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Bottomonium spectroscopy with mixing of η_b states and a light CP-odd Higgs

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The mass of the $\eta_b(1S)$, measured recently by BABAR, is significantly lower than expected from QCD predictions for the $\Upsilon(1S)$ – $\eta_b(1S)$ hyperfine splitting. We suggest that the observed $\eta_b(1S)$ mass is shifted downwards due to a mixing with a CP-odd Higgs scalar A with a mass m_A in the range 9.4 – 10.5 GeV compatible with LEP, CLEO and BABAR constraints. We determine the resulting predictions for the spectrum of the $\eta_b(nS) - A$ system and the branching ratios into $\tau^+\tau^-$ as functions of m_A.

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The BABAR collaboration has recently determined the $\eta_b(1S)$ mass $m_{\eta_b(1S)}$ with an error of only a few MeV in radiative decays $\Upsilon \to \gamma\,\eta_b$ of excited Υ states and the observation of peaks in the photon energy spectrum. The result from $\Upsilon(3S)$ decays is $m_{\eta_b(1S)} = 9388.9^{+3.1}_{-2.3}$ (stat)$\pm$ 2.7 (syst) MeV [1], and the result from $\Upsilon(2S)$ decays is $m_{\eta_b(1S)} = 9392.9^{+4.6}_{-4.8}$ (stat) $\pm$ 1.9 (syst) MeV [2]. The average gives [2]

$$m_{\eta_b(1S)} = 9390.9 \pm 3.1 \text{ MeV} , \qquad (1)$$

implying a hyperfine splitting $E_{hfs}(1S) = m_{\Upsilon(1S)} - m_{\eta_b(1S)}$ of

$$E^{exp}_{hfs}(1S) = 69.9 \pm 3.1 \text{ MeV} . \qquad (2)$$

This result can be compared to predictions from QCD. Recent results based on perturbative QCD are in good agreement with each other and give $E_{hfs}(1S) = 44 \pm 11$ MeV [3] and $E_{hfs}(1S) = 39 \pm 14$ MeV [4] (whereas $E_{hfs}(1S)$ varies over a wider range in phenomenological models [5]); in the following we will use tentatively an average value

$$E^{pQCD}_{hfs}(1S) = 42 \pm 13 \text{ MeV} \qquad (3)$$

which is about two standard deviations away from the experimental result (2). The most recent result from (unquenched) lattice QCD is [6]

$$E^{latQCD}_{hfs}(1S) = 61 \pm 14 \text{ MeV} , \qquad (4)$$

which is within $1\,\sigma$ of (2). However, the hyperfine splitting is quite sensitive to short distances or hard quark momenta. It has been argued in [7] that the perturbative results in [4] can be used for short distance corrections of the Wilson coefficient of the corresponding spin-flip operator measured on the lattice. The additional contribution $\delta^{hard}E_{hfs}(1S)$ to $E^{latQCD}_{hfs}(1S)$ has been estimated as $\delta^{hard}E_{hfs}(1S) \sim -20$ MeV [7], which brings the lattice result (4) in good agreement with the perturbative result (3). Although this conclusion needs to be checked within perturbation theory with lattice regularization, we consider it as a support for the perturbative QCD result (3).

Whereas an explanation of the discrepancy between (2) and (3) within QCD is not excluded at present, we will elaborate below the consequences of an explanation of this discrepancy due to new physics in the form of a mixing of the η_b states with a CP-odd Higgs scalar A with a mass m_A in the range $9.4-10.5$ GeV: as a result of such a mixing, the masses of the η_b - like eigenstates of the full mass matrix can differ considerably from their values in pure QCD without the presence of the CP-odd Higgs A [8, 9, 10], and the mass of the state interpreted as the $\eta_b(1S)$ can be smaller than expected if m_A is somewhat above 9.4 GeV.

