PHYSICAL REVIEW D **93,** 055014 (2016)

Probe of new light Higgs bosons from bottomonium χ_{b0} decay

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(Received 23 October 2015; published 8 March 2016)

We calculate the branching ratios of $\chi_{b0} \to \tau^+\tau^-$ via an s-channel Higgs boson and estimate the sensitivity to this process from $\Upsilon \to \gamma\chi_{b0} \to \gamma\tau^+\tau^-$. We show that future running at the $\Upsilon(3S)$ at a very high luminosity super B factory can put significant constraints on the type-II two-Higgs-doublet model when the discovered 125 GeV Higgs boson is the heavier of the two charge parity (CP)-even scalars.

DOI: 10.1103/PhysRevD.93.055014

I. INTRODUCTION

The search for evidence for new physics beyond the standard model (SM) has been intensifying with the discovery of a scalar boson whose properties are consistent with those expected for the SM Higgs boson. Many extensions of the SM include an extended Higgs sector and there has been considerable experimental effort put into searching for evidence of this. One possibility to constrain such extensions is through the effects of off-shell Higgs bosons in bottom meson decays to third-generation leptons. The decay of $B^+ \to \tau^+\nu$ [1] (or $B \to \tau\nu X$ [2–4]) is sensitive to s-channel exchange of charged Higgs bosons, and $\eta_b \to \tau^+\tau^-$ [5] is sensitive to s-channel exchange of charge parity (CP)-odd neutral Higgs bosons, leading to constraints on parts of parameter space where the boson in question is light and its couplings to down-type fermions are enhanced, such as what occurs in the type-II two-Higgs-doublet model (2HDM) [6].

In this paper we consider the decay of scalar bottomonium $\chi_{b0} \to \tau^+\tau^-$, which is sensitive to s-channel exchange of CP-even neutral Higgs bosons. This process was first proposed as a probe of Higgs bosons in Ref. [7]. The advent of the very high luminosity SuperKEKB e^+e^- collider [8] offers the possibility of using $\chi_{b0} \to \tau^+\tau^-$ to put meaningful constraints on the parameter space of the type-II 2HDM. To explore this possibility we first estimate the χ_{b0} decay constant in Sec. II, and use it to calculate the relevant branching ratios in Sec. III. We find that the process $\Upsilon(3S) \to \gamma\chi_{b0}(2P) \to \gamma\tau^+\tau^-$ is the most promising, though the event rate from SM Higgs exchange is a few orders of magnitude too small to be observed. In Sec. IV we consider the prospects in the type-II 2HDM, in which the scalar couplings to both the $b\bar{b}$ initial state and the $\tau^+\tau^-$ final state can be enhanced. We show that, with 250 fb^{-1} of data on the $\Upsilon(3S)$, this process has the potential to constrain a large region of as-yet-unexcluded parameter space in which the second CP-even Higgs boson is lighter than about 80 GeV. This amount of integrated luminosity represents only a small fraction of the expected ultimate SuperKEKB integrated luminosity of 50 ab^{-1} [8], and could be collected during early running as the accelerator luminosity is ramped up. We briefly conclude in Sec. V.

II. CALCULATION OF THE χ_{b0} DECAY CONSTANT

We use the quark model to calculate the χ_{b0} decay constant (f_{χ_0}), in particular, the *mock meson* approach [9–12] which has been a useful tool in calculating hadronic matrix elements (see also Ref. [13]). The basic premise of the quark model is that hadrons are made of constituent quarks and antiquarks and one solves a Schrodinger (like) equation, typically employing a short distance Lorentz vector one-gluon-exchange interaction and a Lorentz scalar confining interaction, to obtain hadron masses and wave functions which are used to calculate hadron properties (see for example Refs. [14,15]). Quark model predictions have been reasonably reliable in describing the properties of known mesons. Ideally one would like to use *ab initio* lattice QCD calculations to calculate hadron properties but we know of no lattice calculation of f_{χ_0} so we turn to the quark model. We note that quark model predictions [15] give similar agreement to lattice calculations [16] for the related Υ leptonic decays. There are a few predictions of scalar decay constants using the QCD sum-rules approach [17–19] but it is not clear how to relate their conventions to the ones used in this paper making comparisons difficult.

