^2_{h^0}-m^2_{A^0})\cos(2\beta)\cos(\beta-\alpha)\right]\frac{1}{M_W\sin(2\beta)}$
$H^\pm H^\pm A^0$	0
$H^\pm G^\pm H^0$	$(-ig)(m^2_{H^\pm}-m^2_{H^0})\frac{\sin(\beta-\alpha)}{2M_W}$
$H^\pm G^\pm h^0$	$ig(m^2_{H^\pm}-m^2_{h^0})\frac{\cos(\beta-\alpha)}{2M_W}$
$H^\pm G^\pm A^0$	$\pm g\,\frac{(m^2_{H^\pm}-m^2_{A^0})}{2M_W}$
$G^\pm G^\pm H^0$	$(-ig)\frac{m^2_{H^0}\cos(\beta-\alpha)}{2M_W}$
$G^\pm G^\pm h^0$	$(-ig)\frac{m^2_{h^0}\sin(\beta-\alpha)}{2M_W}$
$G^\pm G^\pm A^0$	0

Table 1: Feynman rules for the trilinear couplings involving the Higgs self-interactions and the Higgs and Goldstone boson vertices in the Feynman gauge. The rules are common to both Type I and Type II 2HDM under the conditions explained in the text. We have singled out some null entries associated to CP violation. These null vertices imply the corresponding deletion of some vertex diagrams in Fig. 1.

the particles:

$$(e, M_W, M_Z, m_{h^0}, m_{H^0}, m_{A^0}, m_{H^\pm}, m_f)\,. \tag{7}$$

The remaining parameters, except the Higgs mixing angles, are understood to be given in terms of the latter, e.g. the $SU(2)$ gauge coupling appearing in the previous formulae and in Table 1 is given by $g = e/s_w$, where the sinus of the weak mixing angle is defined through $s_w^2 = 1-M_W^2/M_Z^2$. It should be clear that, as there are no tree-level FCNC decays of the top quark, there is no need to introduce counterterms for the physical inputs in this calculation. In fact, the calculation is carried out in lowest order ("tree level") with respect to the effective tch and tcg couplings and so the sum of all the one-loop diagrams (as well as of certain subsets of them) should be finite in a renormalizable theory, and indeed it is.

3 Numerical analysis

From the previous interaction Lagrangians and Feynman rules it is now straightforward to compute the loop induced FCNC rates for the decays (2) and (3). We shall refrain from listing the lengthy analytical formulae as the computation is similar to the one reported in great detail in Ref. [13]. Therefore, we will limit ourselves to exhibit the final numerical results. The fiducial ratio on which we will apply our numerical computation is the following:

$$B^j(t \to h + c) = \frac{\Gamma^j(t \to h + c)}{\Gamma(t \to W^+ + b) + \Gamma^j(t \to H^+ + b)} \ , \tag{8}$$

for each Type $j = I, II$ of 2HDM and for each neutral Higgs boson $h = h^0$, H^0, A^0. While this ratio is not the total branching fraction, it is enough for most practical purposes and it is useful in order to compare with previous results in the literature. Notice that for $m_{H^\pm} > m_t$ (the most probable situation for Type II 2HDM's, see below) the ratio (8) reduces to $B^j(t \to h + c) = \Gamma^j(t \to h + c)/\Gamma(t \to W^+ + b)$, which is the one that we used in Ref. [13]. It is understood that $\Gamma^j(t \to h + c)$ above is computed from the one-loop diagrams in Fig. 1, with all quark families summed up in the loop. Therefore, consistency in perturbation theory requires to compute $\Gamma(t \to W^+ + b)$ and $\Gamma(t \to H^+ + b)$ in the denominator of (8) only at the tree-level (for explicit expressions see e.g. [4]). As mentioned in Sec. 2, we wish to compare our results for the Higgs channels (2) with those for the gluon channel (3), so that we similarly define

$$B^j(t \to g + c) = \frac{\Gamma^j(t \to g + c)}{\Gamma(t \to W^+ + b) + \Gamma^j(t \to H^+ + b)} \ . \tag{9}$$

We have performed a fully-fledged independent analytical and numerical calculation of $\Gamma^j(t \to g + c)$ at one-loop in the context of 2HDM I and II. Where there is overlapping, we have checked the numerical results of Ref. [8], but we point out that they agree with us only if $\Gamma(t \to H^+ + b)$ is included in the denominator of eq. (9), in contrast to what is asserted in that reference in which $B(t \to g + c)$ is defined without the charged Higgs channel contribution.

We have performed part of the analytical calculation of the diagrams for both processes (2) and (3) by hand and we have cross-checked our results with the help of the numeric and algebraic programs FeynArts, FormCalc and LoopTools [20], with which we have completed the rest of the calculation. In particular, the cancellation of UV divergences in the total amplitudes was also verified by hand. In addition we have checked explicitly the gauge invariance of the total analytical amplitude for the process (3), which is a powerful test.