In such a scenario, the masses of the states interpreted as $\eta_b(2S)$ and $\eta_b(3S)$ can also be affected, and all states can have non-negligible branching ratios into $\tau^+\tau^-$ due to their mixing with A. According to recent results of BABAR [11], the corresponding branching ratio of the observed state is below 8% at 90% confidence level. The branching ratio into $\mu^+\mu^-$ would be smaller by a factor m_μ^2/m_τ^2, and well below the present upper limit [12]. The investigation of these phenomena is the purpose of the present paper.

A relatively light CP-odd Higgs scalar can appear, e. g., in non-minimal supersymmetric extensions of the Standard Model (SM) as the NMSSM (Next-to-Minimal Supersymmetric Standard Model) [10, 13, 14, 15, 16, 17]. Its mass has to satisfy constraints from LEP, where it could have been produced in $e^+e^- \to Z^* \to Z\,H$ and $H \to A\,A$ (where H is a CP-even Higgs scalar). For $m_A > 10.5$ GeV – where A would decay dominantly into $b\,\bar{b}$ – and $m_H < 110$ GeV, corresponding LEP constraints are quite strong [18]. For $2\,m_\tau < m_A < 10.5$ GeV, A would decay dominantly into $\tau^+\tau^-$ and values for m_H down to ~ 86 GeV are allowed [18] even if H couples to the Z boson with the strength of a SM Higgs boson.

In fact, searches for $e^+e^- \to Z^* \to Z\,H$ with $H \to b\,\bar{b}$ indicate a light excess of events (of $\sim 2.3\,\sigma$ significance) for $m_H \sim 95 - 100$ GeV [18], which could be explained by a reduced branching ratio $BR(H \to b\,\bar{b}) \sim 0.1$ and a dominant branching ratio $BR(H \to A\,A) \sim 0.9$ [16, 17] if A decays dominantly into $\tau^+\tau^-$. The possible explanation of this excess of events at LEP is an additional motivation for a CP-odd Higgs scalar with a mass below 10.5 GeV. Allowing for m_H somewhat below 100 GeV, such a scenario would also alleviate the "little fine tuning problem" of supersymmetric extensions of the SM [15, 16, 17].

Finally CLEO [19] and BABAR [11] have searched for a light CP-odd scalar A with $A \to \tau^+\tau^-$ in radiative $\Upsilon(1S)$ and $\Upsilon(3S)$ decays respectively, which mainly constrains the range $m_A \lesssim 9.4$ GeV. In this work we focus on the range 9.4 GeV $\lesssim m_A < 10.5$ GeV, which would be the most relevant for strong $\eta_b - A$ mixing effects as advocated in [10].

Apart from m_A, such mixing effects depend in a calculable way on the model-dependent coupling of A to b quarks. The corresponding coupling, normalized with respect to the coupling of the SM Higgs scalar to b quarks, will be denoted by X_d. In models with two Higgs doublets H_u and H_d, where H_u couples to up quarks and H_d to down quarks and leptons, one has [10, 13, 14, 15, 16, 17]

$$X_d = \cos\theta_A \ \tan\beta \tag{5}$$

where $\cos\theta_A$ denotes the SU(2) doublet component of the CP-odd scalar A, and $\tan\beta = \langle H_u\rangle / \langle H_d\rangle$. For $\tan\beta \gg 1$, X_d can equally satisfy $X_d \gg 1$ [10].

Below we proceed as follows: i) First we assume that, in the absence of a CP-odd Higgs scalar, $m_{\eta_b^0(1S)}$ would have a value compatible with (3), i.e.

$$m_{\eta_b^0(1S)} \sim 9418 \pm 13 \text{ MeV} . \tag{6}$$

ii) We diagonalise the $\eta_b(nS) - A$ mass matrix ($n = 1, 2, 3$) and require that one eigenvalue coincides with the mass measured by BABAR within errors (1); this condition gives us an allowed strip in the $X_d - m_A$ plane, which depends only weakly on the assumed masses of $\eta_b^0(2S)$ and $\eta_b^0(3S)$.

The resulting values of X_d as function of m_A allow to determine the remaining eigenvalues of the mass matrix, and the decompositions of the eigenvectors in terms of A and $\eta_b^0(nS)$. Finally the A components of the eigenstates allow us to determine their partial widths and to estimate the branching ratios into $\tau^+\tau^-$.