The basic assumption of the mock meson approach is that physical hadronic amplitudes can be identified with the corresponding quark model amplitudes in the weak binding limit of the valence quark approximation. This correspondence is exact in the limit of zero binding and in the hadron rest frame. Away from this limit the amplitudes are not in general Lorentz invariant by terms of order p_i^2/m_i^2. In this approach the mock meson, $\widetilde{M}$, is defined as a state of a free quark and antiquark with the wave function of the physical meson, M:

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$$|\widetilde{M}(\vec{K})\rangle = \sqrt{2\widetilde{M}_M}\int d^3p\Phi_M(\vec{p})\chi_{s\bar{s}}\phi_{q\bar{q}} \times |q[(m_q/\mu)\vec{K}+\vec{p},s]\bar{q}[(m_{\bar{q}}/\mu)\vec{K}-\vec{p},\bar{s}]\rangle, \quad (1)$$

where $\Phi_M(\vec{p})$, $\chi_{s\bar{s}}$ and $\phi_{q\bar{q}}$ are momentum, spin and flavor wave functions respectively, $\mu = m_q + m_{\bar{q}}$, $\widetilde{M}_M$ is the mock meson mass, and $\vec{K}$ is the mock meson momentum. To calculate the hadronic amplitude, the physical matrix element is expressed in terms of Lorentz covariants with Lorentz scalar coefficients A. In the simple cases when the mock meson matrix element has the same form as the physical meson amplitude we simply take $A = \tilde{A}$.

We write the scalar decay constant for the χ_0 meson as

$$\langle 0|\bar{q}q|M(\vec{K})\rangle = if_{\chi_0}, \quad (2)$$

where $M \equiv \chi_0$ and the amplitude has been normalized to one particle per unit volume [13]. To calculate the left-hand side of Eq. (2) we first calculate

$$\langle 0|\bar{q}q|q[(m_q/\mu)\vec{K}+\vec{p},s]\bar{q}[(m_{\bar{q}}/\mu)\vec{K}-\vec{p},\bar{s}]\rangle \quad (3)$$

using free quark and antiquark wave functions, then weight the resulting expression with the meson's momentum space wave function using Eq. (1).

There are typically a number of ambiguities in this approach that we must deal with. For example, there are several different prescriptions for the definition of the meson mass $\widetilde{M}_M$ appearing in Eq. (1). We use the physical mass as $\widetilde{M}_M$ was introduced to give the correct relativistic normalization of the meson's wave function and hence to give the correct kinematics for the process being studied. Fortunately, many of these ambiguities do not show up in calculations of scalar decay constants, and for the heavy b quark the nonrelativistic limit is a good approximation. Discussion of these ambiguities is given in Refs. [10–12].

Evaluating Eq. (2) in the nonrelativistic limit we obtain

$$f_{\chi_0} = -\frac{3\sqrt{3M_{\chi_0}}}{\sqrt{\pi}\tilde{m}_q}R'(0), \quad (4)$$

where $R'(0)$ is the derivative of the radial part of the χ_0 wave function at the origin, M_{χ_0} is the measured χ_0 mass, and $\tilde{m}_q$ is the heavy quark constituent mass.

The radial wave functions are computed using the relativized quark model [14]. For $\chi_{b0}(1P)$ and $\chi_{b0}(2P)$ these wave functions were recently computed in Ref. [15], which found $R'_{\chi_{b0}(1P)}(0) = 2.255\ \text{GeV}^{5/2}$ and $R'_{\chi_{b0}(2P)}(0) = 2.290\ \text{GeV}^{5/2}$. For $\chi_{c0}(1P)$ we compute the wave function following the same procedure and obtain $R'_{\chi_{c0}(1P)}(0) = 0.912\ \text{GeV}^{5/2}$.

