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Flavor Changing Neutral Scalar Currents at $\mu^+\mu^-$ Colliders

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Abstract: The prospect of observing the flavor changing decay $\mathcal{H} \to t\bar{c}$ of a neutral Higgs boson produced via s-channel and its subsequent decay into $t\bar{c}$ is considered at a $\mu^+\mu^-$ collider. Numerical estimates are given in the context of a two Higgs doublet model with flavor changing couplings. It is found that for many values of the model parameters such *tree-level* flavor changing decays will be produced at an observable level. In addition studies of the helicity of the top will allow the determination of the relative strengths of the flavor changing Higgs couplings and these may be measured with about 10^3 events.

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The suppression of flavor changing neutral currents (FCNC) is an important feature of the Standard Model (SM). Thus, the measurement of such currents provides an important test which can discriminate between the SM and various models of new physics. In the SM, the relative largeness of the top mass [1, 2] leads to a measurable rate of FCNC's in the down type quark sector through penguin processes [3]. In fact recent experiments at CLEO have observed the reaction $b \to s\gamma$. At least in part due to the fact that no correspondingly heavy down type quark is thought to exist, similar FCNC processes within the up sector (e.g. $t \to c\gamma$) are highly suppressed in the SM[4]. Since we do not know of a conservation law that enforces the absence of such FCNC's their continual search is clearly warranted. These considerations have, of course, fueled the searches for $\mu \to e\gamma$, $K_L \to \mu e$ etc. for a very long time. The extraordinary mass scale of the top quark has prompted many to advocate that FCNC involving the top quark may well exist [5].

An important class of models where FCNC's can occur among up type quarks are those where flavor changing occurs in an extended neutral Higgs sector. In previous works [6, 7], the observation of FCNC's (due to penguin graphs involving such a Higgs sector) was considered in the processes $t \to c\gamma$ or cZ and e^+e^- (or indeed $\mu^+\mu^-$) $\to \gamma$ or $Z \to t\bar{c}$ respectively. In this Letter we suggest that the tree level coupling of such flavor-changing neutral Higgs bosons [8] to $t\bar{c}$ may be probed by $\mu^+\mu^- \to t\bar{c}$ at suggested muon colliders (MUCs).

Although very much in the notion stage at present, the MUC has been suggested [9]-[12] as a possible lepton collider for energies in the TeV range. The advantage of such a MUC would be that the much heavier muon suffers appreciably less energy loss from synchrotron and beamstrahlung radiation. The obvious disadvantages include the fact that muons eventually decay as well as the new accelerator technology development needed to produce and control such beams to the necessary degree to reach high luminosities.

If MUCs are eventually shown to be a practical and desirable tool for exploring physics in the TeV range, most of the applications would be very similar to electron colliders. One advantage however is that they may be able to produce Higgs bosons directly in the s channel in sufficient quantity to study their properties directly [9, 13, 14, 7]. In particular, a simple but fascinating possibility that we wish to explore here is when such a Higgs, $\mathcal{H}$, has a flavor-changing $\mathcal{H}t\bar{c}$ coupling then the process $\mu^+\mu^- \to t\bar{c}$ will give a signal which should be easy to identify, is likely to take place at an observable

rate and yet has a negligible SM background. Thus the properties of this important coupling can be studied in detail.

The crucial point is that in spite of the fact that the $\mu^+\mu^-\mathcal{H}$ coupling, being proportional to m_μ, is very small, if the MUC is run on the Higgs resonance, $\sqrt{s} = m_\mathcal{H}$, Higgs bosons may be produced at an appreciable rate [9, 13, 14, 7].

At $\sqrt{s} = m_\mathcal{H}$, the cross section for producing $\mathcal{H}$, $\sigma_\mathcal{H}$, normalized to $\sigma_0 = \sigma(\mu^+\mu^- \to \gamma \to e^+e^-)$, is given by:

$$R(\mathcal{H}) = \frac{\sigma_\mathcal{H}}{\sigma_0} = \frac{3}{\alpha_e^2} B_\mu^\mathcal{H} \tag{1}$$

where $B_\mu^\mathcal{H}$ is the branching ratio of $\mathcal{H} \to \mu^+\mu^-$ and α_e is the electromagnetic coupling.