For the $\eta_b^0(1S) - \eta_b^0(2S) - \eta_b^0(3S) - A$ mass matrix we make the ansatz

$$\mathcal{M}^2 = \begin{pmatrix} m^2_{\eta_b^0(1S)} & 0 & 0 & \delta m_1^2 \\ 0 & m^2_{\eta_b^0(2S)} & 0 & \delta m_2^2 \\ 0 & 0 & m^2_{\eta_b^0(3S)} & \delta m_3^2 \\ \delta m_1^2 & \delta m_2^2 & \delta m_3^2 & m_A^2 \end{pmatrix} . \tag{7}$$

The diagonal elements $m^2_{\eta_b^0(nS)}$ are assumed to be known (within errors) from QCD with (6) for $m_{\eta_b^0(1S)}$. The results depend only weakly on the $\eta_b^0(2S)$ and $\eta_b^0(3S)$ masses, for which we take [3] $m_{\eta_b^0(2S)} = 10002$ MeV, $m_{\eta_b^0(3S)} = 10343$ MeV. (We neglect finite width effects in $\mathcal{M}^2$, and take real matrix elements.)

The off-diagonal elements δm_n^2 can be computed in the framework of a non-relativistic quark potential model in terms of the radial wave functions at the origin [8, 9]. These can be considered as identical for vector and pseudocalar states, and be determined from the measured $\Upsilon \to e^+e^-$ decay widths. Substituting recent values for these widths (see [10] for details) we find

$$\begin{aligned} \delta m_1^2 &\simeq (0.14 \pm 10\%) \text{ GeV}^2 \times X_d , \\ \delta m_2^2 &\simeq (0.11 \pm 10\%) \text{ GeV}^2 \times X_d , \\ \delta m_3^2 &\simeq (0.10 \pm 10\%) \text{ GeV}^2 \times X_d . \end{aligned} \tag{8}$$

We estimated the errors from higher order QCD corrections to the relation between the radial wave functions at the origin and $\Gamma(\Upsilon(nS) \to e^+e^-)$ [8, 9] to be $\sim 10\%$. (These errors play only a minor role for our results.)

In order that one eigenvalue of $\mathcal{M}^2$ coincides with the BABAR result (1) subsequently denoted as m_{obs}, m_A in (7) has to satisfy

$$\begin{aligned} m_A^2 \ = \ & m_{obs}^2 + \frac{\delta m_1^4}{m^2_{\eta_b^0(1S)} - m^2_{obs}} \\ & + \frac{\delta m_2^4}{m^2_{\eta_b^0(2S)} - m^2_{obs}} + \frac{\delta m_3^4}{m^2_{\eta_b^0(3S)} - m^2_{obs}} . \end{aligned} \tag{9}$$

Once m_A is expressed in terms of X_d, X_d remains the only unknown parameter in $\mathcal{M}^2$. Varying m_{obs} within the errors in (1), $m_{\eta_b^0(1S)}$ within the errors in (6) and δm_i^2 within the errors in (8), we obtain for X_d as a function of m_A the result shown in Fig. 1. (Recall that m_A – or, the heaviest eigenstate of $\mathcal{M}^2$ with a large A component – has to be below 10.5 GeV in order to satisfy LEP constraints in the presence of a Higgs scalar with a mass around 100 GeV.)

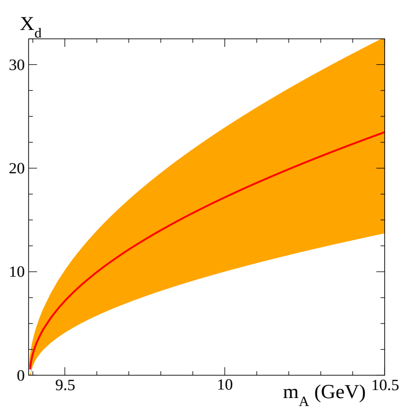

FIG. 1: X_d as a function of m_A (in GeV) such that one eigenvalue of $\mathcal{M}^2$ coincides with the BABAR result (1).

