Probing light pseudoscalar, axial vector states through $\eta_b \to \tau^+\tau^-$

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In this paper, we explore the decay $\eta_b \to \tau^+\tau^-$ as a probe for a light pseudoscalar or a light axial vector state. We estimate 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 limit of $\sim 8\%$, are possible in models with a light pseudoscalar or a light axial vector state. As we do not include possible mixing effects between the light pseudoscalar and the η_b, our results should be reliable when the pseudoscalar mass is away from the η_b mass.

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I. INTRODUCTION

It is widely anticipated that physics beyond the standard model (SM) or new physics (NP) will be discovered soon at experiments such as the LHC. This NP might contain new gauge bosons, additional Higgs bosons beyond the SM Higgs, or new quarks and leptons. It is generally believed that these new particles will be heavy with masses from the weak scale $\sim$100 GeV to a TeV. However, light scalars and vector bosons with masses in the GeV range or even lower are not ruled out. For instance, light scalar states coming from a primary Higgs with non-SM decays can be consistent with existing experimental constraints [1]. One of the ways to probe these light states is to look at decays of particles with masses in the 10 GeV range, such as the Υ. Data from the present and future B factories can be used to search for these states and/or to put constraints on models that predict such states.

The pseudoscalar $b\bar{b}$ bound state in the 1S configuration, the η_b, was recently observed. Two research groups in BABAR observed it in two different experiments. First, it was seen in the decay of $\Upsilon(3S) \to \gamma\eta_b$ [2] with a signal significance greater than 10 standard deviations (σ). The η_b was observed 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 $m(\eta_b) = 9388.9^{+3.1}_{-2.3} \cdot (\text{stat}) \pm 2.7(\text{syst})$ MeV. Soon after, it was also seen in $\Upsilon(2S) \to \gamma\eta_b$ [3] 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. In the past, since the discovery of the $\Upsilon(nS)$ resonances [4] in 1977, various experimental environments [5–7] have been used to seek the ground state η_b, but without success. Many theoretical models have attempted to predict the mass of η_b. Lattice NRQCD [8,9] predicts the hyperfine splitting to be $E^{\text{lat}}_{\text{hfs}} = 61 \pm 14$ MeV and, correspondingly, the mass to be $m_{\eta_b} = 9383(4)(2)$ MeV, which is in agreement with the experimental results. The calculations of perturbative QCD-based models [9,10] predict the hyperfine splitting to be $E^{\text{QCD}}_{\text{hfs}} = 39 \pm 11(\text{th})^{+9}_{-8}(\delta\alpha_s)$ MeV, which is smaller than the measured values. Experiments at BABAR have also searched for a low-mass Higgs boson in $\Upsilon(3S) \to \gamma A^0$, $A^0 \to \tau^+\tau^-$ [11] with data sample containing 122×10^6 $\Upsilon(3S)$ events. In the same analysis, constraint on the branching ratio for $\eta_b \to \tau^+\tau^-$ was reported as $\mathcal{BR}(\eta_b \to \tau^+\tau^-) < 8\%$ at 90% confidence level (C.L.).

In this paper, we will be interested in probing light scalar and spin-1 states via η_b decays. As the η_b is a pseudoscalar, a light pseudoscalar and a spin-1 state with axial vector coupling can directly couple to η_b. We will assume the pseudoscalar to couple to the mass of the fermion, as is usually the case for Higgs coupling to fermions. Hence, the η_b, which is a $b\bar{b}$-bound state, has advantages over the η_c and η/η' mesons, which are $c\bar{c}$ and $q\bar{q}(q = u, d, s)$-bound states, respectively. The η_b is expected to be a sensitive probe of a light axial vector state. This follows from the fact that the longitudinal polarization of the axial vector, $\epsilon^\mu_L \sim k^\mu$, when k^μ, the momentum of the vector boson is much larger than its mass. Consequently, the effective axial vector-fermion pair coupling is proportional to the fermion mass for the longitudinal polarization.