As mentioned above, a highly relevant parameter is $\tan\beta$, which must be restricted to the approximate range

$$0.1 < \tan\beta \lesssim 60 \tag{10}$$

in perturbation theory[5]. It is to be expected from the various couplings involved in the processes under consideration that the low $\tan\beta$ region could be relevant for both the Type I and Type II 2HDM's. In contrast, the high $\tan\beta$ region is only potentially important for the Type II. However, the eventually relevant regions of parameter space are also determined by the value of the mixing angle α, as we shall see below.

Of course there are several restrictions that must be respected by our numerical analysis. Above all the quadratic violations of $SU(2)$ "custodial symmetry" must be within

[5]Some authors [21] claim that perturbativity allows $\tan\beta$ to reach values of order 100 and beyond, and these are still used in the literature. We consider it unrealistic and we shall not choose $\tan\beta$ outside the interval (10). Plots versus $\tan\beta$, however, will indulge larger values just to exhibit the dramatic enhancements of our FCNC top quark rates.

experimental bound. Therefore, the one-loop corrections to the ρ-parameter from the 2HDM sector cannot deviate from the reference SM contribution in more than one per mil [3]:

$$|\delta\rho^{2HDM}| \leqslant 0.001\,. \tag{11}$$

From the analytical expression for $\delta\rho$ in the general 2HDM we have introduced this numerical condition in our codes. Moreover, non-SUSY charged Higgs bosons from Type II models are subject to a very important indirect bound from radiative B-meson decays, specifically the experimental measurement by CLEO of the branching fraction $BR(B \to X_s\,\gamma)$ – equivalently $BR(b \to s\,\gamma)$ at the quark level [22]. At present the data yield:

$$BR(b \to s\,\gamma) = (3.15 \pm 0.35 \pm 0.32 \pm 0.26) \times 10^{-4}. \tag{12}$$

The charged Higgs contribution to $BR(b \to s\,\gamma)$ is positive; hence the larger is the experimental rate the smaller can be the charged Higgs boson mass. From the various analysis in the literature one finds $m_{H^\pm} > (165 - 200)\ GeV$ for virtually any $\tan\beta \gtrsim 1$ [23, 24]. This bound does not apply to Type I models because at large $\tan\beta$ the charged Higgs couplings are severely suppressed, whereas at low $\tan\beta$ we recover the previous unrestricted situation of Type II models. Therefore, in principle the top quark decay $t \to H^+ + b$ is still possible in 2HDM I; but also in 2HDM II, if $m_{H^\pm}$ lies near the lowest end of the previous bound, and in this case that decay can contribute to the denominator of eqs. (8)-(9).

One may also derive lower bounds to the neutral Higgs masses for these models [25]. For instance, one may use the Bjorken process $e^+e^- \to Z + h^o$ and the associated Higgs boson pair production $e^+e^- \to h^0(H^0) + A^0$ to obtain the following bounds in most of the parameter space: $m_{h^0} + m_{A^0} \gtrsim 100\,GeV$ or $\gtrsim 150\,GeV$ depending on whether we accept any value of $\tan\beta$ or we impose $\tan\beta > 1$ respectively [26]. In each of these cases there is a light mass corner in parameter space both in the CP-even and in the CP-odd mass ranges around $m_{h,^0A^0} = 20 - 30\ GeV$ [26]. Notwithstanding, as it is shown by the fit analysis of precision electroweak data in Ref. [24], in the large $\tan\beta$ region a light h^0 is statistically correlated with a light $H^\pm$, so that this situation is not favored by the aforementioned bound from $b \to s\,\gamma$. Moreover, since our interest in Type II models is mainly focused in the large $\tan\beta$ regime, the corner in the light CP-even mass range is a bit contrived. At the end of the day one finds that, even in the worst situation, the strict experimental limits still allow generic 2HDM neutral scalar bosons as light as $70\,GeV$ or so. As we said, most of these limits apply to Type II 2HDM's, but we will conservatively apply them to Type I models as well.

Finally, for both models we have imposed the condition that the (absolute value) of the trilinear Higgs self-couplings do not exceed the maximum unitarity limit tolerated for the SM trilinear coupling:

$$|\lambda_{HHH}| \leqslant \left|\lambda^{(SM)}_{HHH}(m_H = 1\,TeV)\right| = \frac{3\,g\,(1\,TeV)^2}{2\,M_W}\,. \tag{13}$$

The combined set of independent conditions turns out to be quite effective in narrowing down the permitted region in the parameter space, as can be seen in Figs. 2-5 where we plot the fiducial FCNC rates (8)-(9) versus the parameters (4). The cuts in some of these curves just reflect the fact that at least one of these conditions is not fulfilled.

