

Effects of an $H - \mu - \tau$ Coupling in Quarkonium Lepton Flavor Violation Decays

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Abstract

In this work, we study the consistency of a possible nonvanishing coupling $H\mu\tau$ smaller than 0.25% as pointed by the CMS Collaboration [1], with measured lepton flavor violation processes involving quarkonium. The study we perform is in a model-independent scenario assuming only the presence of such a coupling. We show that the most promising channel to confirm this excess is to look for the lepton flavor tau decay into f_0 and μ where the experimental limit could be strongly improved with the new B factories as Belle II.

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1. INTRODUCTION

In 2015, CMS and ATLAS Collaborations have observed a slight excess of signal events with a significance of 2.4 standard deviations which can be interpreted as Higgs particles decaying into a muon and tau leptons:

$$Br(H \rightarrow \mu\tau) = \begin{cases} 0.84^{+0.39}_{-0.37} \% [2], \\ 0.77 \pm 0.62 \% [3]. \end{cases} \quad (1)$$

Recently, CMS and ATLAS Collaborations studied the upper limit on this branching ratio to obtain the following results:

$$\begin{aligned} Br(H \rightarrow \mu\tau)^{CMS} &< 0.25\% [1], \\ Br(H \rightarrow \mu\tau)^{ATLAS} &< 0.28\% [4], \end{aligned} \quad (2)$$

respectively¹. Using the reported value of Higgs mass to be 125 GeV [5], this requires the square of the coupling $H\mu\tau$ to be smaller than 1.43×10^{-3} . In new physics models, it is very challenging to generate such a lepton flavor violation without affecting other observables [6] but this value of the lepton flavor violating coupling is not in contradiction to the experimental upper limits for $\tau \rightarrow \mu\gamma$ and $\tau \rightarrow 3\mu$. Indeed, using the formalism in [7, 8], we obtain the values listed in Table 1 for the corresponding branching ratios. The expected value from $H \rightarrow \mu\tau$ is done assuming that the Higgs couplings to charged leptons are given by SM values [9, 10].

The computation of Higgs-induced lepton flavor violation (LFV) in channel decays involving only charged leptons is even more intricate as the smallness of the lepton Yukawa couplings implies that higher loop contributions can be bigger than tree-level ones [11, 8]. To avoid this problem, we shall study the effect of the LFV $H \rightarrow \tau\mu$ coupling in processes involving quarkonium. In [12], the effects of heavy sterile Majorana neutrinos in LFV decays of vector quarkonia have been studied. Sterile Majorana neutrinos induces $\gamma l_i l_j$ and $Z l_i l_j$ LFV couplings at one loop and $WW l_i l_j$ LFV couplings at tree level.

These couplings produce LFV effects in quarkonia decay studied in [12]. Another effect, not studied there, is to produce a non-vanishing $H l_i l_j$ LFV coupling at one level, whose effects in quarkonium decay were not analyzed in [12].

In this paper, we assume that there are no specific models for new physics behind the $H\tau\mu$ coupling. We systematically study the LFV decays of quarkonia involving the $H\tau\mu$ coupling and no other LFV coupling, considering the phenomenological value of its 1.43×10^{-3} pointed by the CMS and ATLAS Collaborations. We shall show that even if the expected branching ratios are still below the experimental limit, some of them could be accessible to the next generation of B factories as Belle II.

2. GENERAL QUARKONIUM FORMALISM

The Higgs decay into heavy quarkonia has been extensively studied, for instance, [14, 15, 12, 16, 17, 18, 19] and references therein. The description of processes involving the annihilation or creation of heavy quarkonia can be systematically done in the framework of nonrelativistic quantum chromodynamics (NRQCD) [20, 21]. This is a systematic expansion in terms of α_s and the quarks relative velocity v with a clear separation of the perturbative phenomena occurring at the scale m_Q and the nonperturbative ones occurring at the scale $m_Q v$. The nonperturbative effects are encoded in universal matrix elements with a well-defined hierarchy in the v expansion. The novelty of this systematic approach is that, for some processes, color-octet configurations of the created or annihilated quark-antiquark pair yield contributions of the same order as the old color-singlet contributions to a given order in the α_s and v expansion.