Now the masses of all 4 eigenstates of $\mathcal{M}^2$ can be computed, which are shown together with the error bands from m_{obs}, $m_{\eta_b^0(1S)}$ and δm_i^2 (in orange/grey) in Fig. 2 as functions of m_A. Henceforth we denote the 4 eigenstates of $\mathcal{M}^2$ by η_i, $i = 1 \ldots 4$ where, by construction, $m_{\eta_1} \equiv m_{obs}$. For clarity we have indicated in Fig. 2 our assumed values for $m_{\eta_b^0(nS)}$ as horizontal dashed lines. For m_A not far above 9.4 GeV (where X_d is relatively small) the effects of the mixing on the states $\eta_b^0(2S)$ and $\eta_b^0(3S)$ are negligible, but for larger m_A the spectrum can differ considerably from the one expected without the presence of A.

Now we consider the branching ratios of the eigenstates into $\tau^+\tau^-$, which are induced by their A-components.

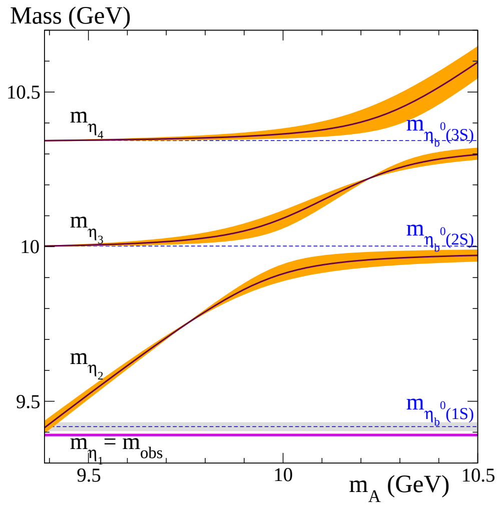

FIG. 2: The masses of all eigenstates as function of m_A.

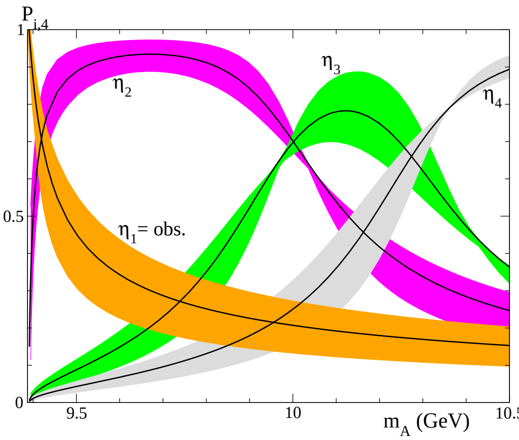

FIG. 3: The A-components $|P_{i,4}|$ for all 4 eigenstates as functions of m_A.

The decomposition of the eigenstates into the states before mixing can be written as

$$\eta_i = P_{i,1}\, \eta_b^0(1S) + P_{i,2}\, \eta_b^0(2S) + P_{i,3}\, \eta_b^0(3S) + P_{i,4}\, A \,. \tag{10}$$

It turns out that the coefficients $P_{i,j}$ can be expressed analytically in terms of the eigenvalues $m^2_{\eta_i}$ of $\mathcal{M}^2$:

$$\begin{aligned} P_{i,4} &= \left[1 + \frac{\delta m_1^4}{(m^2_{\eta_b^0(1S)} - m^2_{\eta_i})^2} \right. \\ &\quad \left. + \frac{\delta m_2^4}{(m^2_{\eta_b^0(2S)} - m^2_{\eta_i})^2} + \frac{\delta m_3^4}{(m^2_{\eta_b^0(3S)} - m^2_{\eta_i})^2}\right]^{-1/2} ; \\ P_{i,j} &= \frac{-\delta m_j^2}{m^2_{\eta_b^0(jS)} - m^2_{\eta_i}} P_{i,4} \quad \text{for } j = 1,2,3 \,. \end{aligned} \tag{11}$$

In Fig. 3 we show our results for the A-components $P_{i,4}$ for all 4 eigenstates together with the error bands from m_{obs}, $m_{\eta_b^0(1S)}$ and δm_i^2.