Inserting these values into Eq. (4) and using constituent quark masses $\tilde{m}_b = 4.977$ GeV and $\tilde{m}_c = 1.628$ GeV [14], we obtain

$$f_{\chi_{c0}(1P)} = -3.03\ \text{GeV}^2,$$
$$f_{\chi_{b0}(1P)} = -4.17\ \text{GeV}^2,$$
$$f_{\chi_{b0}(2P)} = -4.31\ \text{GeV}^2. \quad (5)$$

III. $\chi_0 \to \ell^+\ell^-$ DECAY AND SM EVENT RATES

Now that we have an estimate for the scalar quarkonium decay constants we can obtain expressions for their decay widths and branching ratios. The matrix element for a scalar meson χ_0 decaying to two leptons $\ell^+\ell^-$ via an s-channel SM Higgs boson is given by

$$\mathcal{M}^H = \langle \ell^+\ell^-|\frac{m_\ell}{v}\bar{\ell}\ell|0\rangle\frac{i}{M_H^2}\langle 0|\frac{m_q}{v}\bar{q}q|\chi_0\rangle = -\left(\frac{m_q m_\ell}{v^2 M_H^2}\right)f_{\chi_0}\bar{u}(p_{\ell^-})v(p_{\ell^+}), \quad (6)$$

where $v^2 = 1/\sqrt{2}G_F$ is the SM Higgs vacuum expectation value, M_H is the Higgs mass, we have neglected $M_{\chi_0}^2$ relative to M_H^2 in the propagator, and in the second line we have used Eq. (2). The partial width is then given by

$$\Gamma^H(\chi_0 \to \ell^+\ell^-) = \frac{M_{\chi_0}}{8\pi}\left[1-\frac{4m_\ell^2}{M_{\chi_0}^2}\right]^{3/2}\left(\frac{m_q m_\ell}{v^2 M_H^2}\right)^2 f_{\chi_0}^2. \quad (7)$$

This expression is in agreement with those given by Refs. [7,20] with the exception that those papers take the current and constituent quark masses to be equal to each other.

There is also a contribution to the $\chi_0 \to \ell^+\ell^-$ decay through a two-photon intermediate state. Following Ref. [5], we estimate the partial width for this one-loop process using the optical theorem,

$$\Gamma^{2\gamma}(\chi_0 \to \ell^+\ell^-) \simeq \frac{\alpha^2}{2\beta_\ell}\left[\frac{m_\ell}{M_{\chi_0}}\ln\frac{(1+\beta_\ell)}{(1-\beta_\ell)}\right]^2\Gamma(\chi_0 \to \gamma\gamma), \quad (8)$$

where α is the electromagnetic fine structure constant,

$$\Gamma(\chi_0 \to \gamma\gamma) = \frac{4\pi\alpha^2}{81M_{\chi_0}^3}f_{\chi_0}^2, \quad \text{and} \quad \beta_\ell = \sqrt{1-\frac{4m_\ell^2}{M_{\chi_0}^2}}. \quad (9)$$

We first consider $\chi_{c0}(1P) \to \mu^+\mu^-$ (the decay to $\tau^+\tau^-$ is kinematically forbidden). Taking $M_{\chi_{c0}(1P)} = 3.415$ GeV, $m_c = 1.27$ GeV, $m_\mu = 0.10566$ GeV, and $M_H = 125$ GeV [21], we obtain $\Gamma^H(\chi_{c0}(1P) \to \mu^+\mu^-) = 2.5 \times 10^{-20}$ GeV. Combining with the total width $\Gamma^{\text{tot}}_{\chi_{c0}(1P)} = 10.5 \pm 0.6$ MeV

[21], we obtain the branching ratio from SM Higgs exchange,

$$\mathrm{BR}^H(\chi_{c0}(1P) \to \mu^+\mu^-) = 2.4 \times 10^{-18}, \qquad (10)$$

which is clearly too small to be observed due to the small values of the muon and charm-quark masses. The competing two-photon intermediate state in fact yields a much larger contribution,

$$\mathrm{BR}^{2\gamma}(\chi_{c0}(1P) \to \mu^+\mu^-) \simeq 2 \times 10^{-10}. \qquad (11)$$

The situation is not quite so dire for the χ_{b0} decays. Taking $M_{\chi_{b0}(1P)} = 9.860$ GeV, $M_{\chi_{b0}(2P)} = 10.232$ GeV, $m_b = 4.67$ GeV, and $m_\tau = 1.77682$ GeV [21], we obtain the SM Higgs-exchange contributions, $\Gamma^H(\chi_{b0}(1P) \to \tau^+\tau^-) = 4.3 \times 10^{-16}$ GeV and $\Gamma^H(\chi_{b0}(2P) \to \tau^+\tau^-) = 4.8 \times 10^{-16}$ GeV.