If the Higgs is very narrow, the exact tuning to the resonance implied in equation (1) may not in general be possible. Let us suppose then that the energy of the beam has a finite spread described by δ:

$$m_\mathcal{H}^2(1-\delta) < s < m_\mathcal{H}^2(1+\delta) \tag{2}$$

where we assume that s is uniform about this range. The effective rate of Higgs production will thus be given by:

$$\tilde{R}(\mathcal{H}) = \left[\frac{\Gamma_\mathcal{H}}{m_\mathcal{H}\delta} \arctan \frac{m_\mathcal{H}\delta}{\Gamma_\mathcal{H}}\right] R(\mathcal{H}) \tag{3}$$

We now consider an extended Higgs sector which admits FCNCs. In refs. [6, 7], for instance, a minimal FCNC Higgs model with two Higgs doublets ϕ_1, ϕ_2 is considered. We assume, without loss of generality, that ϕ_1 is aligned with the vev so that

$$< \phi_1 >= \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}, \quad < \phi_2 >= 0 \tag{4}$$

where $v = (\sqrt{2}G_F)^{-\frac{1}{2}}$. There are three neutral mass eigenstates denoted by H, h, and A which are [6, 7]

$$\begin{aligned} H &= \sqrt{2}[(\mathrm{Re}\phi_1^0 - v)\cos\alpha + \mathrm{Re}\phi_2^0 \sin\alpha] \\ h &= \sqrt{2}[(-\mathrm{Re}\phi_1^0 - v)\sin\alpha + \mathrm{Re}\phi_2^0 \cos\alpha] \\ A &= -\sqrt{2}\mathrm{Im}\phi_2^0 \end{aligned} \tag{5}$$

where the mixing angle α is a parameter determined by the Higgs potential.

The Lagrangian for the Higgs-fermion interaction is [6, 7]:

$$\begin{aligned}\mathcal{L} &= \lambda^U_{ij}\bar{Q}_i\tilde{\phi}_1 U_j + \lambda^D_{ij}\bar{Q}_i\phi_1 D_j + \lambda^L_{ij}\bar{L}_i\phi_1 E_j \\ &+ \xi^U_{ij}\bar{Q}_i\tilde{\phi}_2 U_j + \xi^D_{ij}\bar{Q}_i\phi_2 D_j + \xi^L_{ij}\bar{L}_i\phi_2 E_j \ + \text{h.c.}\end{aligned} \tag{6}$$

Here the $\lambda^{U,D,L}_{ij}$ couplings turn out to be proportional respectively to the quark and lepton mass matrices, while the ξ_{ij} couplings are arbitrary and flavor non-diagonal. For definiteness, we will assume that the magnitude of the parameters ξ_{ij} are as suggested by the ansatz of [15],

$$|\xi_{ij}| \approx g\frac{\sqrt{m_i m_j}}{M_w} \tag{7}$$

Let us now consider that a Higgs $\mathcal{H}$ of mass $m_{\mathcal{H}}$ is under study at a MUC. For illustrative purposes we take $\mathcal{H} = h$ in the above model where $\alpha = 0$ (case 1) or $\pi/4$ (case 2). The main distinction between the two cases is that in case 2 the decays $\mathcal{H} \to ZZ,\ WW$ are possible while in case 1 they are not. Thus case 1 is very similar to $\mathcal{H} = A$. In general the coupling of h to $f\bar{f}$ is:

$$C_{hff} = -\frac{g}{2}\frac{m_f}{m_W}\sin\alpha + \frac{\text{Re}\xi_{ff} + i\gamma_5\text{Im}\xi_{ff}}{\sqrt{2}}\cos\alpha \equiv \frac{gm_f}{2m_W}\chi_f e^{i\gamma_5\lambda_f} \tag{8}$$

while the coupling to ZZ and WW is given by:

$$C_{hZZ} = \frac{g\sin\alpha}{\cos\theta_W}m_Z g^{\mu\nu} \qquad C_{hWW} = g\sin\alpha m_W g^{\mu\nu} \tag{9}$$

Finally the flavor changing Higgs $-t\bar{c}$ coupling is given by:

$$C_{htc} = \frac{1}{\sqrt{2}}\left[\xi_{tc}P_R + \xi^\dagger_{ct}P_L\right]\cos\alpha \equiv \frac{g\sqrt{m_t m_c}}{2m_W}(\chi_R P_R + \chi_L P_L) \tag{10}$$

where χ_L and χ_R are in general complex numbers and of order unity if (7) applies.

The decay rates to these modes given the above couplings can be readily calculated at tree level by using the results that exist in the literature [16]:

$$
\begin{aligned}
\Gamma(\mathcal{H} \to t\bar{t}) &= \frac{3g_W^2 m_t^2 m_{\mathcal{H}}}{32\pi m_W^2}\beta_t\left[\beta_t^2 + (1-\beta_t^2)\sin\lambda_t\right]\chi_t^2 \\
\Gamma(\mathcal{H} \to b\bar{b}) &= \frac{3g_W^2 m_b^2 m_{\mathcal{H}}}{32\pi m_W^2}\chi_b^2 \\
\Gamma(\mathcal{H} \to ZZ) &= \frac{g^2}{128\pi}\frac{m_{\mathcal{H}}^3}{m_Z^2}\beta_Z(\beta_Z^2 + 12\frac{m_Z^4}{m_{\mathcal{H}}^4})\sin^2\alpha \\
\Gamma(\mathcal{H} \to WW) &= \frac{g^2}{64\pi}\frac{m_{\mathcal{H}}^3}{m_W^2}\beta_W(\beta_W^2 + 12\frac{m_W^4}{m_{\mathcal{H}}^4})\sin^2\alpha
\end{aligned} \tag{11}
$$

where $\beta_i = \sqrt{1 - 4m_i^2/m_{\mathcal{H}}^2}$.