In this work, we are interested in the order of magnitude of the branching ratios of the considered processes and will focus on the color-singlet contributions which can be calculated using the old quarkonium techniques described in [22, 23]. A more refined analysis can be done in the most promising channels but this is beyond the scope of the present work.

The invariant amplitude for the annihilation of color-singlet quarkonium in a $^{2S+1}L_J$ angular momentum config-

¹We will concentrate on the stronger limit coming from CMS.

	Experimental bound [13]	Expected from $H \rightarrow \mu\tau$
$\tau \rightarrow \mu\gamma$	4.4×10^{-8}	1.3×10^{-9}
$\tau \rightarrow 3\mu$	2.1×10^{-8}	1×10^{-10}

TABLE 1: LFV in τ LFV decays involving charged leptons.

uration $\bar{Q}Q[{}^{2S+1}L_J] \rightarrow X$ is given by [22, 23]

$$\mathcal{M}[\bar{Q}Q[{}^{2S+1}L_J] \rightarrow X] = \int \frac{d^4q}{(2\pi)^4} \text{Tr}[\mathcal{O}(Q, q)\chi(Q, q)], \quad (3)$$

where $\mathcal{O}(Q, q)$ is the operator entering amplitude for the corresponding free quarks transition:

$$\begin{aligned} \mathcal{M}[\bar{Q}(\frac{Q}{2} - q, s_2), Q(\frac{Q}{2} + q, s_1) \rightarrow X] \\ = \bar{v}(\frac{Q}{2} - q, s_2)\mathcal{O}(Q, q)u(\frac{Q}{2} + q, s_1), \end{aligned} \quad (4)$$

and $\chi(Q, q)$ denotes the wave function for the $\bar{Q}Q[{}^{2S+1}L_J]$ bound state:

$$\begin{aligned} \chi(Q, q) = \sum_{M, S_z} 2\pi\delta(q^0 - \frac{\mathbf{q}^2}{2m_Q})\psi_{LM}(\mathbf{q}) \\ \cdot P_{S, S_z}(Q, q)\langle LM; SS_z | JJ_z \rangle. \end{aligned} \quad (5)$$

Here, P_{S, S_z} stands for the spin projectors:

$$\begin{aligned} P_{S, S_z}(Q, q) = \sqrt{\frac{N_c}{m_Q}} \sum_{s_1, s_2} u(\frac{Q}{2} + q, s_1) \\ \cdot \bar{v}(\frac{Q}{2} - q, s_2)\langle \frac{1}{2}s_1; \frac{1}{2}s_2 | SS_z \rangle \\ = \sqrt{\frac{N_c}{m_Q}} \left(\frac{1}{2\sqrt{2}m_Q} \right) \left(\frac{Q}{2} + \not{q} + m_Q \right) \\ \cdot \left\{ \begin{array}{l} \gamma^5 \\ \varepsilon(Q, S_z) \end{array} \right\} \left(\frac{Q}{2} + \not{q} - m_Q \right) \\ \text{for } \left\{ \begin{array}{l} S = 0 \\ S = 1 \end{array} \right\}, \end{aligned} \quad (6)$$

where $\varepsilon(Q, S_z)$ denotes the polarization vector of the spin one system.

For s -wave quarkonium, the wave function is rapidly damped in the relative momentum q and the leading terms are given by $P_{S, S_z}(Q, 0)$ and $\mathcal{O}(Q, 0)$. In the zero-binding approximation, the quarkonium mass M is given by $M \approx 2m_Q$ and the amplitude reads

$$\begin{aligned} \mathcal{M}[\bar{Q}Q[{}^{2S+1}S_J] \rightarrow X] \\ = \frac{R(0)}{\sqrt{4\pi}} \sqrt{\frac{3}{4M}} \text{Tr} \left[\mathcal{O}(Q, 0) \left\{ \begin{array}{l} \gamma^5 \\ \not{\varepsilon}(Q, S_z) \end{array} \right\} (Q - M) \right] \\ \text{for } \left\{ \begin{array}{l} S = 0 \\ S = 1 \end{array} \right\}, \end{aligned} \quad (7)$$

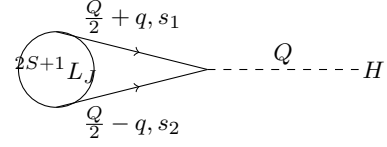


FIGURE 1: Feynman diagrams for the Higgs-quarkonium coupling.