In this work, we will study the process $\eta_b \to \tau^+\tau^-$ mediated by a pseudoscalar (A^0) or an axial vector (U). In the SM, this process can only go through a Z exchange at tree level and is highly suppressed, with a branching ratio $\sim 4 \times 10^{-9}$. There is also a higher order contribution to $\eta_b \to \tau^+\tau^-$ in the SM, via two intermediate photons. The branching ratio for this process is also tiny: $\sim 10^{-10}$. Hence, a measurement of $\text{BR}[\eta_b \to \tau^+\tau^-]$ larger than the SM rate will be a signal of new states. One can also probe the states $A^0(U)$ in Υ decays. To search for light $A^0(U)$ states in Υ decays, one generally considers the decay chains, $\Upsilon \to A^0(U)\gamma(A^0(U) \to \tau^+\tau^-)$ [11]. In other words, the $A^0(U)$ is assumed to be produced on-shell. One

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then looks for a peak in the invariant mass of the τ pairs. The experimental measurement/constraint of BR[$\Upsilon \to A^0(U)\gamma$] $\times$ BR[$A^0(U) \to \tau^+\tau^-$] can be converted into a measurement/constraint on the coupling of the $A^0(U)$ to $b\bar{b}$, and hence on model parameters, if the BR[$A^0(U) \to \tau^+\tau^-$] is used as an input [12]. Clearly, as $m_{A^0(U)} > m_\Upsilon$, the $A^0(U)$ can no longer be produced on-shell and the rate for $\Upsilon \to \tau^+\tau^-\gamma$ will fall; consequently, the constraints on the model parameters will be weaker. Note that the constraint $m_{A^0} < 2m_B$ needs to be assumed in the very particular case where the CP-even Higgs mass $m_h < 114$ GeV and $h \to 2A^0$ dominates over $h \to 2m_b$ [1]. In general, $m_{A^0} > 2m_B$ is also possible. We will just assume the existence of light pseudoscalar and axial vector states close to the η_b mass, but they can have masses that are greater than or less than $2m_b$.

The η_b has only been seen in the radiative decays $\Upsilon \to \gamma\eta_b$. Hence, the decay $\eta_b \to \tau^+\tau^-$ has only been studied via the decay $\Upsilon \to \tau^+\tau^-\gamma$. However, the decay $\eta_b \to \tau^+\tau^-$ can be studied independently from the process $\Upsilon \to \tau^+\tau^-\gamma$, as the η_b can be produced from various other processes such as two-photon collisions, $\gamma\gamma \to \eta_b$ [6], and in two-parton collisions [7,13], in hadron colliders like the Tevatron and the LHC. The process $\eta_b \to \tau^+\tau^-$ has several advantages over Υ decays in probing $A^0(U)$ states, especially when $A^0(U)$ is off-shell, which is always the case when $m_{A^0(U)} > m_\Upsilon$. First, unlike the η_b, which can couple directly to $A^0(U)$, the Υ can only couple to $A^0(U)$ in conjunction with another state—usually a photon. Hence, the Υ couplings are second order, and therefore it can decay only to the $\tau^+\tau^-\gamma$ state with a rate much smaller than the rate for $\eta_b \to \tau^+\tau^-$. However, the Υ states are narrower than the η_b, which may partially compensate the larger rate for $\eta_b \to \tau^+\tau^-$ relative to $\Upsilon \to \tau^+\tau^-\gamma$ in the branching ratio measurements. Secondly, an important distinction between $\Upsilon \to \tau^+\tau^-\gamma$ and $\eta_b \to \tau^+\tau^-$ is that the former decay can also proceed as a radiative decay in the SM, while the latter decay is highly suppressed in the SM, as indicated above. Adapting the expression used to estimate the SM branching ratio for $J/\psi \to e^+e^-\gamma$ [14], with the γ emitted from the final-state electrons, to the decay $\Upsilon \to \tau^+\tau^-\gamma$, the rate for this decay in the SM is

$$d\Gamma_{\Upsilon\to\tau^+\tau^-\gamma} = d\Gamma_{\Upsilon\to\tau^+\tau^-}\beta'^3 \frac{2\alpha}{\pi}\frac{dE'_\gamma}{E'_\gamma}\frac{s'}{s} \times \frac{1-\cos^2\theta'_{\gamma\tau}}{(1-\beta'^2\cos^2\theta'_{\gamma\tau})^2}d\Omega'_\gamma, \quad (1)$$

with

$$d\Gamma_{\Upsilon\to\tau^+\tau^-} = \frac{3}{3+\lambda}(1+\lambda\cos^2\theta'_\tau)\Gamma_{\Upsilon\to\tau^+\tau^-}\frac{d\Omega'_\tau}{4\pi}. \quad (2)$$

