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## Reduced branching ratio for $H \rightarrow AA \rightarrow 4\tau$ from $A - \eta_b$ mixing

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ABSTRACT: Models with an extended Higgs sector, as the NMSSM, allow for scenarios where the Standard Model-like CP-even Higgs boson  $H$  decays dominantly as  $H \rightarrow AA \rightarrow 4\tau$  where  $A$  is a light CP-odd Higgs boson. Tight constraints on this scenario in the form of lower bounds on  $M_H$  have recently been published by the ALEPH group. We show that, due to  $A - \eta_b$  mixing, the branching ratio  $H \rightarrow AA \rightarrow 4\tau$  is strongly reduced for  $M_A$  in the range  $9 - 10.5$  GeV. This is the range of  $M_A$  in which the tension between the observed  $\eta_b(1S)$  mass and its prediction based on QCD can be resolved due to mixing, and which is thus still consistent with a light CP-even Higgs boson  $H$  satisfying LEP constraints with a mass well below 114 GeV. This result is practically independent from the coupling of  $A$  to  $b$  quarks.

KEYWORDS: Supersymmetry Phenomenology

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**1 Introduction**

One of the main goals of the Large Hadron Collider (LHC) is the detection of the Higgs boson, or of at least one of several Higgs bosons if corresponding extensions of the Standard Model (SM) are realized in nature. These searches depend crucially on the Higgs production cross sections and the Higgs decays.

In the case of the SM, the production cross sections and decay branching ratios are quite well known as functions of the still unknown Higgs mass [1]. In the Minimal Supersymmetric Standard Model (MSSM) with its extended Higgs sector, these quantities have been studied as well and it seems that at least one of the Higgs bosons cannot be missed at the LHC [2]. There exist, however, well motivated scenarios with somewhat more extended Higgs sectors, as the Next-to-Minimal Supersymmetric Standard Model (NMSSM, see [3, 4] for recent reviews), where the Higgs decays can differ strongly from both the SM and the MSSM. It is very important to be aware of the possibility of such unconventional Higgs decays; the absence of a signal in standard Higgs search channels may otherwise be completely misinterpreted.

The Higgs sector of the NMSSM consists of two SU(2) doublets  $H_u$  and  $H_d$  (as in the MSSM), and one additional gauge singlet  $S$ . Due to its coupling  $\lambda SH_u H_d$  in the superpotential, a vacuum expectation value (vev)  $s$  of  $S$  generates a supersymmetric mass term  $\mu_{\text{eff}} = \lambda s$  for  $H_u$  and  $H_d$ . Since  $s$  and hence  $\mu_{\text{eff}}$  are naturally of the order of the soft Susy breaking terms  $\sim M_{\text{Susy}}$ , this solves the so-called  $\mu$ -problem of the MSSM [5]. (This remains true in the limit  $\lambda, \kappa \rightarrow 0$ , where  $\kappa$  is the singlet self-coupling in the superpotential, leading to  $s \sim M_{\text{Susy}}/\kappa$ , but  $\mu_{\text{eff}} \sim (\lambda/\kappa)M_{\text{Susy}} \sim M_{\text{Susy}}$ .) Furthermore, in its simplest  $Z_3$  invariant version, the superpotential of the NMSSM is scale invariant; it is in fact the simplest phenomenologically acceptable supersymmetric extension of the SM with this property.

The physical neutral Higgs sector in the NMSSM consists of 3 CP even and 2 CP odd states. (Here we do not consider the possibility of CP violation in the Higgs sector.) In general, these states are mixtures of the corresponding CP even or CP odd components of  $H_u$ ,  $H_d$  and  $S$ , without the CP odd Goldstone boson swallowed by the massive  $Z$  boson. Often, one of the CP even states is SM like, i.e. with similar couplings to gauge bosons as the SM Higgs boson (but with possibly enhanced couplings to quarks and leptons), with

a mass bounded from above by  $\sim 140$  GeV [6]. At first sight, the detection at the LHC of this Higgs boson — denoted subsequently by  $H$  for simplicity — seems to be guaranteed, given the lower LEP bound of  $\sim 114$  GeV on masses of Higgs bosons with SM like couplings to the  $Z$  boson and SM like decays.