In the case of $\eta_1 \equiv \eta_{obs}$, only the coefficients $P_{1,1}$ and $P_{1,4}$ differ significantly from 0. This allows to express the branching ratio $BR(\eta_1 \to \tau^+\tau^-)$ as

$$BR(\eta_1 \to \tau^+\tau^-) = \frac{P_{1,4}^2 \Gamma_A^{\tau\tau}}{P_{1,1}^2 \Gamma_{\eta_b^0(1S)} + P_{1,4}^2 \Gamma_A^{tot}} \tag{12}$$

where $\Gamma_A^{\tau\tau}$ is the partial width for $A \to \tau^+\tau^-$, $\Gamma_{\eta_b^0(1S)}$ the width of the state $\eta_b^0(1S)$ without mixing with A, and Γ_A^{tot} the total width of A (without mixing). For $\Gamma_A^{\tau\tau}$ we have

$$\Gamma_A^{\tau\tau} = X_d^2 \frac{G_F m_\tau^2 M}{4\sqrt{2}\pi} \sqrt{1 - 4\frac{m_\tau^2}{M^2}} \tag{13}$$

where we would have to identify M with m_A if A decays on its mass shell. Since the dependence on M originates from phase space integrals, M has to be identified with the mass of the decaying state (which is actually always close to m_A) in our case. Since the $BR(A \to \tau^+\tau^-)$ is typically $\sim 90\%$ (for large $\tan\beta$), we take $\Gamma_A^{tot} \sim 1.1\, \Gamma_A^{\tau\tau}$ in (12).

It is remarkable that, once we insert eqs. (11) and (13) into the expression (12) for $BR(\eta_1 \to \tau^+\tau^-)$, all dependence on X_d cancels: even for $m_A \to 9.4$ GeV where $P_{1,4}^2 \to 1$, one has $\Gamma_A^{tot} \to 0$ from $X_d \to 0$, and hence the first term in the denominator in (12) always dominates. The numerical result is

$$BR(\eta_1 \to \tau^+\tau^-) = (2.4^{+2.3}_{-2.0}) \times 10^{-2} \times \left(\frac{10\,\mathrm{MeV}}{\Gamma_{\eta_b^0(1S)}}\right) . \tag{14}$$

Since the analytic expression for this BR is proportional to $(m^2_{obs} - m^2_{\eta_b^0(1S)})^2$ in our approach, its lowest possible value is quite sensitive to the smallest allowed value for this difference.

Using $\Gamma_{\eta_b^0}(nS)/\Gamma_{\eta_c}(nS) \simeq (m_b/m_c)[\alpha_s(m_b)/\alpha_s(m_c)]^5 \simeq 0.25 - 0.75$ [20] and $\Gamma_{\eta_c}(1S) = 26.7 \pm 3$ MeV [21] we estimate $\Gamma_{\eta_b^0(1S)} \sim 5 - 20$ MeV. Hence the predicted branching ratio is compatible with BABAR upper limit of 8% [11].

Turning to the remaining heavier eigenstates, expressions similar to (12) are always very good approximations for the branching ratios into $\tau^+\tau^-$, since the eigenstates consist essentially of just one $\eta_b^0(nS)$ state and the CP-odd Higgs A. However, the branching ratios vary now with m_{η_i} and hence with m_A; the results are shown graphically in Fig. 4 assuming $\Gamma_{\eta_b^0(1S)} \sim 10$ MeV and $\Gamma_{\eta_b^0(2S)} \sim \Gamma_{\eta_b^0(3S)} \sim 5$ MeV. For larger (smaller) total widths these branching ratios would be somewhat smaller (larger).

