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Study of the $\eta_b \to \tau^+\tau^-$ decay as a probe for light pseudoscalar, axial vector states

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Abstract

At BaBar experiment the branching ratio for $\eta_b \to \tau^+\tau^-$ was constrained at $\mathcal{BR}(\eta_b \to \tau^+\tau^-) < 8\%$ at 90% confidence level (C.L.). In this article we explore the decay $\eta_b \to \tau^+\tau^-$ as a probe for a light pseudoscalar or a light axial vector state. We find the standard model branching ratio for this decay to be $\sim 4 \times 10^{-9}$. We show that considerably larger branching ratios, up to the present experimental constraint, is possible in models with a light pseudoscalar or a light axial vector state.

At the BaBar experiment the pseudoscalar $b\bar{b}$ bound state (η_b) in the 1S configuration was observed. It was seen by two research groups in BaBar in two different experiments. First, it was seen in the decay of $\Upsilon(3S) \to \gamma\eta_b$ [1] with a signal significance greater than 10 standard deviations (σ). The η_b was seen in the photon energy spectrum using (109 ± 1) million $\Upsilon(3S)$ events and the hyperfine $\Upsilon(1S) - \eta_b$ mass splitting was measured to be $71.4^{+2.3}_{-3.1}(\text{stat}) \pm 2.7(\text{syst})$ MeV from the mass $m(\eta_b) = 9388.9^{+3.1}_{-2.3}\,(\text{stat}) \pm 2.7\,(\text{syst})$ MeV. Soon after, it was also observed in $\Upsilon(2S) \to \gamma\eta_b$ [2] by another group in BaBar, and the hyperfine mass splitting was determined to be $67.4^{+4.8}_{-4.6}(\text{stat}) \pm 2.0(\text{syst})$ MeV from the mass $m(\eta_b) = 9392.9^{+4.6}_{-4.8}\,(\text{stat}) \pm 1.9\,(\text{syst})$ MeV. Many experimental environments [3, 4, 5] have been running to search for the ground state η_b but without success. On the other side, various attempts within lattice NRQCD [6, 7] and perturbative QCD based models [7, 8] have been made to predict the mass of η_b. The lattice NRQCD approach was more accurate with the hyperfine splitting prediction of $E^{lat}_{hfs} = 61 \pm 14$ MeV and correspondingly mass $m_{\eta_b} = 9383(4)(2)$ MeV which are in agreement with the experimental results. An experiment at BaBar has searched for a low-mass Higgs boson in $\Upsilon(3S) \to \gamma A^0$, $A^0 \to \tau^+\tau^-$ [9] with data sample containing 122 million $\Upsilon(3S)$ events. In the same paper, the branching ratio for $\eta_b \to \tau^+\tau^-$ was constrained at $\mathcal{BR}(\eta_b \to \tau^+\tau^-) < 8\%$ at 90% confidence level (C.L.).

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In this work [10], we probe light scalar and spin 1 states via η_b decay. We study the process $\eta_b \to \tau^+\tau^-$ mediated by a pseudoscalar (A^0) or an axial vector (U). This decay process in the SM can only be mediated at the tree level through a Z exchange with very small branching ratio $\sim 4 \times 10^{-9}$, or a higher order contribution via two intermediate photons with a tiny branching ratio $\sim 10^{-10}$. Therefore, a measurement of $BR[\eta_b \to \tau^+\tau^-]$ larger than the SM rate would be a signal of new states. Probing $A^0(U)$ in the process $\eta_b \to \tau^+\tau^-$ has several advantages over Υ decays when $A^0(U)$ is off-shell. First, the new physics (NP) branching ratio of $\Upsilon \to \tau^+\tau^-\gamma$ is suppressed by the second order coupling to $A^0(U)$ in conjunction with another state, usually a photon. Secondly, the SM branching ratio of $\Upsilon \to \tau^+\tau^-\gamma$ can be enhanced through a radiative decay, which we estimate it to be $= 4.4 \times 10^{-3}$ with $E_\gamma > 100$ MeV.