However, the lighter of the two CP odd states (denoted by  $A_1$ ) could have a mass  $M_{A_1}$  below half of the mass  $M_H$  of  $H$  [7, 8]. Then,  $H$  would decay dominantly as  $H \rightarrow A_1 A_1$ , since this coupling is typically larger than the coupling of  $H$  to  $b$  quarks [7, 8]. Such a decay of  $H$  would have important consequences both for lower bounds on its mass from searches at LEP, and for its detection at the LHC. Now the  $H$  final decay products depend on  $M_{A_1}$ : for  $M_{A_1} \gtrsim 10.5$  GeV, they consist mainly of 4  $b$  quarks (with some  $2b + 2\tau$  admixture), whereas for  $3.5$  GeV  $\lesssim M_{A_1} \lesssim 10.5$  GeV, they consist mainly of 4  $\tau$  leptons (with some small  $2\tau + 2\mu$  admixture). In fact,  $H \rightarrow 4b$  decays have also been searched for by OPAL and DELPHI at LEP [9, 10] implying  $M_H \gtrsim 110$  GeV if  $H$  has SM like couplings to the  $Z$  boson [11]. On the other hand, LEP constraints on  $H \rightarrow 4\tau$  decays were relatively weak, allowing for  $M_H$  as low as  $\sim 90$  GeV [11].

This led to the scenario advocated in [13–16] (see also [17]) with  $M_H \lesssim 110$  GeV,  $M_{A_1} \lesssim 10.5$  GeV, a dominant (but not exclusive) decay  $H \rightarrow A_1 A_1 \rightarrow 4$  leptons and a low finetuning among the soft Susy breaking parameters due to the relatively low mass of  $H$ . A remaining small branching ratio for  $H \rightarrow 2b$  could explain the  $2\sigma$  excess observed in this channel for  $M_H \sim 100$  GeV [11, 14].

The final state  $H \rightarrow 4\tau$  has recently been reanalysed by the ALEPH group [12] implying upper bounds on  $\xi^2 = \frac{\sigma(e^+e^- \rightarrow ZH)}{\sigma_{\text{SM}}(e^+e^- \rightarrow ZH)} \times BR(H \rightarrow 2A_1) \times BR(A_1 \rightarrow \tau^+ \tau^-)^2$  as function of  $M_H$  and  $M_{A_1}$ . These bounds seem to impose strong constraints on the above scenario, unless  $\sigma(e^+e^- \rightarrow ZH)$  and/or the  $BR(H \rightarrow 2A_1)$  and/or the  $BR(A_1 \rightarrow \tau^+ \tau^-)$  are smaller than naively expected [18].

A light CP odd scalar  $A_1$  would also have important consequences for the physics of  $b\bar{b}$  bound states. These effects depend on the coupling of  $A_1$  to  $b$  quarks. Normalized relative to the coupling of the SM Higgs boson, the coupling of  $A_1$  to  $b$  quarks is given by  $X_d$  with

$$X_d = \cos \theta_A \tan \beta, \tag{1.1}$$

where  $\cos \theta_A$  denotes the SU(2) doublet component of  $A_1$ , and  $\tan \beta$  is the usual ratio of Higgs vevs  $v_u/v_d$ . For  $\tan \beta$  much larger than 1,  $X_d$  could satisfy  $X_d \gg 1$  as well. ( $X_d$  is simultaneously the coupling of  $A_1$  to leptons normalized relative to the coupling of the SM Higgs boson.)

In fact the relation (1.1) is valid for  $A_1$  in any extension of the SM with two Higgs doublets  $H_u$  (coupling exclusively to up-type quarks) and  $H_d$  (coupling exclusively to down-type quarks and leptons), but arbitrary singlets. Our subsequent results depend only on  $M_{A_1}$  and  $X_d$ , and are valid for any such models. In the NMSSM, a light CP odd scalar  $A_1$  can play the role of a pseudo Goldstone boson of an approximate R- or Peccei-Quinn symmetry [7, 8]. Then, however, one always has  $\cos \theta_A \sim 1/\tan \beta$  [4] and hence  $X_d \lesssim 1$ .