Here, E'_γ represents the γ energy, θ'_γ and $\phi'_\gamma(\Omega'_\gamma)$ the γ angles, and θ'_τ and $\phi'_\tau(\Omega'_\tau)$ the τ angles, all in the $\tau^+\tau^-$ c.m. frame. β' is the τ velocity and $\theta'_{\gamma\tau}$ is the angle between the τ and γ directions, also in the $\tau^+\tau^-$ c.m. frame, while s' is the $\tau^+\tau^-$ invariant mass squared and s is the Υ invariant mass squared. The parameter λ is determined from the experimental data to be (0.88 ± 0.19) [14]. Using the branching ratio for $\Upsilon \to \tau^+\tau^- = 2.6 \times 10^{-2}$ [15], we estimate the branching ratio for $\Upsilon \to \tau^+\tau^-\gamma = 4.4 \times 10^{-3}$ with $E_\gamma > 100$ MeV.

Naively, the rate for $\Upsilon \to \tau^+\tau^-\gamma$ through an off-shell A^0, from a type (II) Two Higgs Doublet Model (2HDM), relative to the SM rate for $\Upsilon \to \tau^+\tau^-\gamma$, is $\sim \frac{g^4\tan^4\beta m_b^2 m_\tau^2}{16e^4M_W^4}$. Therefore, for large $\tan\beta \sim 28$, the SM and the NP rates may be comparable. However, given the hadronic uncertainties in estimating the SM and the NP rates for $\Upsilon \to \tau^+\tau^-\gamma$, it will be difficult to distinguish between the NP and the SM contributions. Hence, searching for $A^0(U)$ with $m_{A^0(U)} > m_\Upsilon$ in $\Upsilon \to \tau^+\tau^-\gamma$ will be very difficult because of the large SM background. Note that even in e^+e^- machines like the B factories where the η_b is produced through the decay $\Upsilon \to \gamma\eta_b$, the product of branching ratios BR[$\Upsilon \to \gamma\eta_b$] $\times$ [$\eta_b \to \tau^+\tau^-$] is tiny in the SM because of the highly suppressed BR[$\eta_b \to \tau^+\tau^-$] $\sim 4 \times 10^{-9}$. Using the measured BR[$\Upsilon \to \gamma\eta_b$] $\sim 5 \times 10^{-4}$ [2,3], one obtains BR[$\Upsilon \to \gamma\eta_b$] $\times$ [$\eta_b \to \tau^+\tau^-$] $\sim 2 \times 10^{-12}$, which is very difficult to measure. In the presence of NP, this product of branching ratios is enhanced and can reach $\lesssim 10^{-5}$. Hence, the observation of $\Upsilon \to \gamma\tau^+\tau^-$, with the τ pairs coming from η_b, at branching ratios much larger than the SM expectations, will be a signal for new light states. In summary, the large SM background in $\Upsilon \to \tau^+\tau^-\gamma$ and a tiny SM contribution to $\eta_b \to \tau^+\tau^-$ makes the later decay potentially a better probe for $A^0(U)$ than the former if the decays proceed through the off-shell exchange of $A^0(U)$.

There are good theoretical motivations for the existence of a light CP odd A^0 Higgs boson or an axial vector boson U with masses, m_{A^0} and m_U,respectively, in the GeV range or below. There has been interest in the $m_{A^0} < 2m_B$ region, for which a light Higgs, h, with SM-like WW, ZZ, and fermionic couplings can have mass $m_h \sim 100$ GeV while still being consistent with LEP data by virtue of $h \to A^0A^0$. This scenario could even explain the 2.3σ excess in the $e^+e^- \to Z + 2b$ channel for $M_{2b} \sim 100$ GeV [16]. Such a light pseudoscalar Higgs can naturally arise in extensions of minimal supersymmetric model with additional singlet scalars and fermions (gauge-singlet supermultiplets) known as next-to-minimal supersymmetric standard model (NMSSM) [17]. Constraints on models with a light A^0 state have been studied recently within a 2HDM framework with certain assumptions about the coupling and in NMSSM [12,18,19].