In the SM, we calculate the branching ratio for $\eta_b \to \tau^+\tau^-$ via the Z exchange to be given by

$$\Gamma^Z(\eta_b \to \tau^+\tau^-) \quad = \quad \frac{G_F^2 M_W^4 m_\tau^2 f_{\eta_b}^2 m_{\eta_b}}{16\pi \cos^4\theta_W} \beta_\tau \left(1 - \frac{m_{\eta_b}^2}{M_Z^2}\right)^2 |a_Z|^2, \tag{1}$$

where θ_W denotes the Weinberg angle, $\beta_\tau = \sqrt{1 - \left(\frac{2m_\tau}{m_{\eta_b}}\right)^2}$ is the velocity of the τ lepton in the η_b rest frame and $|a_Z|^2 \equiv 1/((m_{\eta_b}^2 - M_Z^2)^2 + M_Z^2\Gamma_Z^2)$. The decay constant f_{η_b} in Eq. 1 is defined as [11], $\langle 0|\, \bar{b}(0)\gamma_\mu\gamma_5 b(0)\, |\eta_b(q)\rangle = i f_{\eta_b} q_\mu$. The process $\eta_b \to \tau^+\tau^-$ can also go via two

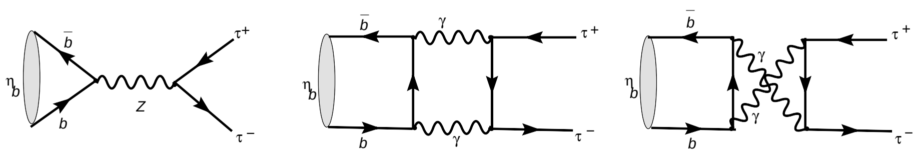

Figure 1: Various processes contributing to $\eta_b \to \tau^+\tau^-$ in the SM.

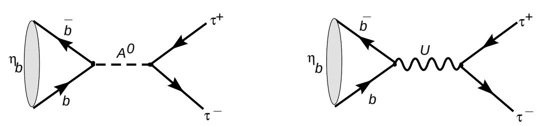

Figure 2: Various processes contributing to $\eta_b \to \tau^+\tau^-$ in NP.

photon intermediate states as shown in Fig. 1. This diagram is dominated by the imaginary part [12] which we can calculate using unitarity [13] to obtain

$$\Gamma^{2\gamma}[\eta_b \to \tau^+\tau^-] \quad \geq \quad \frac{\alpha^2}{2\beta_\tau} \left[\frac{m_\tau}{m_{\eta_b}} \ln \frac{(1+\beta_\tau)}{(1-\beta_\tau)}\right]^2 \Gamma[\eta_b \to \gamma\gamma], \tag{2}$$

where α is the electromagnetic fine structure constant. We find that $\Gamma[\eta_b \to \gamma\gamma] = (\pi\alpha^2 m_{\eta_b} f_{\eta_b}^2/81 m_b^2)$, where we use the heavy quark limit for the b quark. Numerically, one find that $\Gamma^{2\gamma}[\eta_b \to \tau^+\tau^-] \gtrsim 4 \times 10^{-10}$. The 2γ exchange contribution is mostly imaginary relative to the Z

exchange contribution. Therefore, the total width $\Gamma_t[\eta_b \to \tau^+\tau^-]$ can be given to a good approximation as

$$\Gamma_t[\eta_b \to \tau^+\tau^-] \approx \Gamma^Z[\eta_b \to \tau^+\tau^-] + \Gamma^{2\gamma}[\eta_b \to \tau^+\tau^-]. \tag{3}$$