Since the pseudoscalar  $b\bar{b}$  bound states  $\eta_b(nS)$  have the same quantum numbers as a CP odd Higgs  $A_1$ , the states  $\eta_b(nS)$  and  $A_1$  can mix [19–22] with important consequences both for the mass spectrum and the decays of the physical eigenstates. A state  $\eta_b(1S)$  has

been observed in radiative  $\Upsilon(3S)$  and  $\Upsilon(2S)$  decays by BABAR [23, 24], with the result that its mass of  $9390.9 \pm 3.1$  MeV is below the one expected from most QCD predictions for the  $\Upsilon(1S) - \eta_b(1S)$  hyperfine splitting [25–27]. Indeed, such a mass shift could be explained by the mixing of  $\eta_b(1S)$  with  $A_1$  provided  $M_{A_1}$  (before mixing) is in the 9.4 – 10.5 GeV range [22].

On the other hand,  $A_1$  can be searched for in radiative decays  $\Upsilon(nS) \rightarrow \gamma A_1$ ,  $A_1 \rightarrow 2$  leptons. (See [28] for a discussion of  $\eta_b \rightarrow \tau^+ \tau^-$  mediated by  $A_1$ .) Unsuccessful searches by CLEO [29] and BABAR [30, 31] lead to upper bounds on  $X_d$  as function of  $M_{A_1}$ , which have been studied in [18, 21, 32] for  $M_{A_1} \lesssim 9$  GeV where the  $\eta_b(nS) - A_1$  mixing is not very relevant. Notably for  $M_{A_1}$  below the  $2\tau$  threshold, where  $A_1$  has a large branching fraction into two muons, these bounds are quite strong and imply  $X_d \lesssim 0.5$ . Upper bounds on  $X_d$  for 8 GeV  $\lesssim M_{A_1} \lesssim 10.1$  GeV, including effects from  $\eta_b(nS) - A_1$  mixing, have recently been investigated in [33], implying  $X_d \lesssim 2 \dots 7$  depending on  $M_{A_1}$ . (These bounds are consistent with limits from the violation of lepton universality in inclusive  $\Upsilon(nS)$  decays as proposed in [20, 34–36] and studied in [37, 38].)

Possible  $\eta_b(nS) - A_1$  mixings would also affect the ALEPH bounds on  $H \rightarrow 2 A_1 \rightarrow 4\tau$  [12] in the interesting mass range 9 GeV  $\lesssim M_{A_1} \lesssim 10.5$  GeV, since  $A_1$  decaying hadronically through its  $\eta_b$  components would imply a different signature. The corresponding consequences for this process have not been taken into account before; this study is the purpose of the present paper. In fact, our result is quite dramatic: the ALEPH bounds imply practically no constraint on the  $BR(H \rightarrow 2 A_1)$  in the corresponding mass range, since the  $BR(A_1 \rightarrow \tau^+ \tau^-)$  tends to be very small even for small values of  $X_d$ . The origin of this phenomenon can easily be understood qualitatively: the width of the decay  $A_1 \rightarrow \tau^+ \tau^-$  of the pure state  $A_1$ , albeit proportional to  $X_d^2$  (which appears also in the coupling of  $A_1$  to  $\tau$  leptons), is always much smaller than the hadronic width of the  $\eta_b(nS)$  to hadrons given the present upper bounds on  $X_d$ . Hence, even a small admixture of  $\eta_b(nS)$  to any physical eigenstate implies a large hadronic decay width, suppressing the branching ratio of the physical state into  $\tau^+ \tau^-$  and making it very difficult to detect. For  $X_d \lesssim 10$  this effect is approximately independent from  $X_d$ , since both the width for  $A_1 \rightarrow \tau^+ \tau^-$  and the  $\eta_b(nS) - A_1$  mixing are proportional to  $X_d^2$ .

In the next section we study this phenomenon quantitatively, with the result stated above. In section 3 we briefly comment on the impact of our result on future Higgs searches.