Our goal will not be to work in a specific model, but we will assume the couplings of the A^0 to the b quark and the τ lepton to be the same as in the 2HDM. We will assume this 2HDM is part of some extension of the SM. Hence, we will not strictly follow the bounds and constraints obtained in some specific extension of the SM which includes the 2HDM, but will choose values for the parameters in our calculation which are similar to constraints on these parameters in specific NP models. The process $\eta_b \to \tau^+\tau^-$ will proceed through an off-shell A^0, and we will consider both $m_{A^0} < m_{\eta_b}$ and $m_{A^0} > m_{\eta_b}$. In general, there will be mixing between A^0 and the η_b, and as the pseudoscalar state gets close to the η_b mass, the mixing between the states will become important [20]. The calculation of this mixing is model-dependent, and while there are estimates of this mixing in simple quark models, the mixing may be very different in other approaches to the bound state problem in QCD. Hence, we will not take mixing into account in our analysis. Therefore, our results will be reliable when the A^0 mass is away from the η_b mass. We will further assume that the A^0 is narrow and will neglect its width in our calculations. This approximation will be good as long as m_{A^0} is sufficiently away from the η_b mass. When A^0 is produced on-shell, both mixing and width effects will become important, and our results will not be reliable.

There are also models, for example, within SUSY with extra gauged $U(1)$, which have a light axial vector state [21]. These light states can also mediate the process $\eta_b \to \tau^+\tau^-$. Constraints on these models have been studied [22–26]. We will consider $\eta_b \to \tau^+\tau^-$ through the exchange of the axial vector U. To perform our calculations, we will choose the model discussed in Refs. [23,26] and neglect the width of the U boson.

Finally, we note that there are recent dark matter models [27] that also contain light scalar (pseudoscalar) and vector (axial vector) states which may be probed via $\eta_b \to \tau^+\tau^-$. The HyperCP Collaboration has some events for the decay $\Sigma^+ \to p\mu^+\mu^-$ which may be interpreted as evidence for a light pseudoscalar state [28].

This paper is organized in the following manner. In Sec. II, we perform the calculations of the decay $\eta_b \to \tau^+\tau^-$ in the SM and in models with a light pseudoscalar A^0 and a light axial vector U state. In Sec. III, we present the numerical results of the branching ratios for $\eta_b \to \tau^+\tau^-$. Finally, in Sec. IV, we present our conclusion.

II. $\eta_b \to \tau^+\tau^-$ IN THE SM AND NP

In this section, we will study $\eta_b \to \tau^+\tau^-$ in the SM and in models of NP. The η_b is a pseudoscalar and cannot couple to γ directly. Hence, in the SM, $\eta_b \to \tau^+\tau^-$ can only proceed through the exchange of a Z at tree level, and we will calculate the branching ratio for this process in the SM. This decay can also proceed at higher order in the SM through intermediate two-photon states.

In the presence of NP, $\eta_b \to \tau^+\tau^-$ can proceed through the exchange of a light pseudoscalar or a light spin-1 boson with axial vector coupling. We will consider these two NP scenarios in this section. The various tree level contribution to the $\eta_b \to \tau^+\tau^-$ in the SM and NP are shown in Figs. 1 and 2, respectively.

We begin with $\eta_b \to \tau^+\tau^-$ in the SM. We show, in Fig. 1, the decay process $\eta_b \to \tau^+\tau^-$ via the Z boson exchange and through the two-photon intermediate states. The decay rate for the tree level Z exchange process can be obtained as

$$\Gamma^Z(\eta_b \to \tau^+\tau^-) = \frac{G_F^2 M_W^4 m_\tau^2 f_{\eta_b}^2 m_{\eta_b}}{16\pi\cos^4\theta_W}\beta_\tau\left(1 - \frac{m_{\eta_b}^2}{M_Z^2}\right)^2 |a_Z|^2, \tag{3}$$

where θ_W denotes the Weinberg angle, $\beta_\tau = \sqrt{1 - (\frac{2m_\tau}{m_{\eta_b}})^2}$ is the velocity of the τ lepton in the η_b rest frame, and

$$|a_Z|^2 \equiv \frac{1}{(m_{\eta_b}^2 - M_Z^2)^2 + M_Z^2\Gamma_Z^2}. \tag{4}$$

The decay constant f_{η_b} in Eq. (3) is defined as [29]

$$\langle 0|\bar{b}(0)\gamma_\mu\gamma_5 b(0)|\eta_b(q)\rangle = i f_{\eta_b} q_\mu. \tag{5}$$