An important issue is the production rate of the eigenstates η_i in radiative decays of excited Υ states, notably in $\Upsilon(3S) \to \gamma\,\eta_i$. For a pure CP-odd scalar A, the $BR(\Upsilon(3S) \to \gamma\, A)$ is given by the Wilczek formula [22]

$$\frac{BR\,(\Upsilon(nS) \to \gamma\, A)}{BR\,(\Upsilon(nS) \to \mu^+\mu^-)} = \frac{G_F m_b^2 X_d^2}{\sqrt{2}\pi\alpha}\left(1 - \frac{m_A^2}{m^2_{\Upsilon(3S)}}\right) \times F \,. \tag{15}$$

F is an m_A dependent correction factor, which includes three kinds of corrections (the relevant formulas are summarized in [23]): bound state, QCD, and relativistic corrections. Unfortunately these corrections become unreliable for $m_A \gtrsim 8$ GeV; a naïve extrapolation of the known

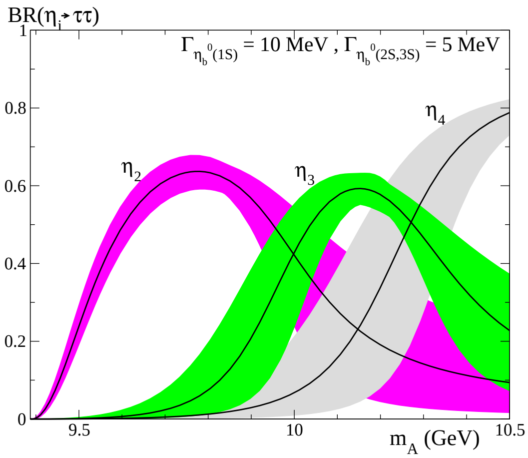

FIG. 4: The branching ratios into $\tau^+\tau^-$ for the eigenstates η_2, η_3 and η_4 as functions of m_A.

corrections leads to a vanishing correction factor F for $m_A \gtrsim 8.8$ GeV [10] as relevant here.

Thus it is difficult to predict the branching ratios $BR(\Upsilon(3S) \to \gamma\,\eta_i)$: if the $BR(\Upsilon(3S) \to \gamma\,A)$ is assumed to vanish, the production of the states η_i has to rely on their $\eta_b^0(nS)$ components. Otherwise, negative interference effects could appear leading to suppressed branching ratios. Hence it cannot be guaranteed that the (kinematically accessible) part of the spectrum shown in Fig. 2 is actually visible in radiative $\Upsilon(3S)$ decays in the form of a peak in the photon energy spectrum [11].

In view of the possibly quite low photon energies and/or large backgrounds, the photons can well escape undetected even if the process $\Upsilon(3S) \to \gamma\,\eta_i$ occurs with a non-negligible rate. In this case, as advocated in [24], the A components of η_i can still manifest themselves through a breakdown of lepton universality in the form of an excess of $\tau^+\tau^-$ final states in $\Upsilon(3S) \to l^+ l^-$ [9, 10]. However, in the case of a vanishing correction factor F in (15), this phenomenon would disappear as well.

On the other hand, if the existence of a CP-odd scalar A and a CP-even scalar H with $m_H \sim 95 - 100$ GeV (and a dominant $H \to A\,A$ branching ratio of $\sim 90\%$) is responsible for the excess of events at LEP as noted above, it becomes important to test this scenario at the LHC: The standard search channels for a SM-like CP-even scalar H would fail, and the final states from $H \to A\,A$ with m_A below the $b\,\bar{b}$ threshold would be difficult to detect. Proposals for a verification of this scenario at the LHC have been made recently in [25, 26, 27].

To conclude, if the mixing with a CP-odd Higgs scalar is responsible for the discrepancy between the BABAR measurement of $m_{\eta_b(1S)}$ and the expectations from QCD, it can manifest itself in the form of a completely distorted spectrum of states as shown in Fig. 2. The branching ratios into $\tau^+\tau^-$ would be non-vanishing, albeit below the present experimental upper limit (for the lowest lying state). These manifestations of a light CP-odd Higgs scalar in Υ physics at (Super) B factories [28] would be complementary to its possible discovery at the LHC.

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