## 2 The $BR(H \rightarrow 4\tau)$ in the presence of $A - \eta_b$ mixing

In this section we consider the mixing of a CP odd Higgs state  $A_1$  (denoted by  $A$  for simplicity) with the states  $\eta_b(1S)$ ,  $\eta_b(2S)$  and  $\eta_b(3S)$  with masses below the  $B\bar{B}$  threshold. The mass squared matrix in the basis  $\eta_b(1S) - \eta_b(2S) - \eta_b(3S) - A$  can be written as [22]

$$\mathcal{M}^2 = \begin{pmatrix} m_{\eta_b(1S)}^2 & 0 & 0 & \delta m_1^2 \\ 0 & m_{\eta_b(2S)}^2 & 0 & \delta m_2^2 \\ 0 & 0 & m_{\eta_b(3S)}^2 & \delta m_3^2 \\ \delta m_1^2 & \delta m_2^2 & \delta m_3^2 & M_A^2 \end{pmatrix}. \quad (2.1)$$

The off-diagonal elements  $\delta m_n^2$  depend on the  $\eta_b(nS)$  wave functions at the origin, and  $X_d$  as given in (1.1) multiplied by the coupling of a SM like Higgs boson to  $b$  quarks [19–22]. Estimating the wave functions at the origin as in [20–22] one obtains

$$\begin{aligned}\delta m_1^2 &\simeq 0.14 \text{ GeV}^2 \times X_d, \\ \delta m_2^2 &\simeq 0.11 \text{ GeV}^2 \times X_d, \\ \delta m_3^2 &\simeq 0.10 \text{ GeV}^2 \times X_d.\end{aligned}\tag{2.2}$$

The errors on these quantities are about 10%, but our subsequent results are not sensitive to the precise numerical values. For the diagonal elements  $m_{\eta_b(nS)}^2$  we take [25]  $m_{\eta_b(2S)} = 10002 \text{ MeV}$ ,  $m_{\eta_b(3S)} = 10343 \text{ MeV}$ .  $m_{\eta_b(1S)}^2$  is determined, for given  $M_A$  and  $X_d$ , by the condition that the state with its mass of  $\sim 9391 \text{ MeV}$  observed in radiative  $\Upsilon(3S)$  and  $\Upsilon(2S)$  decays by BABAR [23, 24] must be identified with one of the eigenstates of  $\mathcal{M}^2$ . Again, our subsequent results depend only weakly on these masses.

It is straightforward to diagonalize the mass matrix (2.1). The 4 eigenstates will be denoted by  $\eta_i$ , which are decomposed into the unmixed states as

$$\eta_i = P_{i,1} \eta_b(1S) + P_{i,2} \eta_b(2S) + P_{i,3} \eta_b(3S) + P_{i,4} A.\tag{2.3}$$

Both the eigenvalues of the mass matrix (2.1) and the mixing coefficients  $P_{i,j}$  in (2.3) depend on the unknown mass  $M_A$ . Let us recall some obvious properties of the eigenvalues and the mixing coefficients: whenever  $M_A$  is far from any of the  $m_{\eta_b(nS)}$ , the mixing will be relatively small (but increasing with  $X_d$ ), and  $A$  will be an approximate mass eigenstate. For fixed  $X_d$ , the closer  $M_A$  is to  $m_{\eta_b(nS)}$ , the larger the  $A - \eta_b(nS)$  mixing will be, resulting in shifts of the eigenvalues of  $\mathcal{M}^2$  w.r.t. its diagonal elements.

We recall that the state with a mass of  $\sim 9391 \text{ MeV}$  observed by BABAR must be identified with one of the eigenstates of  $\mathcal{M}^2$ . Independently from the value of the diagonal element  $m_{\eta_b(1S)}$  of  $\mathcal{M}^2$ , it follows that  $M_A$  cannot be arbitrarily close to  $9391 \text{ MeV}$  unless the mixing (and hence  $X_d$ ) tends to zero. This consideration leads to an upper bound on  $X_d$  depending on  $M_A$ , with  $X_d \rightarrow 0$  for  $M_A \rightarrow 9391 \text{ MeV}$ , and still  $X_d \lesssim 20$  for  $M_A \sim 10 \text{ GeV}$  or  $M_A \sim 8.5 \text{ GeV}$  [21].