The process $\eta_b \to \tau^+\tau^-$ can also go via two-photon intermediate states, as shown in Fig. 1. This diagram is dominated by the imaginary part [30], which we can estimate using unitarity [31] to obtain

$$\Gamma^{2\gamma}[\eta_b \to \tau^+\tau^-] \geq \frac{\alpha^2}{2\beta_\tau}\left[\frac{m_\tau}{m_{\eta_b}}\ln\frac{(1+\beta_\tau)}{(1-\beta_\tau)}\right]^2 \Gamma[\eta_b \to \gamma\gamma], \tag{6}$$

where α is the electromagnetic fine structure constant. One can calculate $\Gamma[\eta_b \to \gamma\gamma]$ as

$$\Gamma[\eta_b \to \gamma\gamma] = \frac{\pi\alpha^2 m_{\eta_b} f_{\eta_b}^2}{81 m_b^2}, \tag{7}$$

where we have used the heavy quark limit for the b quark. Since the 2γ exchange contribution is mostly imaginary relative to the Z exchange contribution, to a good approximation the total width $\Gamma_t[\eta_b \to \tau^+\tau^-]$ is,

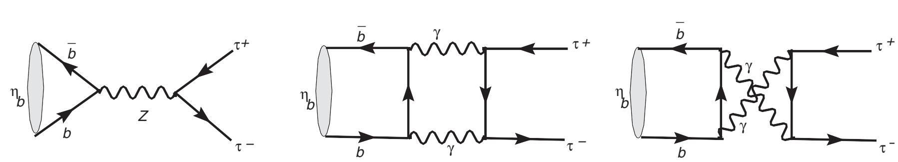

FIG. 1. Various processes contributing to $\eta_b \to \tau^+\tau^-$ in the SM.

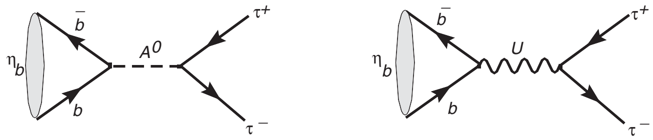

FIG. 2. Various processes contributing to $\eta_b \to \tau^+\tau^-$ in NP.

$$\Gamma_t[\eta_b \to \tau^+\tau^-] \approx \Gamma^Z[\eta_b \to \tau^+\tau^-] + \Gamma^{2\gamma}[\eta_b \to \tau^+\tau^-]. \tag{8}$$

We now turn to NP models and begin with the 2HDM. The couplings of the down-type quarks D and charged leptons ℓ with A^0 in the generic 2HDM model are given by [32]

$$\mathcal{L}_{A^0}^{D,\ell} = \frac{igF_{A^0}}{2M_W}(\bar{D}M_D^{\rm diag}\gamma_5 D + \bar{\ell}M_l^{\rm diag}\gamma_5\ell)A^0, \tag{9}$$

where F_{A^0} is a model-dependent parameter; $M_D^{\rm diag} = (m_d, m_c, m_b)$ and $M_\ell^{\rm diag} = (m_e, m_\mu, m_\tau)$ are the diagonal mass matrices of D and ℓ, respectively. We will 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), $F_{A^0} \equiv -\cot\beta$.

In Fig. 2, we show the decay process $\eta_b \to \tau^+\tau^-$ via the exchange of the CP-odd Higgs scalar A^0. The decay rate for this process can be obtained as

$$\Gamma^{A^0}(\eta_b \to \tau^+\tau^-) = \frac{G_F^2 m_\tau^2 f_{\eta_b}^2 m_{\eta_b}^5}{16\pi}\beta_\tau|a_{A^0}|^2, \tag{10}$$

where the coefficient a_{A^0} depends on the mass m_{A^0} as

$$|a_{A^0}|^2 \equiv \frac{F_{A^0}^4}{(m_{\eta_b}^2 - m_{A^0}^2)^2}. \tag{11}$$

We have assumed that the decay width Γ_{A^0} for the A^0 is negligible. In Eq. (10), we have used

$$\langle 0|\bar{b}(0)\gamma_5 b(0)|\eta_b(q)\rangle = \frac{if_{\eta_b}m_{\eta_b}^2}{2m_b}, \tag{12}$$

where f_{η_b} has been defined in Eq. (5).

Finally, we move to NP models that contain a light spin-1 boson with axial vector couplings. In Fig. 2, we show the decay process $\eta_b \to \tau^+\tau^-$ via the exchange of the light neutral gauge boson U. We write down a model-independent Lagrangian for the U boson, but we assume the structure of the Lagrangian to be similar to the one discussed in Refs. [23–25]. We take the U couplings to the down-type quarks and charged leptons to be given by

$$\mathcal{L}_U^{D,\ell} = f_A^{D,\ell}(\bar{D}\gamma^\mu\gamma_5 D + \bar{\ell}\gamma^\mu\gamma_5\ell)U_\mu, \tag{13}$$

with the axial coupling

$$f_A^{D,\ell} = 2^{-(3/4)}G_F^{1/2}m_U F_U, \tag{14}$$

where m_U denotes the mass of the U boson and F_U denotes a model-dependent parameter. In the specific model [23–25], $F_U \equiv \cos\zeta\tan\beta$.