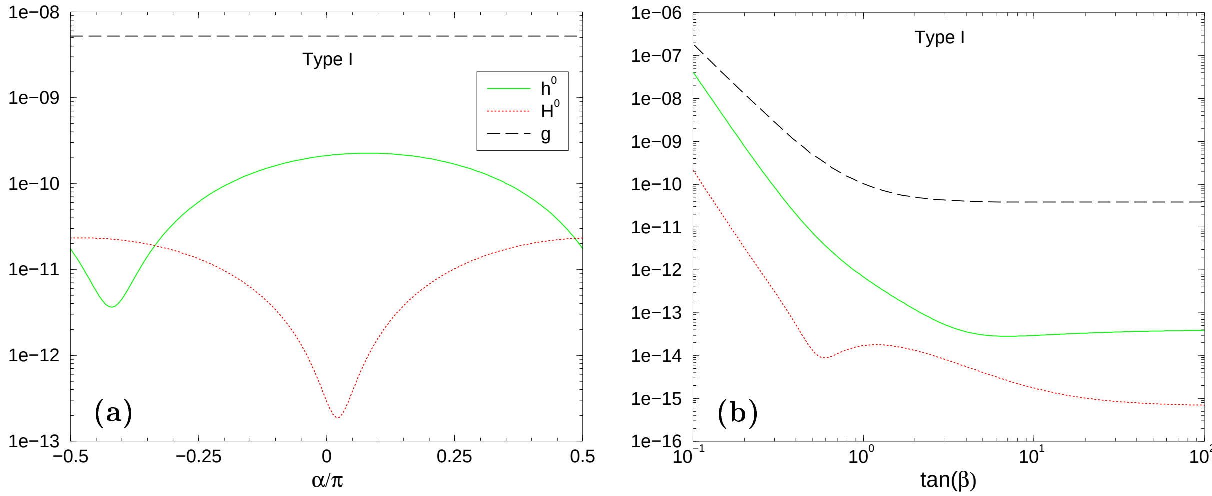

Figure 2: Evolution of the FCNC top quark fiducial ratios (8)-(9) in Type I 2HDM versus: **(a)** the mixing angle α in the CP-even Higgs sector, in units of π; **(b)** $\tan\beta$. The values of the fixed parameters are as in eqs. (14) and (15).

After scanning the parameter space, we see in Figs. 2-3 that the 2HDM I (resp. 2HDM II) prefers low values (resp. high values) of $\tan\alpha$ and $\tan\beta$ for a given channel, e.g. $t \to h^0\, c$. Therefore, the following choice of mixing angles will be made to optimize the presentation of our numerical results:

$$\begin{aligned} \text{2HDM I}: &\ \tan\alpha = \tan\beta = 1/4\ ; \\ \text{2HDM II}: &\ \tan\alpha = \tan\beta = 50\,. \end{aligned} \tag{14}$$

We point out that, for the same values of the masses, one obtains the same maximal FCNC rates for the alternative channel $t \to H^0\, c$ provided one just substitutes $\alpha \to \pi/2 - \alpha$. Equations (14) define the eventually relevant regions of parameter space and, as mentioned above, depend on the values of the mixing angles α and β, namely $\beta \simeq \alpha \simeq 0$ for Type I and $\beta \simeq \alpha \simeq \pi/2$ for Type II.

Due to the $\alpha \to \pi/2 - \alpha$ symmetry of the maximal rates for the CP-even Higgs channels, it is enough to concentrate the numerical analysis on one of them, but one has to keep in mind that the other channel yields the same rate in another region of parameter space. Whenever a mass has to be fixed, we choose conservatively the following values for both models:

$$m_{h^0} = 100\, GeV\,,\ m_{H^0} = 150\, GeV\,,\ \ m_{A^0} = m_{H^\pm} = 180\, GeV\,. \tag{15}$$

Also for definiteness, we take the following values for some relevant SM parameters in our calculation:

$$m_t = 175\, GeV\,,\ m_b = 5\, GeV\,,\ \alpha_s(m_t) = 0.11\,,\ \ V_{cb} = 0.040\,, \tag{16}$$

and the remaining ones are as in [3]. Notice that our choice of m_{A^0} prevents the decay $t \to\ A^0\, c$ from occurring, and this is the reason why it does not appear in Fig. 2. The variation of the results with respect to the masses is studied in Figs. 4-5. In particular,