Next we turn to the decays of the eigenstates  $\eta_i$ , starting with the decays of the states before mixing.  $A$  will decay dominantly into  $A \rightarrow \tau^+ \tau^-$ , with a partial width  $\Gamma_A^{\tau\tau}$  given by

$$\Gamma_A^{\tau\tau} = X_d^2 \frac{G_F m_\tau^2 M_A}{4\sqrt{2}\pi} \sqrt{1 - 4 \frac{m_\tau^2}{M_A^2}} \sim X_d^2 \times 1.9 \cdot 10^{-2} \text{ MeV} \times \left( \frac{M_A}{10 \text{ GeV}} \right).\tag{2.4}$$

We determine the  $BR(A \rightarrow \tau^+ \tau^-)$  from NMHDECAY [39, 40] inside NMSSMTools [41] (assuming  $\tan \beta \sim 5$ ), which gives  $BR(A \rightarrow \tau^+ \tau^-) \sim 0.9 - 0.75$  with increasing  $M_A$ , the remaining  $BR$  originating from  $A$  decays into  $c\bar{c}$  quarks and gluons. (A smaller  $BR(A \rightarrow \tau^+ \tau^-)$ , as advocated for some parameter choices in [18], would only amplify our subsequent conclusions.) Hence we take  $\Gamma_A^{\text{tot}} \sim (1.1 - 1.33) \times \Gamma_A^{\tau\tau}$ .

The states  $\eta_b(nS)$  (before mixing) would decay nearly exclusively into hadrons (like the states  $\eta_c(nS)$ ). Using the formalism in [42], the widths of the states  $\eta_b(nS)$  can be estimated

from the widths of the corresponding  $\Upsilon$  states and the  $\eta_b(nS)$  masses. Subsequently we take  $\Gamma_{\eta_b(1S)} = 11.8 \text{ MeV}$ ,  $\Gamma_{\eta_b(2S)} = 5.4 \text{ MeV}$  and  $\Gamma_{\eta_b(3S)} = 3.9 \text{ MeV}$ . Note that, unless  $X_d \gtrsim 10$ , these widths are much larger than  $\Gamma_A^{\tau\tau}$ . (We recall that, for  $M_A \lesssim 10.1 \text{ GeV}$ ,  $X_d \lesssim 2 \dots 7$  due to constraints from  $\Upsilon(nS) \rightarrow \gamma A_1$ ,  $A_1 \rightarrow 2 \text{ leptons}$  [33].)

In terms of these widths and the mixing coefficients, the  $BR(\eta_i \rightarrow \tau^+ \tau^-)$  of the eigenstates  $\eta_i$  are given by [22]

$$BR(\eta_i \rightarrow \tau^+ \tau^-) = \frac{P_{i,4}^2 \Gamma_A^{\tau\tau}}{\left( \sum_{n=1}^3 P_{i,n}^2 \Gamma_{\eta_b(nS)} \right) + P_{i,4}^2 \Gamma_A^{\text{tot}}} . \quad (2.5)$$

Let us consider the state  $\eta_i$  with the largest  $A$  component, i.e. the largest coefficient  $P_{i,4}^2$ . (Since, essentially,  $A$  mixes with just one of the  $\eta_b(nS)$  states depending on  $M_A$ , there exists always one state with  $P_{i,4}^2 \gtrsim 0.5$ .) Its  $BR(\eta_i \rightarrow \tau^+ \tau^-)$  is smaller than  $0.9 - 0.75$  due to the terms  $\sim \Gamma_{\eta_b(nS)}$  in the denominator of (2.5). In fact, even if  $P_{i,n}^2 \ll 1$ , these terms are often numerically dominant due to  $\Gamma_{\eta_b(nS)} \gg \Gamma_A^{\text{tot}}$ , implying a considerable reduction of the  $BR(\eta_i \rightarrow \tau^+ \tau^-)$ . For  $X_d \lesssim 5$ , the result is nearly independent from  $X_d$ , since  $\Gamma_A^{\tau\tau}$  and  $\Gamma_A^{\text{tot}}$  as well as  $P_{i,n}^2$  are proportional to  $X_d^2$ , and  $X_d^2$  cancels out.