Again, we will be interested in $F_U > 1$. The decay rate for $\eta_b \to \tau^+\tau^-$ can be obtained as

$$\Gamma^U(\eta_b \to \tau^+\tau^-) = \frac{G_F^2 m_\tau^2 f_{\eta_b}^2 m_{\eta_b}}{16\pi}\beta_\tau(m_U^2 - m_{\eta_b}^2)^2 F_U^4|a_U|^2, \tag{15}$$

where

$$|a_U|^2 = \frac{1}{(m_{\eta_b}^2 - m_U^2)^2 + m_U^2\Gamma_U^2}. \tag{16}$$

Equation (16) can be expanded as

$$|a_U|^2 = \frac{1}{(m_{\eta_b}^2 - m_U^2)^2}(1 - x^2 + \ldots), \tag{17}$$

if $x = \frac{\Gamma_U/m_U}{(1-m_{\eta_b}^2/m_U^2)} < 1$.

Neglecting x, Eq. (15) reduces to

$$\Gamma^U(\eta_b \to \tau^+\tau^-) = \frac{G_F^2 m_\tau^2 f_{\eta_b}^2 m_{\eta_b}}{16\pi}\beta_\tau F_U^4. \tag{18}$$

Thus, Eq. (18) shows that the decay width for $\eta_b \to \tau^+\tau^-$ does not depend on m_U in the approximation of neglecting the width of the U boson. This result is easy to understand. If one increases the mass of the U, then the matrix element for $\eta_b \to \tau^+\tau^-$ is suppressed due to propagator effects. However, the coupling, which is proportional to m_U, increases to compensate for this suppression. The fact that the width for $\eta_b \to \tau^+\tau^-$ is independent of m_U only holds because the η_b is a pseudoscalar.

The result of Eq. (18) does not make sense, as m_U gets sufficiently larger as the couplings in Eq. (14) become nonperturbative. Requiring the couplings to be ≤ 1, one gets the constraints $m_U \leq \frac{4M_W}{gF_U}$. Hence, for $F_U \sim 50$, one can get m_U to be in the GeV range.

It is interesting to note that in the up sector, the behavior for the decay width is different. The coupling of the vector boson to the up-type quark, U, is given by

$$\mathcal{L}_U = f_A^{UP}\bar{U}\gamma^\mu\gamma_5 U U_\mu, \tag{19}$$

with the axial coupling of the up-type quarks

$$f_A^{UP} = 2^{-(3/4)}G_F^{1/2}m_U F'_U. \tag{20}$$

In the model of Refs. [23–25], $F'_U \equiv \cos\zeta\cot\beta$.

For instance, the branching ratio $\mathcal{BR}(\eta_c \to \mu^+\mu^-)$ does not depend on m_U or on $\tan\beta$ and is given as

$$\Gamma^U(\eta_c \to \mu^+\mu^-) = \frac{G_F^2 m_\mu^2 f_{\eta_c}^2 m_{\eta_c}}{16\pi}\bar{\beta}_\tau\cos^4\zeta, \tag{21}$$

where $\bar{\beta}_\tau = \sqrt{1 - (\frac{2m_\mu}{m_{\eta_c}})^2}$ and f_{η_c} is the η_c decay constant. We can see from Eq. (21) that the branching ratio