The decay rate to $t\bar{c}$ is thus:

$$
\Gamma(\mathcal{H} \to t\bar{c}) = \frac{3g_W^2 m_t m_c m_{\mathcal{H}}}{32\pi m_W^2}\left(\frac{(m_{\mathcal{H}}^2 - m_t^2)^2}{m_{\mathcal{H}}^4}\right)\left(\frac{|\chi_R|^2 + |\chi_L|^2}{2}\right) \tag{12}
$$

and, $\Gamma(\mathcal{H} \to t\bar{c}) = \Gamma(\mathcal{H} \to c\bar{t})$ at the tree level that we are considering for now. The decay rate to $\mu^+\mu^-$ which we require in equation (1) is

$$
\Gamma(\mathcal{H} \to \mu^+\mu^-) = \frac{g_W^2 m_\mu^2 m_{\mathcal{H}}}{32\pi m_W^2}\chi_\mu^2; \qquad B_\mu^{\mathcal{H}} = \Gamma(\mathcal{H} \to \mu^+\mu^-)/\Gamma_{\mathcal{H}} \tag{13}
$$

For the purpose of numerical estimates let us take the following sample choices of parameters:

- Case 1: $\alpha = \lambda_c = \lambda_t = 0$, $\chi_\mu = \chi_b = \chi_t = 1$ and $\chi_L = \chi_R = 1$
- Case 2: $\alpha = \pi/4$, $\lambda_c = \lambda_t = 0$, $\chi_\mu = \chi_b = \chi_t = 1$ and $\chi_L = \chi_R = 1$

In figure 1 we plot $\tilde{R}(\mathcal{H})$ with $\delta = 0$, 10^{-3} and 10^{-2} in the two cases as well as

$$
\tilde{R}_{tc} = \tilde{R}(\mathcal{H})\,(B_{t\bar{c}}^{\mathcal{H}} + B_{c\bar{t}}^{\mathcal{H}}) \tag{14}
$$

Note that in case 1 if $m_{\mathcal{H}}$ is below the $t\bar{t}$ threshold $\tilde{R}_{tc}$ is about $.01 - 1$ and in fact tc makes up a large branching ratio. Above the $t\bar{t}$ threshold $\tilde{R}_{tc}$ drops. For case 2 the branching ratio is smaller due to the WW and ZZ threshold at about the same mass as the tc threshold and so R_{tc} is

around 10^{-3}. For a specific example if $m_{\mathcal{H}} = 300GeV$, then $\sigma_0 \approx 1pb$. For a luminosity of $10^{34}cm^{-2}s^{-1}$, a year of $10^7 s$ (1/3 efficiency) and for $\delta = 10^{-2}$ case 1 will produce about $5 \times 10^3(t\bar{c} + \bar{t}c)$ events and case 2 will produce about 150 events. Given the distinctive nature of the final state and the lack of a Standard Model background, sufficient luminosity should allow the observation of such events.

If such events are observed one would like to extract the values of χ_L and χ_R. What is measured initially at a $\mu^+\mu^-$ collider is $\tilde{R}_{tc}$. One is required to know the total width of the $\mathcal{H}$ and the energy spread of the beam in order to translate this into $\Gamma(\mathcal{H} \to t\bar{c})$. This then allows the determination of $|\chi_L|^2 + |\chi_R|^2$. To get information separately on the two couplings we note that the total helicity of the top is:

$$\mathbf{H}_t = -\mathbf{H}_{\bar{t}} = \frac{|\chi_R|^2 - |\chi_L|^2}{|\chi_R|^2 + |\chi_L|^2} \tag{15}$$

from which one may therefore infer $|\chi_L|$ and $|\chi_R|$. Unfortunately in the limit of small m_c the helicity of the c-quark is conserved hence the relative phase of χ_L and χ_R may not be determined since the two couplings do not interfere.