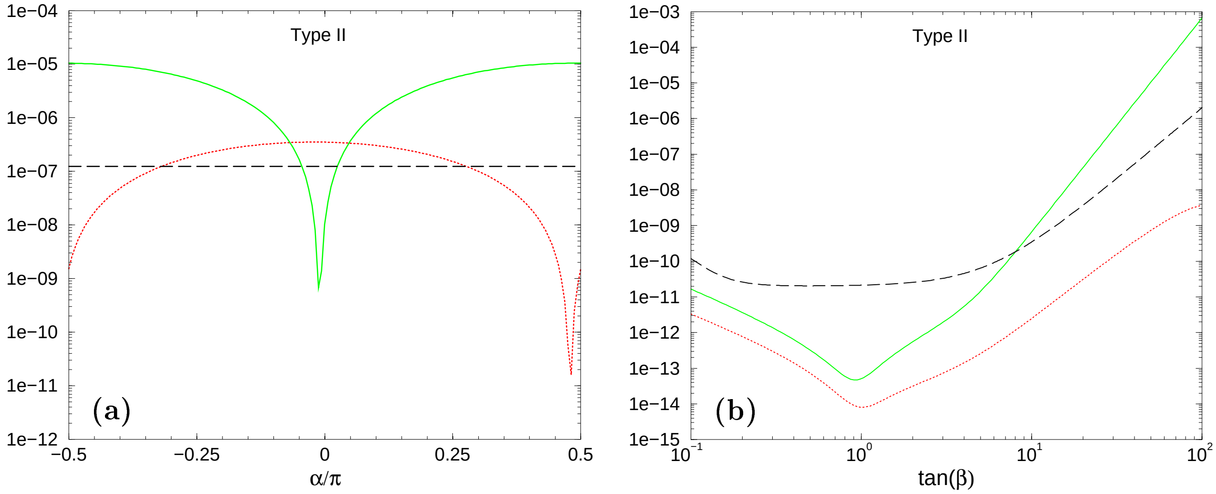

Figure 3: As in Fig. 2, but for the 2HDM II. The plot in (b) continues above the bound in eq. (10) just to better show the general trend.

in Fig. 4 we can see the (scanty) rate of the channel $t \to A^0\,c$ when it is kinematically allowed. In general the pseudoscalar channel is the one giving the skimpiest FCNC rate. This is easily understood as it is the only one that does not have trilinear couplings with the other Higgs particles (Cf. Table 1). While it does have trilinear couplings involving Goldstone bosons, these are not enhanced. The crucial role played by the trilinear Higgs self-couplings in our analysis cannot be underestimated as they can be enhanced by playing around with both (large or small) $\tan\beta$ *and* also with the mass splittings among Higgses. This feature is particularly clear in Fig. 4a where the rate of the channel $t \to h^0\,c$ is dramatically increased at large m_{A^0}, for fixed values of the other parameters, and preserving our list of constraints. Similarly would happen for $t \to H^0\,c$ in the corresponding region $\alpha \to \pi/2 - \alpha$.

From Figs. 2a and 2b it is pretty clear that the possibility to see FCNC decays of the top quark into Type I Higgs bosons is plainly hopeless even in the most favorable regions of parameter space – the lowest (allowed) $\tan\beta$ end. In fact, the highest rates remain neatly down 10^{-6}, and therefore they are (at least) one order of magnitude below the threshold sensibility of the best high luminosity top quark factory in the foreseeable future (see Section 4). We remark, in Fig. 2, that the rate for the reference decay $t \to g\,c$ in the 2HDM I is also too small but remains always above the Higgs boson rates. Moreover, for large $\tan\beta$ one has, as expected, $B^I\,(\,t \to g\,c) \to B^{SM}\,(\,t \to g\,c) \simeq 4 \times 10^{-11}$ because in this limit all of the charged Higgs couplings in the 2HDM I (the only Higgs couplings involved in this decay) drop off. Due to the petty numerical yield from Type I models we refrain from showing the dependence of the FCNC rates on the remaining parameters.

Fortunately, the meager situation just described does not replicate for Type II Higgs bosons. For, as shown in Figs. 3a and 3b, the highest potential rates are of order 10^{-4}, and so there is hope for being visible. In this case the most favorable region of parameter space is the high $\tan\beta$ end in eq. (10). Remarkably, there is no need of risking values over and around 100 (which, as mentioned above, are sometimes still claimed as perturbative!) to obtain the desired rates. But it certainly requires to resort to models whose hallmark is a large value of $\tan\beta$ of order or above $m_t/m_b \gtrsim 35$. As for the dependence of the FCNC