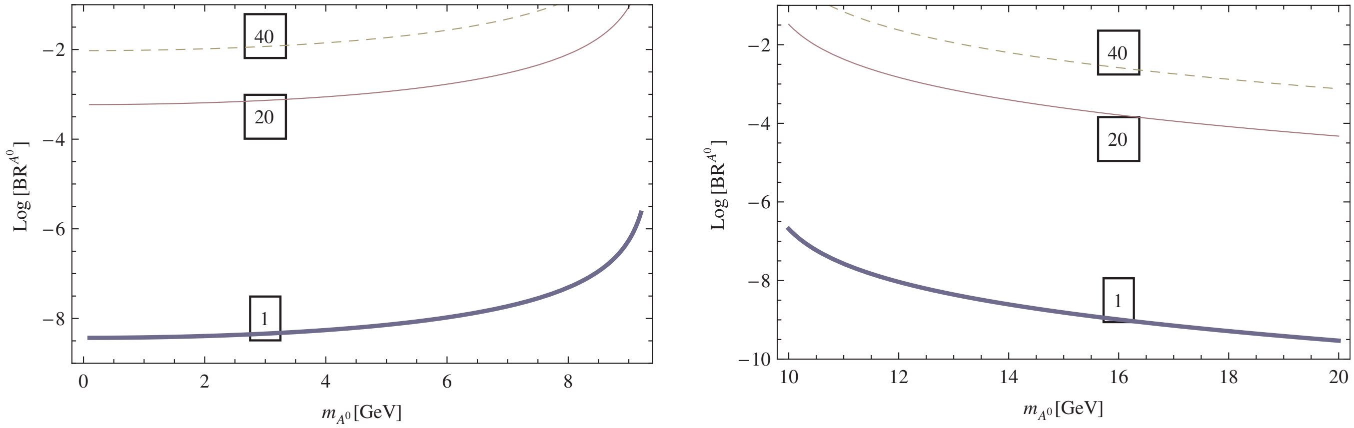

FIG. 3 (color online). The logarithm of $\mathcal{BR}^{A^0}(\eta_b \to \tau^+\tau^-)$ as a function of m_{A^0} for different values of F_{A^0} and $m_{A^0} \in [0.1, 20]$ GeV.

$\mathcal{BR}(\eta_c \to \mu^+\mu^-)$ is much smaller than $\mathcal{BR}(\eta_b \to \tau^+\tau^-)$ if $\tan\beta > 1$ because of the absence of the factor $\tan^4\beta$ in the rate for $\eta_c \to \mu^+\mu^-$.

III. NUMERICAL ANALYSIS

In this section, we present our numerical results. We take the average $\eta_b(1S)$ mass to be $m_{\eta_b} = 9390.8 \pm 3.2$ MeV [3], the decay constant $f_{\eta_b} = (705 \pm 27)$ MeV [33], and the width to be $\Gamma_{\eta_b} \approx 10$ MeV [34].

In the SM, at tree level, $\eta_b \to \tau^+\tau^-$ goes through the exchange of a Z boson, and we obtain a tiny branching ratio, $\mathcal{BR}^Z(\eta_b \to \tau^+\tau^-) = 3.8 \times 10^{-9}$. In our calculation, we have used $\Gamma_Z = 2.4952 \pm 0.0023$ GeV [15]. For the two-photon contribution to $\eta_b \to \tau^+\tau^-$, we obtain, using Eqs. (6) and (7) and Eq. (7), $\mathcal{BR}^{2\gamma}[\eta_b \to \tau^+\tau^-] \geq 4.6 \times 10^{-10}$ for $m_b = 4.8$ GeV. Using Eq. (8), the total branching ratio for $\eta_b \to \tau^+\tau^-$ is $\approx 4.3 \times 10^{-9}$.

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. The branching ratio, $\mathcal{BR}^{A^0}$, is plotted for various values of the A^0 mass, which we take from 0.1 to 20 GeV, and for various values of F_{A^0}. As the mass of the A^0 approaches the mass of the η_b, the branching ratio increases and blows up at $m_{A^0} = m_{\eta_b}$. This behavior clearly does not represent the physical situation, because in this region, the width of the A^0 and mixing effects of the A^0 with η_b become important and regularize the A^0 contribution. We observe in Fig. 3 that the branching ratio $\sim F_{A^0}^4$ is very sensitive to F_{A^0}. The branching ratio is relatively less sensitive to the mass m_A^0. We see from the plots in Fig. 3 that the branching ratio for $\eta_b \to \tau^+\tau^-$, through the A^0 exchange, can be considerably larger than the SM branching ratios and can vary from $\sim 10^{-8}$ to the experimental bound of 8% for $F_{A^0} = 40$. Since we have neglected the width and mixing effects, our predictions are no longer reliable as the mass of the A^0 approaches the mass of the η_b. The mixing effects are model-dependent and, as an example, for the model for mixing employed in Ref. [20], the effects of mixing are important in the m_{A^0} mass range of 9.4–10.5 GeV. We see from Fig. 3 that even outside this range, the branching ratio for $\eta_b \to \tau^+\tau^-$ can