with M denoting the quarkonium physical mass and

$$\int \frac{d^3q}{(2\pi)^3} \psi_{00}(\mathbf{q}) = \frac{R(0)}{\sqrt{4\pi}}. \quad (8)$$

A similar calculation of the invariant amplitude for the production of color-singlet quarkonium, $X \rightarrow \bar{Q}Q[{}^{2S+1}S_J] + Y$ yields

$$\begin{aligned} \mathcal{M}(X \rightarrow \bar{Q}Q[{}^{2S+1}S_J] + Y) \\ = -\frac{R(0)}{\sqrt{4\pi}} \sqrt{\frac{3}{4M}} \text{Tr} \left[\mathcal{O}(Q, 0) \left\{ \begin{array}{l} \gamma^5 \\ \not{\varepsilon}(Q, S_z) \end{array} \right\} (Q + M) \right] \\ \text{for } \left\{ \begin{array}{l} S = 0 \\ S = 1 \end{array} \right\}. \end{aligned} \quad (9)$$

For p -wave quarkonium, the wave function at the origin vanishes and the leading term for the annihilation amplitude is given by the first term in the expansion in q of Equation (3). A straightforward calculation yields

$$\begin{aligned} \mathcal{M}[\bar{Q}Q[{}^{2S+1}P_J] \rightarrow X] \\ = -i \sum_{M, S_z} \langle 1M; SS_z | JJ_z \rangle \varepsilon_\alpha(M) \sqrt{\frac{3}{4\pi}} R'(0) \\ \cdot \text{Tr} \left[\mathcal{O}^\alpha(Q, 0) P_{S, S_z}(Q, 0) + \mathcal{O}(Q, 0) P_{S, S_z}^\alpha(Q, 0) \right], \end{aligned} \quad (10)$$

where

$$A^\alpha(Q, q) \equiv \frac{\partial A(Q, q)}{\partial q_\alpha}, \quad (11)$$

and in this case,

$$\int \frac{d^3q}{(2\pi)^3} q^\alpha \psi_{1M}(\mathbf{q}) = -i \sqrt{\frac{3}{4\pi}} R'(0) \varepsilon_\alpha(M). \quad (12)$$

The polarization vector $\varepsilon_\alpha(M)$ satisfies the following relations:

$$\sum_{M, S_z} \langle 1M; 1S_z | 00 \rangle \varepsilon_\alpha(M) \varepsilon_\beta(S_z) = -g_{\alpha\beta} + \frac{Q_\alpha Q_\beta}{M^2}, \quad (13)$$

$$\sum_{M, S_z} \langle 1M; 1S_z | 1J_z \rangle \varepsilon_\alpha(M) \varepsilon_\beta(S_z) = \frac{-i}{M} \frac{1}{\sqrt{2}} \varepsilon_{\alpha\beta\mu\nu} Q^\mu \varepsilon^\nu(J_z), \quad (14)$$

$$\sum_{M, S_z} \langle 1M; 1S_z | 2J_z \rangle \varepsilon_\alpha(M) \varepsilon_\beta(S_z) = \varepsilon_{\alpha\beta}(J_z). \quad (15)$$

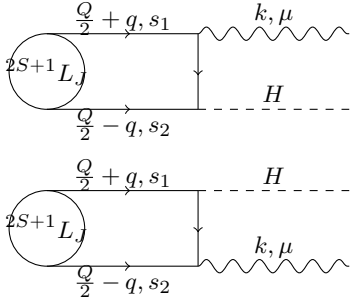


FIGURE 2: Feynman diagrams for the Higgs-quarkonium-photon coupling.

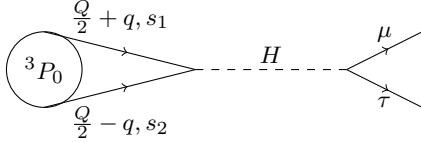


FIGURE 3: Diagram for the $\bar{Q}Q[{}^3P_0] \rightarrow \mu\tau$ decay.