In the NP model, we consider the decay process $\eta_b \to \tau^+\tau^-$ mediated via the pseudoscalar Higgs boson A^0 in the 2HDM as in Fig. 2(a). The couplings of the down-type quarks D and charged leptons ℓ with A^0 in the generic 2HDM model are given by [14]

$$\mathcal{L}_{A^0}^{D,\ell} = \frac{igF_{A^0}}{2M_W}(\bar{D}M_D^{diag}\gamma_5 D + \bar{\ell}M_\ell^{diag}\gamma_5\ell)A^0, \tag{4}$$

where F_{A^0} is a model-dependent parameter, $M_D^{diag} = (m_d, m_c, m_b)$ and $M_\ell^{diag} = (m_e, m_\mu, m_\tau)$ are the diagonal mass matrices of D and ℓ, respectively. We consider $F_{A^0} > 1$ in our analysis. In the case of 2HDM type (II) $F_{A^0} \equiv tan\beta$ while in 2HDM type (I) $F_{A^0} \equiv -\cot\beta$ where $tan\beta = v_2/v_1$ is the ratio between the VEVs of the two Higgs doublets. The decay rate for this process was obtained as,

$$\Gamma^{A^0}(\eta_b \to \tau^+\tau^-) = \frac{G_F^2 m_\tau^2 f_{\eta_b}^2 m_{\eta_b}^5}{16\pi}\beta_\tau |a_{A^0}|^2, \tag{5}$$

where the coefficient a_{A^0} depends on the mass m_{A^0} as, $|a_{A^0}|^2 \equiv (F_{A^0}^4/(m_{\eta_b}^2 - m_{A^0}^2)^2)$. We assume that the decay width Γ^{A^0} for the A^0 is negligible. In Eq. 5, we use $\langle 0|\,\bar{b}(0)\gamma_5 b(0)\,|\eta_b(q)\rangle = (if_{\eta_b} m_{\eta_b}^2/2m_b)$.

Some specific models contain a light axial vector state, U. We introduce the U state as a mediator for the process $\eta_b \to \tau^+\tau^-$. We take the U couplings to the down-type quarks and charged leptons to be given by

$$\mathcal{L}_U^{D,\ell} = f_A^{D,\ell}(\bar{D}\gamma^\mu\gamma_5 D + \bar{\ell}\gamma^\mu\gamma_5\ell)U_\mu, \tag{6}$$

with the axial coupling $f_A^{D,\ell} = 2^{-\frac{3}{4}} G_F^{\frac{1}{2}} m_U F_U$, where m_U denotes the mass of U-boson and F_U denotes a model-dependent parameter. In the specific model [15, 16, 17], $F_U \equiv \cos\zeta\tan\beta$. We consider $F_U > 1$ in our analysis. The decay rate for $\eta_b \to \tau^+\tau^-$ can be obtained as

$$\Gamma^U(\eta_b \to \tau^+\tau^-) = \frac{G_F^2 m_\tau^2 f_{\eta_b}^2 m_{\eta_b}}{16\pi}\beta_\tau F_U^4. \tag{7}$$

The above equation shows an interesting result; if the width of the U is neglected the $\Gamma^U(\eta_b \to \tau^+\tau^-)$ will be independent of the mass of U-boson. This result follows from the fact that the axial U-boson coupling to fermions is proportional to the mass m_U and the fact that η_b is a pseudoscalar. The result of Eq. 7 is not reliable as m_U gets considerably large and the coupling becomes non-perturbative. One gets the constraints $m_U \leq \frac{4M_W}{gF_U}$ requiring the coupling to be ≤ 1. Hence, for $F_U \sim 50$ one find m_U to be in the GeV range.

In Fig. 3, we plot the logarithm of the branching ratio for $\eta_b \to \tau^+\tau^-$ mediated by the pseudoscalar A^0 in a generic 2HDM model. We use various values of the A^0 mass between 0.1 to 20 GeV to plot the branching ratio $\mathcal{BR}^{A^0}$. As the mass of the A^0 approaches the mass of the η_b the branching ratio increases and blows up at $m_{A^0} = m_{\eta_b}$. We observe that the