To estimate the branching ratios we need the total widths for the $\chi_{b0}(1P)$ and $\chi_{b0}(2P)$, which have not been measured. To estimate them we follow Ref. [15], which combined the measured branching ratios for $\chi_{b0} \to \gamma\Upsilon(1S)$ with the predicted partial widths for these transitions, yielding $\Gamma^{\mathrm{tot}}_{\chi_{b0}(1P)} = 1.35$ MeV and $\Gamma^{\mathrm{tot}}_{\chi_{b0}(2P)} = (247 \pm 93)$ keV [15]. Note that there is modeling dependence in these estimates as we have relied on the results of the relativized quark model to obtain the radiative transition widths. Furthermore, the experimental values for the $\chi_{b0} \to \gamma\Upsilon(1S)$ branching ratios have experimental errors, which are particularly large for the $\chi_{b0}(2P)$ state. Combining our $\chi_{b0} \to \tau^+\tau^-$ partial width calculation with these total width estimates we obtain the branching ratios from SM Higgs exchange,

$$\mathrm{BR}^H(\chi_{b0}(1P) \to \tau^+\tau^-) = 3.1 \times 10^{-13},$$
$$\mathrm{BR}^H(\chi_{b0}(2P) \to \tau^+\tau^-) = (1.9 \pm 0.5) \times 10^{-12}. \qquad (12)$$

The competing two-photon intermediate state again yields a larger contribution,

$$\mathrm{BR}^{2\gamma}(\chi_{b0}(1P) \to \tau^+\tau^-) \simeq 1 \times 10^{-9},$$
$$\mathrm{BR}^{2\gamma}(\chi_{b0}(2P) \to \tau^+\tau^-) \simeq 6 \times 10^{-9}. \qquad (13)$$

The final step is to estimate event rates to see if these processes are actually measurable. We base our estimates on the production of the $\Upsilon(2S)$ and $\Upsilon(3S)$ at the high luminosity SuperKEKB e^+e^- collider followed by a radiative decay to the χ_{b0} states. The $e^+e^- \to \Upsilon(2S)$ cross section averaged over the Belle and *BABAR* measurements is about 6.5 nb [22]. We assume $\mathcal{L} = 250$ fb^{-1} of integrated luminosity which will yield 1.6×10^9 $\Upsilon(2S)$s. The branching ratio for $\Upsilon(2S) \to \gamma\chi_{b0}(1P)$ is $(3.8 \pm 0.4)\%$ [21], which would yield 6.2×10^7 $\chi_{b0}(1P)$s. Combining this with the branching ratios in Eqs. (12)–(13) yields 1.9×10^{-5} $\gamma\tau^+\tau^-$ signal events from SM Higgs exchange and about 0.07 events from the two-photon intermediate state.

Likewise we can estimate the number of χ_{b0}s that would be produced from $\Upsilon(3S)$ decay. The $\Upsilon(3S)$ e^+e^- production cross section is 4 nb [22], and with $\mathcal{L} = 250$ fb^{-1} of integrated luminosity this yields 10^9 $\Upsilon(3S)$s. The branching ratio for $\Upsilon(3S) \to \gamma\chi_{b0}(2P)$ is $(5.9 \pm 0.6)\%$ and for $\Upsilon(3S) \to \gamma\chi_{b0}(1P)$ is $(0.27 \pm 0.04)\%$ [21], which yield 5.9×10^7 $\chi_{b0}(2P)$s and 2.7×10^6 $\chi_{b0}(1P)$s. Combining these with the branching ratios in Eqs. (12)–(13) yields 1.1×10^{-4} and 8.5×10^{-7} $\gamma\tau^+\tau^-$ signal events via an s-channel SM Higgs boson, for the decays via $\chi_{b0}(2P)$ and $\chi_{b0}(1P)$ respectively. The number of signal events from the two-photon intermediate state is about 0.3 and 0.003 for the decays via $\chi_{b0}(2P)$ and $\chi_{b0}(1P)$ respectively.

Clearly these event numbers are too small to be able to observe $\chi_{b0} \to \tau^+\tau^-$ decays mediated by the SM Higgs boson. However, if the fermion-Higgs couplings were enhanced these decays might become observable. We explore this possibility in the following section.