3. HIGGS-QUARKONIUM COUPLING

The Higgs to $\bar{Q}Q[{}^{2S+1}L_J]$ quarkonium coupling is obtained from the diagram in Figure (1) which yields the following operator:

$$\mathcal{O}(Q, q) = i \frac{m_Q}{v}, \quad (16)$$

where v stands for the Higgs vacuum expectation value. Using this operator in the previous formulae, it is easy to show that the only nonvanishing coupling of the Higgs to quarkonium is to $S = 1$ and $J = 0$ p -wave quarkonium, in which case, we obtain

$$\mathcal{M}[\bar{Q}Q[{}^3P_0] \rightarrow H] = \frac{3R'(0)}{v} \sqrt{\frac{3M}{\pi}}. \quad (17)$$

Notice that this coupling is proportional to the derivative of the wave function at the origin, which, according to the NRQCD rules, is suppressed by a v^2 factor with respect to the wave function at the origin. This makes the radiative transitions involving s -wave quarkonium configurations of the same order as the nonradiative ones involving p -wave quarkonium configurations. The radiation changes the quarkonium quantum numbers allowing the corresponding quarkonium to couple to the Higgs. The calculation of Higgs-mediated lepton flavor violating radiative transitions involving s -wave quarkonium requires to work out Higgs-quarkonium-photon coupling. This transition is induced by the diagrams in Figure (2).

From these diagrams, we identify the transition operator as

$$\begin{aligned} \mathcal{O}(Q, q) &= iee_Q \frac{m_Q}{v} \\ &\cdot \left[\not{\epsilon}(k) \frac{\frac{Q}{2} + \not{q} + \not{k} + m_Q}{(\frac{Q}{2} + q + k)^2 - m_Q^2} - \frac{\frac{Q}{2} - \not{q} + \not{k} - m_Q}{(\frac{Q}{2} - q + k)^2 - m_Q^2} \not{\epsilon}(k) \right], \end{aligned} \quad (18)$$

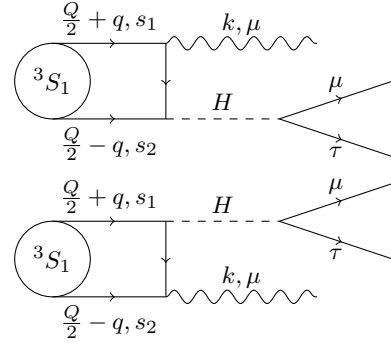


FIGURE 4: Feynman diagrams for the ${}^3S_1 \rightarrow \mu\tau$ decay.

where e_Q stands for the heavy quark charge in the units of e . For s -wave $J = 0$ from Equation (7), we obtain

$$\begin{aligned} \mathcal{M}[H \rightarrow \bar{Q}Q[{}^1S_0]\gamma] &= -i \frac{ee_Q MR(0)}{4v} \sqrt{\frac{3}{4\pi M}} \\ &\cdot \text{Tr} \left[\left(\frac{\not{\epsilon}(k)(\frac{Q}{2} + \not{k}) - (\frac{Q}{2} + \not{k})\not{\epsilon}(k)}{(\frac{Q}{2} + k)^2 - m_Q^2} \right) \gamma^5 (\not{Q} - M) \right] = 0. \end{aligned} \quad (19)$$

Similarly, for s -wave $J = 1$, we get

$$\begin{aligned} \mathcal{M}[H \rightarrow \bar{Q}Q[{}^3S_1]\gamma] &= i \frac{ee_Q m_Q R(0)}{2v} \sqrt{\frac{3}{4\pi M}} \\ &\cdot \text{Tr} \left[\left(\frac{\not{\epsilon}(k)(\frac{Q}{2} + \not{k}) - (\frac{Q}{2} + \not{k})\not{\epsilon}(k)}{(\frac{Q}{2} + k)^2 - m_Q^2} \right) \not{\eta}(Q) (\not{Q} - M) \right], \end{aligned} \quad (20)$$

where $\eta(Q)$ stands for the polarization vector of the quarkonium. A straightforward calculation yields

$$\mathcal{M}[H \rightarrow \bar{Q}Q[{}^3S_1]\gamma] = \frac{ee_Q R(0)}{v} \sqrt{\frac{3M}{\pi}} T_{\mu\nu} \epsilon^\mu \eta^\nu \quad (21)$$