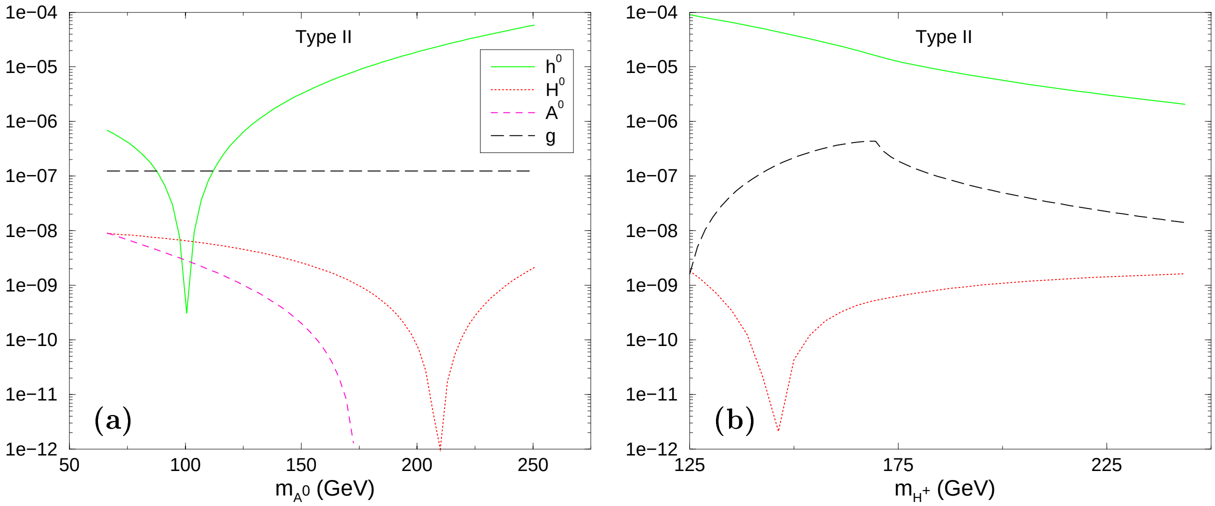

Figure 4: Evolution of the FCNC top quark fiducial ratios (8)-(9) in Type II 2HDM versus: **(a)** the CP-odd Higgs boson mass m_{A^0}; **(b)** the charged Higgs boson mass $m_{H^\pm}$. The values of the fixed parameters are as in eqs. (14) and (15). The plot in (b) starts below the bound $m_{H^\pm} > 165\,GeV$ mentioned in the text to better show the general trend.

rates on the various Higgs boson masses (Cf. Figs. 4-5) we see that for large m_{A^0} the decay $t \rightarrow h^0\,c$ can be greatly enhanced as compared to $t \rightarrow g\,c$; and of course, once again, the same happens with $t \rightarrow H^0\,c$ in the alternative region $\alpha \rightarrow \pi/2 - \alpha$. We also note (from the combined use of Figs. 3b, 4a and 4b) that in the narrow range where $t \rightarrow H^+\,b$ could still be open in the 2HDM II, the rate of $t \rightarrow h^0\,c$ becomes the more visible the larger and larger is $\tan\beta$ and m_{A^0}. Indeed, in this region one may even overshoot the 10^{-4} level without exceeding the upper bound (10) while also keeping under control the remaining constraints, in particular eq. (11). Finally, the evolution of the rates (8)-(9) with respect to the two CP-even Higgs boson masses is shown in Figs. 5a and 5b. The neutral Higgs bosons themselves do not circulate in the loops (Cf. Fig. 1) but do participate in the trilinear couplings (Cf. Table 1) and so the evolution shown in some of the curves in Fig. 5 is due to both the trilinear couplings and to the phase space exhaustion.

Turning now to the light scalar and pseudoscalar corners in parameter space mentioned above, it so happens that, after all, they prove to be of little practical interest in our case. Ultimately this is due to the quadratic Higgs boson mass differences entering $(\delta\rho)^{2HDM}$ which make very difficult to satisfy the bound (11). The reason being that for Type II models the limit $m_{H^\pm} \gtrsim 165\,GeV$ from $b \rightarrow s\,\gamma$ implies that the constraint (11) cannot be preserved in the presence of light neutral Higgses. In actual fact the analysis shows that if e.g. one fixes $m_{h^0} = 20 - 30\ GeV$, then the minimum m_{A^0} allowed by $\delta\rho$ is $100\,GeV$ and the maximum rate (8) is of order 10^{-6}. Conversely, if one chooses $m_{A^0} = 20 - 30\ GeV$, then the minimum m_{h^0} allowed by $\delta\rho$ is $120\,GeV$ and the maximum rate (8) is near 10^{-4}. Although in the last case the maximum rate is higher than in the first case, it is just of the order of the maximum rate already obtained outside the light mass corners of parameter space. On the other hand, these light mass regions do not help us in Type I models either. Even though for these models we do not have the $b \rightarrow s\,\gamma$ bound on the charged Higgs, we still have the direct LEP 200 bound $m_{H^\pm} \gtrsim 78.7\,GeV$ [27] which is of course weaker than the CLEO bound. As a consequence the $\delta\rho$ constraint can be satisfied in the 2HDM I for