Of course the helicity of the t cannot be observed directly, however following the discussion of [17] one may obtain it from the decay distributions of the top. In particular if X is a particle arising in top decay let us define the forward-backwards asymmetry

$$A_X = \frac{\Gamma(\cos\theta_X > 0) - \Gamma(\cos\theta_X < 0)}{\Gamma(\cos\theta_X > 0) + \Gamma(\cos\theta_X < 0)} \tag{16}$$

where θ_X is the angle between $\vec{P}_X$ and $-\vec{P}_{\mathcal{H}}$ in the t rest frame. For each particular choice of X we define ϵ_X to be the correlation with the polarization defined by:

$$\epsilon_X = 3 < \cos\theta_X^t > \tag{17}$$

where θ_X^t is the angle between X and the spin axis of a polarized top.

In terms of ϵ_X the asymmetry A_X is thus given by:

$$A_X = \frac{1}{2}\epsilon_X \mathbf{H}_t. \tag{18}$$

Let us now consider the following decays [17]:

- 1) for $t \to Wb, W \to l^+\nu_l$ where $l = e, \mu$ then $\epsilon_l = 1$ and the branching fraction for this case is $B_1 \sim \frac{2}{9}$.
- 2) For $t \to Wb, W \to$ hadrons then $\epsilon_W = (m_t^2 - 2m_W^2)/(m_t^2 + 2m_W^2) \approx 0.39$ and the branching fraction for this is $B_2 \sim \frac{7}{9}$.

The number of $t\bar{c}$ events needed to observe the top helicity with a significance of 3-σ is [17]:

$$N_{3\sigma} = \frac{36}{\mathcal{E}_t^2 \mathbf{H}_t^2} \approx \frac{107}{\mathbf{H}_t^2} \tag{19}$$

where

$$\mathcal{E}_t = \sqrt{B_1\epsilon_l^2 + B_2\epsilon_W^2} \approx .58 \tag{20}$$

Thus at least 10^2 events are required to begin to measure the helicity of the top and hence the relative strengths of χ_L and χ_R. In the above numerical examples it is clear that for some combinations of parameters, particularly if the luminosity is $10^{34} cm^{-2} s^{-1}$, sufficient events to measure the helicity may be present.

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Figure Captions

Figure 1: The value of $\tilde{R}(\mathcal{H})$ is shown as a function of $m_{\mathcal{H}}$ for scenario 1 (dash-dot) and for scenario 2 (dots). The value of $\tilde{R}_{tc}$ is shown in case 1 for $\delta = 0$ (upper solid curve); $\delta = 10^{-3}$ (middle solid curve) and $\delta = 10^{-2}$ (lower solid curve). The value of $\tilde{R}_{tc}$ is shown in case 2 for $\delta = 0$ (upper dashed curve) and $\delta = 10^{-2}$ (lower dashed curve).

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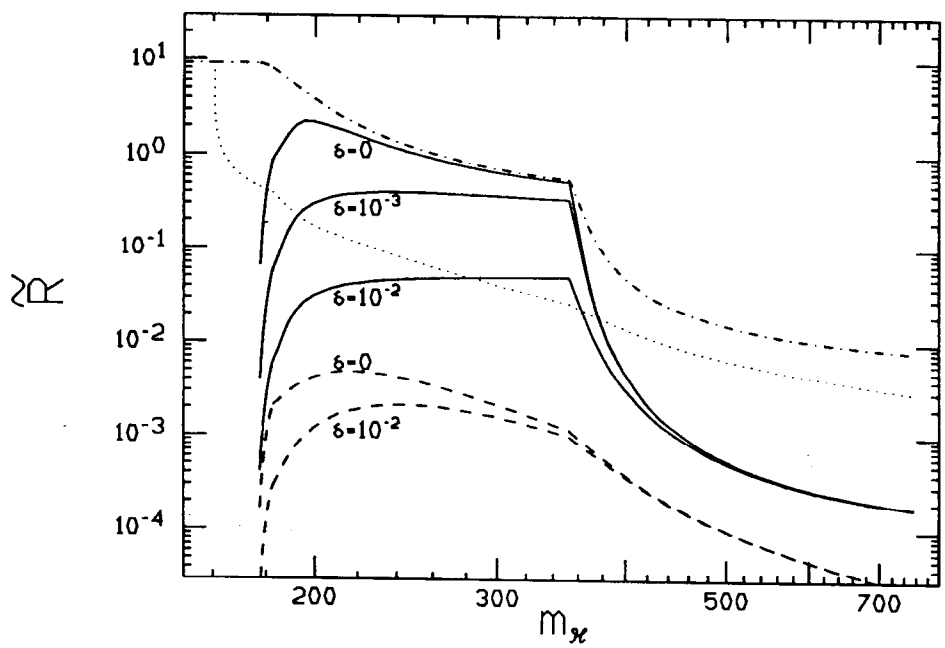
10^1
10^0
10^{-1}
10^{-2}
10^{-3}
10^{-4}
Ωh^2
δ=0
δ=10⁻³
δ=10⁻²
δ=0
δ=10⁻²
200
300
500
700
m_χ

Figure 1