with

$$T_{\mu\nu} = g_{\mu\nu} - \frac{Q^\mu k^\nu}{Q \cdot k}. \quad (22)$$

4. $\bar{Q}Q[{}^3P_0] \rightarrow \mu\tau$

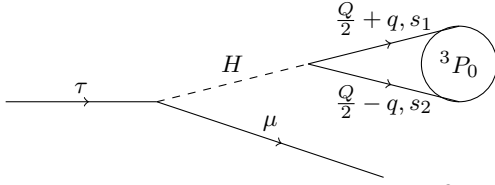
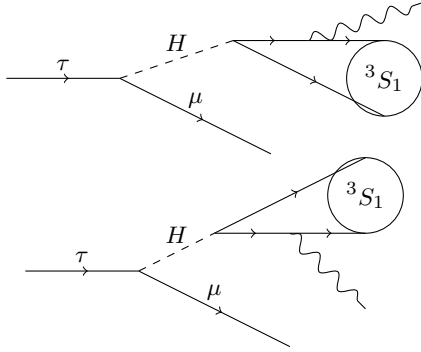
Now, we focus on the Higgs-mediated LFV processes. We start with the $\bar{Q}Q[{}^3P_0] \rightarrow \mu\tau$ decay through the diagram in Figure (3). A direct calculation yields the following decay width:

$$\Gamma[\bar{Q}Q[{}^3P_0] \rightarrow \mu^- \tau^+] = \frac{27y^2 |R'(0)|^2 M^2}{8\pi^2 v^2 m_H^4} \left(1 - \frac{m_\tau^2}{M^2} \right), \quad (23)$$

where we neglected the muon mass and y stands for the $H\mu\tau$ coupling.

5. $\bar{Q}Q[{}^3S_1] \rightarrow \mu\tau\gamma$

Next, we go through the corresponding radiative process $\bar{Q}Q[{}^3S_1] \rightarrow \mu\tau\gamma$. This decay proceeds through the diagrams in Figure (4).


FIGURE 5: Feynman diagram for the $\tau \rightarrow \mu \bar{Q}Q[{}^3P_0]$ decay.

FIGURE 6: Feynman diagram for the $\tau \rightarrow \mu \bar{Q}Q[{}^3S_1]\gamma$ decay.

Neglecting the muon mass, we obtain the following decay width:

$$\Gamma[\bar{Q}Q[{}^3S_1] \rightarrow \mu^- \tau^+ \gamma] = \frac{\alpha e_Q^2 y^2 |R(0)|^2 M^4}{12\pi^3 v^2} \frac{M^4}{m_H^4} f\left(\frac{m_\tau^2}{M^2}\right), \quad (24)$$

where

$$f(x) = 1 - 6x + 3x^2 + 2x^3 - 6 \ln(x). \quad (25)$$

6. $\tau \rightarrow \mu \bar{Q}Q[{}^3P_0]$

The $H\mu\tau$ coupling can also mediate LFV decays of the tau meson involving light quarkonium. Although this is beyond the scope of the systematic NRQCD expansion due to the light quark mass, we still can use the quarkonium techniques taking care of extracting the corresponding nonperturbative pieces from the appropriate experimental data. The first possible decay is $\tau \rightarrow \mu \bar{Q}Q[{}^3P_0]$ which goes through the diagram shown in Figure (5). The decay width is given by

$$\Gamma(\tau^- \rightarrow \mu^- \bar{Q}Q[{}^3P_0]) = \frac{27y^2 |R'(0)|^2 m_\tau M}{16\pi^2 v^2} \frac{m_\tau M}{m_H^4} \left(1 - \frac{M^2}{m_\tau^2}\right), \quad (26)$$

where we neglected the muon mass.

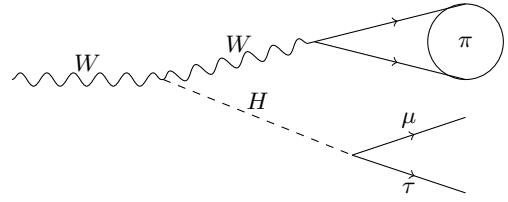
7. $\tau \rightarrow \mu \bar{Q}Q[{}^3S_1]\gamma$