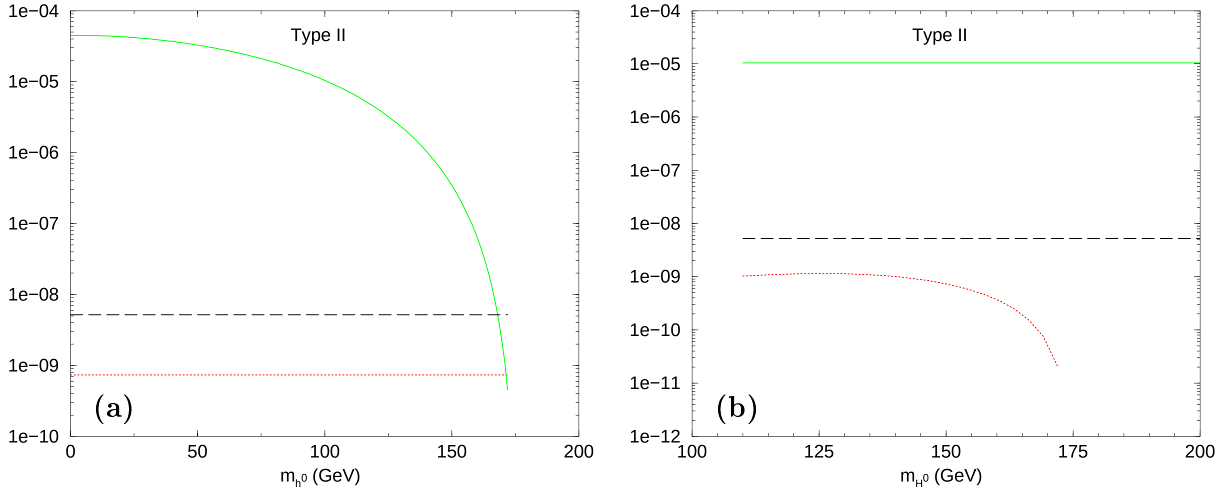

Figure 5: As in Fig. 4, but plotting versus: **(a)** the lightest CP-even Higgs boson mass m_{h^0}; **(b)** the heaviest CP-even Higgs boson mass m_{H^0}.

neutral Higgs bosons lighter than in the corresponding 2HDM II case, and one does get some enhancement of the FCNC rates. Specifically, one may reach up to 10^{-6}. However, the maximal rates (8) for the 2HDM I Higgs bosons are so small (see Figs. 2a-2b) that this order of magnitude enhancement is rendered immaterial. The upshot is that the top quark FCNC processes are not especially sensitive to the potential existence of a very light Higgs boson in either type of 2HDM.

4 Discussion and conclusions

The sensitivities to FCNC top quark decays for $100\,fb^{-1}$ of integrated luminosity in the relevant colliders are estimated to be [28]:

$$\begin{aligned} \mathbf{LHC}: & B(t \to c\,X) \gtrsim 5 \times 10^{-5} \\ \mathbf{LC}: & B(t \to c\,X) \gtrsim 5 \times 10^{-4} \\ \mathbf{TEV33}: & B(t \to c\,X) \gtrsim 5 \times 10^{-3}\ . \end{aligned} \tag{17}$$

From these experimental expectations and our numerical analysis it becomes patent that whilst the Tevatron will remain essentially blind to this kind of physics the LHC and the LC will have a significant potential to observe FCNC decays of the top quark beyond the SM. Above all there is a possibility to pin down top quark decays into neutral Higgs particles, eq. (2), within the framework of the general 2HDM II provided $\tan\beta \gtrsim m_t/m_b \sim 35$. The maximum rates are of order 10^{-4} and correspond to the two CP-even scalars. This conclusion is remarkable from the practical (quantitative) point of view, and also qualitatively because the top quark decay into the SM Higgs particle is, in notorious contradistinction to the 2HDM II case, the less favorable top quark FCNC rate in the SM. On the other hand, we deem practically hopeless to see FCNC decays of the top quark in a general 2HDM I for which the maximum rates are of order 10^{-7}. This order of magnitude cannot be enhanced unless one allows $\tan\beta \ll 0.1$, but the latter possibility is unrealistic because perturbation theory breaks down and therefore one cannot make any prediction within our approach.

We have made a parallel numerical analysis of the gluon channel $t \to c\, g$ in both types of 2HDM's. We confirm that this is another potentially important FCNC mode of the top quark in 2HDM extensions of the SM [8] but, unfortunately, it still falls a bit too short to be detectable. The maximum rates for this channel lie below 10^{-6} in the 2HDM I (for $\tan\beta > 0.1$) and in the 2HDM II (for $\tan\beta < 60$), and so it will be hard to deal with it even at the LHC.