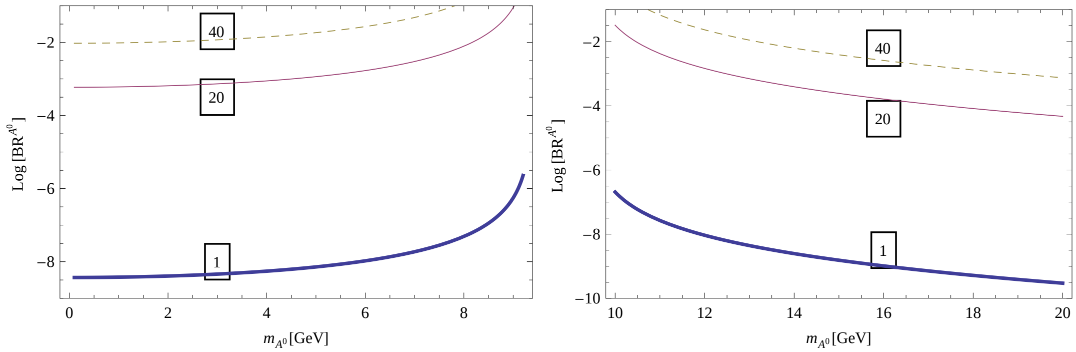

Figure 3: The logarithm of $\mathcal{BR}^{A^0}(\eta_b \to \tau^+\tau^-)$ as a function of m_{A^0} for different values of F_{A^0} and $m_{A^0} \in [0.1, 20]$ GeV.

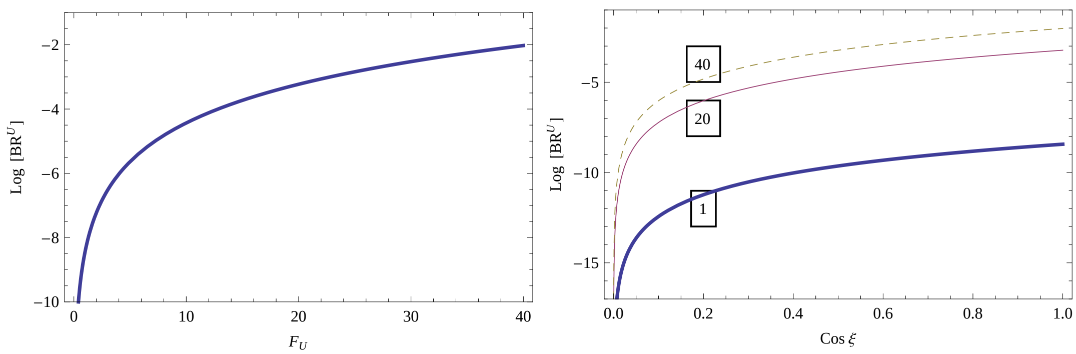

Figure 4: On the left panel: The logarithm of $\mathcal{BR}^{U}(\eta_b \to \tau^+\tau^-)$ as a function of F_U. On the right panel: The logarithm of $\mathcal{BR}^{U}(\eta_b \to \tau^+\tau^-)$ as a function of $\cos\zeta$ for different values of $\tan\beta$ and $\cos\zeta \in [0, 1]$.

branching ratio $\sim F_{A^0}^4$ is more sensitive to F_{A^0} than to the mass m_{A^0}. We note from the plots in Fig. 3 that $\mathcal{BR}^{A^0}$ can be significantly larger than the SM branching ratios and can vary from $\sim 10^{-8}$ to the experimental bound of 8 % for $F_{A^0} = 40$. We note from Fig. 3 that outside the range of the η_b mass, the branching ratio for $\eta_b \to \tau^+\tau^-$ can be significant and we expect the same to be true also in the mass range where mixing effects are important.

In Fig. 4, we plot the logarithm of the branching ratio for $\eta_b \to \tau^+\tau^-$ versus F_U. Working in a specific model [15, 16, 17] $F_U \equiv \cos\zeta \tan\beta$. We plot the branching ratio versus the invisibility factor $\cos\zeta$ for different values of $\tan\beta$. Again we observe in a model with axial vector boson that the branching ratio can vary over a wide range and can be much larger than the SM prediction.

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