The corresponding radiative decay is $\tau \rightarrow \mu \bar{Q}Q[{}^3S_1]\gamma$. The Feynman diagrams for this process are given in Figure (6). Neglecting the muon mass, we obtain the following decay width:

$$\Gamma(\tau^- \rightarrow \mu^- \bar{Q}Q[{}^3S_1]\gamma) = \frac{\alpha e_Q^2 y^2 |R(0)|^2 m_\tau^3 M}{32\pi^3 v^2} \frac{m_\tau^3 M}{m_H^4} h\left(\frac{M^2}{m_\tau^2}\right), \quad (27)$$

where

$$h(x) = (1-x)^3 + \frac{3}{2}x[(1-x)(3-x) + 2\ln(x)]. \quad (28)$$


FIGURE 7: Feynman diagram for the $W \rightarrow \tau\mu\pi$ decay.

Finally, although the calculation does not require to use the quarkonium techniques, it is interesting to estimate the effects of the $H\mu\tau$ coupling in LFV decay of gauge bosons. As a sample, we calculate $W \rightarrow \mu\tau\pi$. The nonperturbative piece of this decay is related to the pion decay constant. This decay is induced by the diagram in Figure(7).

The amplitude for $W(Q, \varepsilon) \rightarrow \mu(p_1)\tau(p_2)\pi(p_3)$ is

$$\mathcal{M} = \frac{yg^2 V_{ud} f_\pi}{2\sqrt{2}m_W} \frac{p_3 \cdot \varepsilon(Q)}{(p_1 + p_2)^2 - m_H^2} \bar{u}(p_2)v(p_1), \quad (29)$$

where g is the weak coupling constant. The resulting decay width is

$$\Gamma(W^- \rightarrow \mu^- \tau^+ \pi^-) = \frac{f_\pi^2 g^2 \lambda^2 V_{ud}^2}{73728\pi^3 m_W} f(a, b), \quad (30)$$

where

$$f(a, b) = \frac{1}{a^4} \left[a^2 (b^2 - 1) (24a^6 - 6a^4 (4b^2 + 7) + a^2 (b^2 + 17) (2b^2 + 1) - 6b^2) - 6 (a^2 - 1)^2 (4a^6 - a^4 (6b^2 + 1) + 2a^2 b^4 + b^4) \cdot \ln\left(\frac{a^2 - 1}{a^2 - b^2}\right) - 12b^4 \ln(b) \right], \quad (31)$$

where $a = \frac{m_H}{m_W}$, $b = \frac{m_\tau}{m_W}$, and we neglected the pion and muon masses.

8. NUMERICAL RESULTS

The results for the studied decays depend on the color-singlet matrix elements $R(0)$ for 3S_1 quarkonium and $R'(0)$ for 3P_0 quarkonium. The same matrix elements appear in the leptonic decay of the first- and two-photon decays of the latter. We use the available experimental results on these decays to extract the phenomenological value of the matrix elements. The only matrix element that cannot be calculated this way is $R'(0)$ for the ${}^3P_0 \bar{b}b$ quarkonium χ_{b0} . There is no available experimental data on the $\chi_{b0} \rightarrow \gamma\gamma$ transition, but $R'(0)$ has been calculated in several potential models summarized in [24] yielding $R'(0) \approx 1 \text{ GeV}^5$, and we will use this value in our calculations.

The leptonic decay of vector quarkonia is induced by the diagram in Figure (8).

The corresponding decay width is

$$\Gamma(\bar{Q}Q[{}^3S_1] \rightarrow l^+ l^-) = \frac{4\alpha^2 e_Q^2 |R(0)|^2}{M^2}, \quad (32)$$

where we neglected the lepton mass. The two-photon decay of 3P_0 quarkonium proceeds through the diagrams in Figure (9).

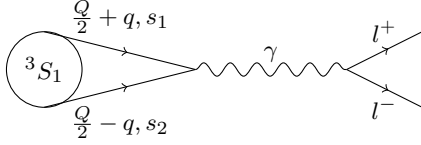


FIGURE 8: Diagram for the $\bar{Q}Q[{}^3S_1] \rightarrow l^+l^-$ decay.

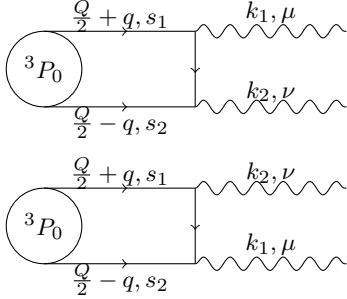


FIGURE 9: Feynman diagrams for the two-photon decay of 3P_0 quarkonium.