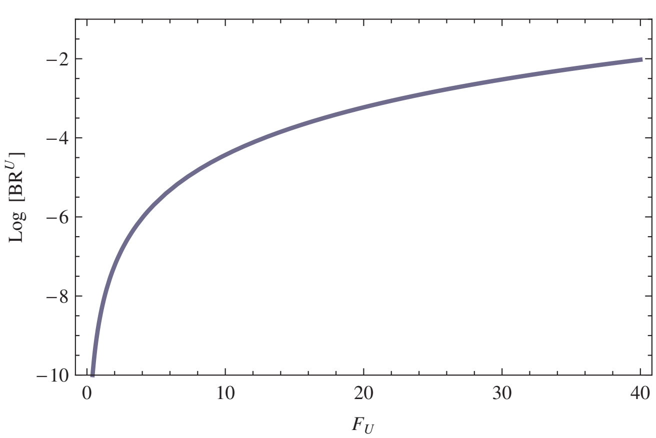

FIG. 4 (color online). The logarithm of $\mathcal{BR}^U(\eta_b \to \tau^+\tau^-)$ as a function of F_U.

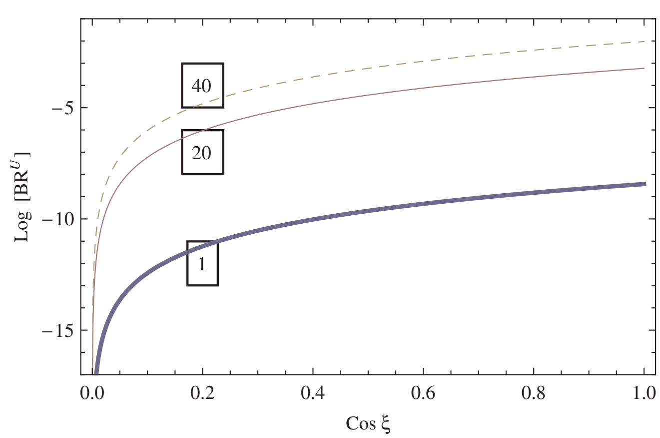

FIG. 5 (color online). The logarithm of $\mathcal{BR}^U(\eta_b \to \tau^+\tau^-)$ as a function of $\cos\zeta$ for different values of $\tan\beta$ and $\cos\zeta \in [0, 1]$.

be significant, and we expect the same also to be true in the mass range, where mixing effects are important.

As discussed in the previous section, the branching ratio for the decay $\mathcal{BR}^U(\eta_b \to \tau^+\tau^-)$ is independent of the mass of the gauge boson U in the approximation of neglecting the width of the U boson. We next plot in Fig. 4 the logarithm of the branching ratio for $\eta_b \to \tau^+\tau^-$ versus F_U. Working in a specific model [23–25] $F_U \equiv \cos\zeta \tan\beta$, we plot the branching ratio versus the invisibility factor $\cos\zeta$ for different values of $\tan\beta$ in Fig. 5. Again, we observe that the branching ratio can vary over a wide range and can be much larger than the SM prediction.

IV. CONCLUSION

In this paper, we explored the decay $\eta_b \to \tau^+\tau^-$ as a probe for a light pseudoscalar or a light axial vector state. We estimated the SM branching ratios for $\eta_b \to \tau^+\tau^-$ via the Z exchange and the two-photon intermediate state and found it to be very small at $\sim 4 \times 10^{-9}$. We then considered the decay process $\eta_b \to \tau^+\tau^-$ mediated via the pseudoscalar Higgs boson A^0 in a 2HDM-type NP model. We found that the branching ratio for $\eta_b \to \tau^+\tau^-$ can be substantially larger than the SM prediction and can reach the experimental bound of 8%. Working in a specific model containing a light axial vector state, U, a similar result was obtained for the branching ratio of $\eta_b \to \tau^+\tau^-$. We also obtained an interesting result—that the $\mathcal{BR}^U(\eta_b \to \tau^+\tau^-)$ is independent of the mass of the U boson if the width of the U is neglected. This result followed from the fact that the axial U boson couplings to fermions were proportional to the mass m_U and the fact that η_b is a pseudoscalar. A constraint on the U boson mass could be obtained by requiring its coupling to fermions to be ≤ 1. In light of the results obtained in the paper, an experimental measurement of the branching ratio for $\eta_b \to \tau^+\tau^-$ is strongly desirable, as this measurement might reveal the presence of light, $\sim$GeV, pseudoscalar, or axial vector states. The experimental measurements of $\eta_b \to \tau^+\tau^-$ may be feasible at planned high-luminosity B factories and at hadron colliders such as the Tevatron and the LHC, specifically if the branching ratios are much larger than the SM rate.

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