In figure 1 we show the  $BR(\eta_i \rightarrow \tau^+ \tau^-)$  for the state  $\eta_i$  with the largest  $A$  component as function of  $M_A$  for  $X_d = 1$ . Depending on  $M_A$ , this state corresponds to  $\eta_1 \dots \eta_4$ , which is indicated by the various colors. For  $X_d = 1$ , the mass of this state is practically identical to  $M_A$ . Usually, the branching ratios into  $\tau^+ \tau^-$  of the remaining states are negligibly small. Note that, whenever  $M_A$  is close to any of the masses  $m_{\eta_b(nS)}$ , the mixing becomes strong ( $P_{i,4}^2 \sim P_{i,n}^2 \sim 1/2$ ) leading to  $BR(\eta_i \rightarrow \tau^+ \tau^-) \sim \Gamma_A^{\tau\tau} / \Gamma_{\eta_b(nS)}$ , which is very small. (As stated above, we must have  $X_d \rightarrow 0$  for  $M_A \rightarrow 9391 \text{ MeV}$ . This upper bound is applied to  $X_d$  for  $M_A \sim 9391 \text{ MeV}$  in figure 1, but  $X_d = 1$  is used for all other values of  $M_A$ .) Remarkably, even if  $M_A$  is not close to any of the masses  $m_{\eta_b(nS)}$ , the suppression of the  $BR(\eta_i \rightarrow \tau^+ \tau^-)$  is still quite strong due to the terms  $\sim \Gamma_{\eta_b(nS)}$  in the denominator of (2.5), and  $BR(\eta_i \rightarrow \tau^+ \tau^-) \lesssim 0.65$  for any  $M_A$  in the range  $9 - 10.5 \text{ GeV}$ .

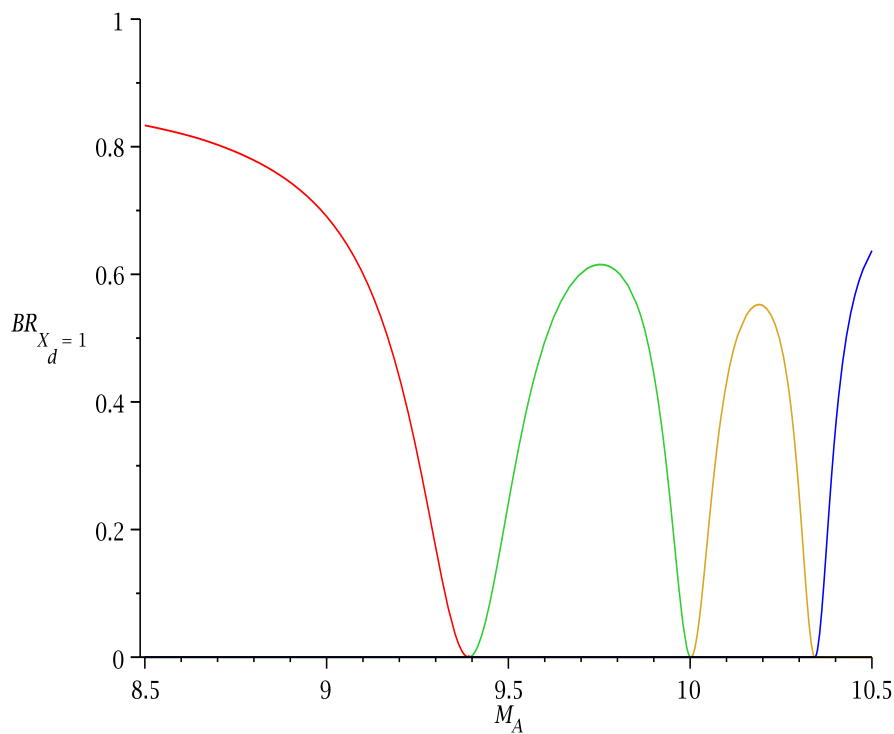
Finally we have to re-interpret the decay  $H \rightarrow AA \rightarrow 4\tau$  in the presence of  $A - \eta_b(nS)$  mixing: now this process corresponds to  $\sum_{i,j=1}^4 (H \rightarrow \eta_i \eta_j \rightarrow 4\tau)$ . The coupling of the states  $\eta_i$  to  $H$  (originating from the coupling of  $A$  to  $H$ ) is proportional to  $P_{i,4}$ , and we can write

$$\sum_{i,j=1}^4 BR(H \rightarrow \eta_i \eta_j \rightarrow 4\tau) = BR(H \rightarrow AA) \times R ,$$