The decay width is

$$\Gamma({}^3P_0 \rightarrow \gamma\gamma) = \frac{432\alpha^2 e_Q^2 |R'(0)|^2}{M^4}. \quad (33)$$

In Table 2, we give the results for the matrix elements of the different quarkonia.

There is no available information on the total width of the χ_{b0} ; thus, we report the decay width when we use $R'_{\chi_b}(0) = 1 \text{ GeV}^5$ from quark model calculations [24]:

$$\Gamma(\chi_{b0} \rightarrow \mu\tau) = 5.5 \times 10^{-17} \text{ GeV}. \quad (34)$$

The branching ratios of the remaining decays are calculated using the estimates for the nonperturbative matrix elements in Table 2. We list in Table 3 the so-obtained branching ratios. Here, the branching ratios include a factor 2 to account for the two charge states where appropriate; e.g., $BR(\chi_{c0} \rightarrow \mu\tau) = BR(\chi_{c0} \rightarrow \mu^- \tau^+) + BR(\chi_{c0} \rightarrow \mu^+ \tau^-)$.

As said previously, the quarkonium formalism described above cannot be directly applied to $f_0(980)$ as it is composed of light quarks, even assuming that it is an $\bar{s}s$ [25, 26, 27, 28, 29, 30, 31, 32]. A way to describe it can be done using the methods given in [33, 34, 35, 36, 37] and references therein. For $\tau \rightarrow \mu f_0$, the form of Equation (26) is still valid. So using the $f_0(980)$ form factors as given in [38, 39, 40], one obtains that

$$\tau \rightarrow \mu f_0(980) \approx 10^{-11}$$

compared to the experimental limit ($< 3.4 \times 10^{-8}$ [13]). This value is very close to the value obtained using the quarkonium formalism to $f_0(980)$ extracting the effective form factor $|R'(0)|^2 (\text{GeV}^5)$ from the $f_0 \rightarrow \gamma\gamma$ branching ratio (0.29×10^{-6} [13]):

$$|R'(0)|^2 (\text{GeV}^5) = 1.08 \times 10^{-4}.$$

Compared to other channels, it seems that the $\tau \rightarrow \mu f_0(980)$ is the most promising channel to be observed in the next experiment generations as BELLE II.

Process	Γ_{exp} (GeV)	$ R(0) ^2$ (GeV ³)	$ R'(0) ^2$ (GeV ⁵)
$Y \rightarrow e^+e^-$	1.28×10^{-6}	4.856	-
$J/\psi \rightarrow e^+e^-$	5.54×10^{-6}	0.560	-
$\phi \rightarrow e^+e^-$	1.26×10^{-6}	5.53×10^{-2}	-
$\chi_c^0 \rightarrow \gamma\gamma$	2.34×10^{-6}	-	3.10×10^{-2}

TABLE 2: Numerical values of the nonperturbative matrix elements extracted from the leptonic and two-photon decays of quarkonia.

Process	Branching Ratio	Exp. bound
$\chi_{c0} \rightarrow \mu\tau$	1.5×10^{-17}	
$Y \rightarrow \mu\tau\gamma$	5.7×10^{-14}	
$J/\psi \rightarrow \mu\tau\gamma$	5.1×10^{-17}	
$\tau \rightarrow \mu\phi\gamma$	1.7×10^{-14}	
$W \rightarrow \mu\tau\pi$	3.2×10^{-17}	

TABLE 3: Branching ratios for lepton flavor violation decays involving the $H\mu\tau$ coupling.

9. CONCLUSION

A model-independent study of LFV processes involving quarkonium has been performed in light of the possibility of having a nonvanishing $H\tau\mu$ coupling smaller than 0.25% as pointed by the CMS Collaboration. The expressions for the LFV quarkonium decays (including the LFV radiative decays) have been presented. The numerical values for the branching ratio are usually small. In conclusion, it could be difficult to observe the LFV in quarkonium decay in the present generation of collider except for the $f_0(980)$ case which is not far from the limit that BELLE II could reach in the near future. A better experimental knowledge of the different quarkonium states is also needed to improve our determination of the involved form factors and of the underlying dynamics.

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