We are thus led to the conclusion that the Higgs channels (2), more specifically the CP-even ones, give the highest potential rates for top quark FCNC decays in a general 2HDM II. Most significant of all: they are *the only* FCNC decay modes of the top quark, within the simplest renormalizable extensions of the SM, that have a real chance to be seen in the next generation of high energy, high luminosity, colliders.

The former conclusions are similar to the ones derived in Ref. [13] for the MSSM case, but there are some conspicuous differences on which we wish to elaborate a bit in what follows [29]. First, in the general 2HDM II the two channels $t \to (h^0, H^0)\; c$ give the same maximum rates, provided we look at different (disjoint) regions of the parameter space. The $t \to A^0\; c$ channel is, as mentioned, negligible with respect to the CP-even modes. Hereafter we will discard this FCNC top quark decay mode from our discussions within the 2HDM context. On the other hand, in the MSSM there is a most distinguished channel, viz. $t \to h^0\; c$, which can be high-powered by the SUSY stuff all over the parameter space. In this framework the mixing angle α becomes stuck once $\tan\beta$ and the rest of the independent parameters are given, and so there is no possibility to reconvert the couplings between h^0 and H^0 as in the 2HDM. Still, we must emphasize that in the MSSM the other two decays $t \to H^0\; c$ and $t \to A^0\; c$ can be competitive with $t \to h^0\; c$ in certain portions of parameter space. For example, $t \to H^0\; c$ becomes competitive when the pseudoscalar mass is in the range $110\, GeV < m_{A^0} \lesssim 170\, GeV$ [13]. The possibility of having more than one FCNC decay (2) near the visible level is a feature which is virtually impossible in the 2HDM II. Second, the reason why $t \to h^0\; c$ in the MSSM is so especial is that it is the only FCNC top quark decay (2) which is always kinematically open throughout the whole MSSM parameter space, while in the 2HDM all of the decays (2) could be, in the worse possible situation, dead closed. Nevertheless, this is not the most likely situation in view of the fact that all hints from high precision electroweak data seem to recommend the existence of (at least) one relatively light Higgs boson [27, 30]. This is certainly an additional motivation for our work, as it leads us to believe that in all possible (renormalizable) frameworks beyond the SM, and not only in SUSY, we should expect that at least one FCNC decay channel (2) could be accessible. Third, the main origin of the maximum FCNC rates in the MSSM traces back to the tree-level FCNC couplings of the gluino [13]. These are strong couplings, and moreover they are very weakly restrained by experiment. In the absence of such gluino couplings, or perhaps by further experimental constraining of them in the future, the FCNC rates in the MSSM would boil down to just the electroweak (EW) contributions, to wit, those induced by charginos, squarks and also from SUSY Higgses. The associated SUSY-EW rate is of order 10^{-6} at most, and therefore it is barely visible, most likely hopeless even for the LHC. In contrast, in the general 2HDM the origin of the contributions is purely EW and the maximum rates are two orders of magnitude higher than the full SUSY-EW effects in the MSSM. It means that we could find ourselves in the following situation. Suppose that the FCNC couplings of the gluino get severely restrained in the future and that we come to observe a few FCNC decays of the top quark into Higgs bosons, perhaps at the LHC and/or the LC. Then we would immediately conclude that these Higgs bosons could not be SUSY-MSSM, whilst they could perhaps be CP-even members of a 2HDM II. Fourth, the gluino effects are basically insensitive to $\tan\beta$, implying that the maximum MSSM rates are achieved equally well for low, intermediate or high values of $\tan\beta$, whereas the maximum 2HDM II rates (comparable to the MSSM ones) are attained only for high $\tan\beta$.

The last point brings about the following question: what could we possibly conclude if the gluino FCNC couplings were not further restricted by experiment and the tagging of certain FCNC decays of the top quark into Higgs bosons would come into effect? Would still be possibly to discern whether the Higgs bosons are supersymmetric or not? The answer is, most likely yes, provided certain additional conditions would be met.

There are many possibilities and corresponding strategies, but we will limit ourselves to point out some of them. For example, let us consider the type of signatures involved in the tagging of the Higgs channels. In the favorite FCNC region (14) of the 2HDM II, the combined decay $t \to h\ c \to c b\bar{b}$ is possible only for h^0 or for H^0, but not for both – Cf. Fig. 3a – whereas in the MSSM, h^0 together with H^0, are highlighted for $110\,GeV < m_{A^0} < m_t$, with no preferred $\tan\beta$ value. And similarly, $t \to A^0\ c$ is also non-negligible for $m_{A^0} \lesssim 120\,GeV$ [13]. Then the process $t \to h\ c \to c b\bar{b}$ gives rise to high p_T charm-quark jets and a recoiling $b\bar{b}$ pair with large invariant mass. It follows that if more than one distinctive signature of this kind would be observed, the origin of the hypothetical Higgs particles could not probably be traced back to a 2HDM II.