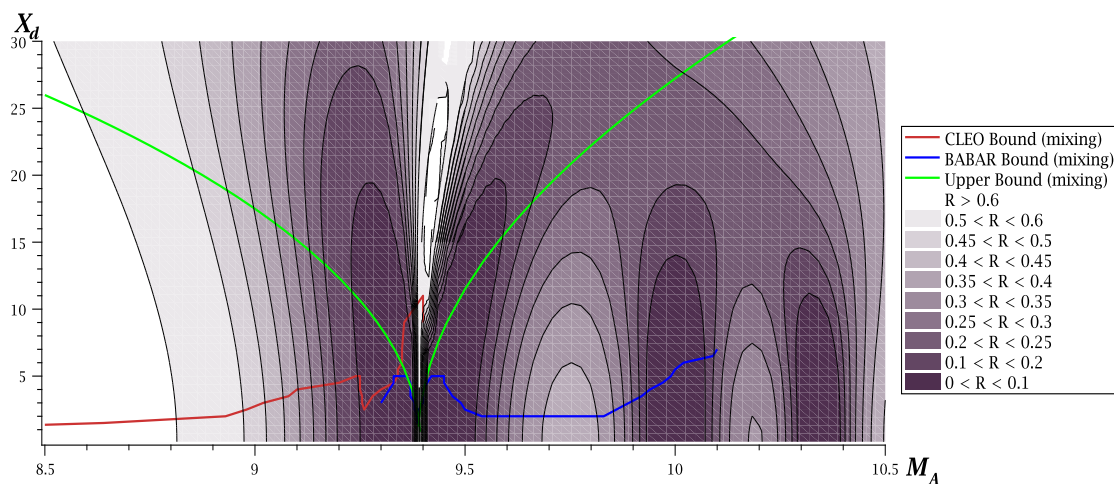
$$R = \left[ \sum_{i=1}^4 P_{i,4}^2 \times BR(\eta_i \rightarrow \tau^+ \tau^-) \right]^2 . \quad (2.6)$$

We can compute  $R$  as function of  $M_A$  and  $X_d$ , and the result is shown in figure 2.

In figure 2 we also show upper bounds on  $X_d$  from CLEO (red), BABAR (blue) and from the condition that one eigenstate of the mass matrix (2.1) has a mass of  $9391 \text{ MeV}$  and  $m_{\eta_b(1S)}$  (before mixing) is within a range  $9360 - 9445 \text{ MeV}$  covered by QCD predictions



**Figure 1.** The  $BR(\eta_i \rightarrow \tau^+ \tau^-)$  for the state  $\eta_i$  with the largest  $A$  component as function of  $M_A$  for  $X_d = 1$ . The colors indicate which state  $\eta_i$  is concerned (red  $\rightarrow \eta_1$ , green  $\rightarrow \eta_2$ , brown  $\rightarrow \eta_3$ , blue  $\rightarrow \eta_4$ ).



**Figure 2.** The function  $R$ , defined in (2.6), in the plane  $X_d$  vs.  $M_A$ . Also indicated are upper bounds on  $X_d$  from CLEO (red), BABAR (blue) and from the condition that one eigenstate of the mass matrix (2.1) has a mass of 9391 MeV with  $m_{\eta_b(1S)}$  within a reasonable range (green).