One might worry that in the case of h^0 and H^0 they could also (in principle) decay into electroweak gauge boson pairs $h^0, H^0 \to V_{ew}\overline{V}_{ew}$, which in some cases could be kinematically possible. But this is not so in practice for the 2HDM II if we stick to our favorite scenario, eq. (14). In fact, we recall that the decay $h^0 \to V_{ew}\overline{V}_{ew}$ is not depressed with respect to the SM Higgs boson case provided $\beta - \alpha = \pi/2$, and similarly for $H^0 \to V_{ew}\overline{V}_{ew}$ if $\beta - \alpha = 0$. However, neither of these situations is really pinpointed by FCNC physics because we have found $\beta \simeq \pi/2$ in the most favorable region of our numerical analysis, and moreover α was also seen there to be either $\alpha \simeq \pi/2$ (for h^0) or 0 (for H^0), so both decays $h^0, H^0 \to V_{ew}\overline{V}_{ew}$ are suppressed in the regions where the FCNC rates of the parent decays $t \to (h^0, H^0)\ c$ are maximized. Again, at variance with this situation, in the MSSM case $H^0 \to V_{ew}\overline{V}_{ew}$ is perfectly possible – not so $h^0 \to V_{ew}\overline{V}_{ew}$ due to the aforementioned upper bound on m_{h^0} – because $\tan\beta$ has no preferred value in the most favorable MSSM decay region of $t \to H^0\ c$. Therefore, detection of a high p_T charm-quark jet against a $V_{ew}\overline{V}_{ew}$ pair of large invariant mass could only be advantageous in the MSSM, not in the 2HDM. Similarly, for $\tan\beta \gtrsim 1$ the decay $H^0 \to h^0\ h^0$ (with real or virtual h^0) is competitive in the MSSM [31] in a region where the parent FCNC top quark decay is also sizeable. Again this is impossible in the 2HDM II and therefore it can be used to distinguish the two (SUSY and non-SUSY) Higgs frames.

Finally, even if we place ourselves in the high $\tan\beta$ region both for the MSSM and the 2HDM II, then the two frameworks could still possibly be separated provided that two Higgs masses were known, perhaps one or both of them being determined from the tagged Higgs decays themselves, eq. (2). Suppose that $\tan\beta$ is numerically known (from other processes or from some favorable fit to precision data), then the full spectrum of MSSM Higgs bosons would be approximately determined (at the tree level) by only knowing one Higgs mass, a fact that could be used to check whether the other measured Higgs mass becomes correctly predicted. Of course, the radiative corrections to the MSSM Higgs mass relations can be important at high $\tan\beta$ [19], but these could be taken into account from the approximate knowledge of the relevant sparticle masses obtained from the best fits available to the precision measurements within the MSSM. If there were significant departures between the predicted mass for the other Higgs and the measured one, we would probably suspect that the tagged FCNC decays into Higgs bosons should correspond to a non-supersymmetric 2HDM II.

At the end of the day we see that even though the maximum FCNC rates for the MSSM and the 2HDM II are both of order 10^{-4} – and therefore potentially visible – at some point on the road it should be possible to disentangle the nature of the Higgs model behind the FCNC decays of the top quark. Needless to say, if all the recent fuss at CERN [27] about the possible detection of a Higgs boson would eventually be confirmed, this could still be interpreted as the discovery of one neutral member of an extended Higgs model. Obviously the combined Higgs data from LEP 200 and the possible discovery of FCNC top quark decays into Higgs bosons at the LHC/LC would be an invaluable cross-check of the purportedly new phenomenology.

We emphasize our most essential conclusions in a nutshell: i) Detection of FCNC top quark decay channels into a neutral Higgs boson would be a blazing signal of physics beyond the SM; ii) There is a real chance for seeing rare events of that sort both in generic Type II 2HDM's and in the MSSM. The maximum rates for the leading FCNC processes (2) and (3) in the 2HDM II (resp. in the MSSM) satisfy the relations

$$BR(t \to g\,c) < 10^{-6}(10^{-5}) < BR(t \to h\,c) \sim 10^{-4}\,, \tag{18}$$

where it is understood that h is h^0 or H^0, but not both, in the 2HDM II; whereas h is most likely h^0, but it could also be H^0 and A^0, in the MSSM ; iii) Detection of more than one Higgs channel would greatly help to unravel the type of underlying Higgs model.

The pathway to seeing new physics through FCNC decays of the top quark is thus potentially open. It is now an experimental challenge to accomplish this program using the high luminosity super-colliders round the corner.

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