(green). Hence, for  $M_A \lesssim 10.1$  GeV, only small values of  $X_d$ , where  $R$  is nearly independent from  $X_d$  (as explained above), are of interest. Like the  $BR(\eta_i \rightarrow \tau^+ \tau^-)$ ,  $R$  varies strongly with  $M_A$ . It follows from  $R \sim \sum_{i=1}^4 P_{i,A}^4 \times BR(\eta_i \rightarrow \tau^+ \tau^-)^2$  that  $R$  never exceeds 0.4 for  $M_A$  in the range 9 – 10.5 GeV, and  $R \sim (\Gamma_A^{\tau\tau} / \Gamma_{\eta_b(nS)})^2$  (which is tiny) as soon as  $M_A$  is near any of the masses  $m_{\eta_b(nS)}$ . Now the quantity  $\xi^2$  constrained by ALEPH (see figure 6 in [12]) must be interpreted as  $\xi^2 = \xi'^2 \times R$ ,  $\xi'^2 = \frac{\sigma(e^+e^- \rightarrow ZH)}{\sigma_{\text{SM}}(e^+e^- \rightarrow ZH)} \times BR(H \rightarrow 2A)$ . It follows that  $\xi'^2$  is left unconstrained at least for  $M_H \gtrsim 98$  GeV and  $M_A$  in the range 9 – 10.5 GeV, as well as for any lower value of  $M_H$  as long as  $M_A$  is in the range where  $R$  in figure 2 is below 0.2, corresponding essentially to a  $BR(\eta_i \rightarrow \tau^+ \tau^-)$  in figure 1 below  $\sim 0.5$  (but depending slightly on  $X_d$ ). Since, in addition, one always has  $\xi'^2 \lesssim 1$  even if the process  $H \rightarrow 2A$  is kinematically allowed (since the  $BR(H \rightarrow b\bar{b})$  is never exactly zero), scenarios with  $M_H \lesssim 98$  GeV are consistent with the ALEPH constraints as well for most values of  $M_A$  in the range 9 – 10.5 GeV.

### 3 Conclusions and outlook

After the publication of the ALEPH analysis [12] it seemed that the attractive scenario with a light CP-even Higgs boson  $H$  and a mass  $M_H$  well below 114 GeV, decaying dominantly as  $H \rightarrow 2A \rightarrow 4\tau$ , was tightly constrained. We have shown that these constraints are absent for  $M_H \gtrsim 98$  GeV and  $M_A$  in the range 9 – 10.5 GeV, and in the case of lower values of  $M_H$  for most values of  $M_A$  in this range. The origin is a reduced  $BR(A \rightarrow \tau^+ \tau^-)$  caused by  $A - \eta_b(nS)$  mixing, leading to dominant hadronic decays of the physical eigenstates. This window for  $M_A$  is of particular interest, since it contains the region in which the tension between the observed  $\eta_b(1S)$  mass and its prediction based on QCD can be resolved [22, 33] through this mixing. We emphasize that we did not make particular assumptions on the SU(2) doublet component  $\cos\theta_A$ , on  $\tan\beta$  or on the coupling  $X_d$  (see (1.1)) of  $A$  to  $b$  quarks since, at least for small mixing angles,  $X_d^2$  cancels out in the expression (2.5) for the  $BR(\eta_i \rightarrow \tau^+ \tau^-)$  for the mass eigenstates.

For small  $X_d$  and a correspondingly small  $A - \eta_b(nS)$  mixing, this result seems counter-intuitive at first sight. However, the point is that already a small admixture of any  $\eta_b(nS)$  state to the mass eigenstate  $\eta_i$  suffices such that the mass eigenstate  $\eta_i$  decays dominantly hadronically, since the corresponding hadronic widths of  $\eta_b(nS)$  are much larger than  $\Gamma_A^{\tau\tau}$ . This remains true for small  $X_d$ , since then  $\Gamma_A^{\tau\tau}$  becomes small as well.

The consequences of this scenario for Higgs searches at the LHC would be quite dramatic, since the dominant Higgs decay mode would be  $H \rightarrow 2A \rightarrow$  hadrons and, like in the scenarios discussed in [43–45], the  $H$  signal would be buried under the QCD background. Moreover, dominant hadronic decays of the mass eigenstate  $\eta_i$  would also handicap searches for  $A$  via central exclusive production [46] at hadron colliders, or via the  $\mu^+ \mu^-$  final state as proposed in [47] and studied, using early LHC data, in [48]. It remains to look for  $A$  in radiative  $\Upsilon$  decays, but corresponding searches have also to be interpreted carefully taking mixing effects into account [19–21, 33].



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