IV. SIGNAL RATES IN THE TWO-HIGGS-DOUBLET MODEL

A. Resonant signal

To explore the possibility of enhanced $\chi_{b0} \to \tau^+\tau^-$ decays, we consider the type-II 2HDM [6]. In this model the scalar couplings to b quarks and τ leptons can be simultaneously enhanced for large values of the parameter $\tan\beta$, which is defined as the ratio of vacuum expectation values of the two Higgs doublets. The model contains two CP-even neutral scalars, which we call H_{125} and H_{new}. We identify H_{125} with the discovered Higgs boson at 125 GeV. The Higgs-exchange matrix element in Eq. (6) gets modified by the presence of the second Higgs resonance, yielding

$$\left(\frac{m_b m_\tau}{v^2 M_H^2}\right)^2 \to \left[\frac{m_b m_\tau}{v^2}\left(\frac{\kappa_b^{125}\kappa_\tau^{125}}{M_H^2} + \frac{\kappa_b^{\mathrm{new}}\kappa_\tau^{\mathrm{new}}}{M_{\mathrm{new}}^2 - M_{\chi_{b0}}^2}\right)\right]^2, \qquad (14)$$

where M_{new} is the mass of the second scalar H_{new} and the κ factors represent the couplings of the two scalars to b quarks or τ leptons normalized to the corresponding coupling of the SM Higgs boson [23]. We have kept the $p^2 = M_{\chi_{b0}}^2$ dependence in the second diagram because we are interested in low M_{new}.

Setting the couplings of the 125 GeV Higgs boson equal to their SM values (i.e., working in the *alignment limit* [24]), the branching ratios in Eq. (12) are modified by the multiplicative factor

$$\left[1 + \frac{M_H^2}{M_{\mathrm{new}}^2 - M_{\chi_{b0}}^2}\tan^2\beta\right]^2. \qquad (15)$$

The number of signal events grows with increasing $\tan\beta$ and decreasing $M_{\rm new}$. For a large enough enhancement, the $H_{\rm new}$-exchange contribution will dominate over the SM two-photon intermediate state process. As we will see, a detectable signal will require a large number $\gg 1$ of signal events, so that the SM two-photon contribution can be neglected.

B. A continuum signal?

There is also a continuum signal from $\Upsilon \to \gamma H^*_{\rm new} \to \gamma\tau^+\tau^-$, in which the photon is not monoenergetic (we do not consider $M_{\rm new} \lesssim 10$ GeV, which is excluded by searches for $\Upsilon \to \gamma H_{\rm new}$ [25,26]). This can be computed from the on-shell $\Upsilon \to \gamma H_{\rm new}$ decay width by taking the Higgs off shell. In particular, neglecting the SM Higgs contribution we have

$$\Gamma(\Upsilon \to \gamma H^*_{\rm new} \to \gamma\tau^+\tau^-) = \frac{1}{\pi}\int_{4m_\tau^2}^{M_\Upsilon^2} \frac{dQ^2 Q\Gamma(\Upsilon \to \gamma H^*_{\rm new})\Gamma(H^*_{\rm new} \to \tau\tau)}{(Q^2 - M^2_{\rm new})^2 + M^2_{\rm new}\Gamma^2_{\rm new}}, \quad (16)$$

where $\Gamma_{\rm new}$ is the total width of $H_{\rm new}$ and the partial widths in the numerator are computed by setting the $H_{\rm new}$ mass equal to the $\tau^+\tau^-$ invariant mass Q. We use [25][1]

$$\Gamma(\Upsilon \to \gamma H^*_{\rm new}) = \frac{(m_b\kappa_b^{\rm new})^2}{2\pi\alpha v^2}\left[1 - \frac{Q^2}{M^2_\Upsilon}\right]\Gamma(\Upsilon \to \mu^+\mu^-) \quad (17)$$

and

$$\Gamma(H^*_{\rm new} \to \tau^+\tau^-) = \frac{(m_\tau\kappa_\tau^{\rm new})^2 Q}{8\pi v^2}\left[1 - \frac{4m_\tau^2}{Q^2}\right]^{3/2}. \quad (18)$$

By integrating numerically we can compare the number of continuum signal events to the number of resonant signal events through the intermediate χ_{b0}. Taking $M_{\rm new}$ to be well above the Υ mass, the dependence on $M_{\rm new}$ and the coupling enhancement factors $\kappa^{\rm new}_{b,\tau}$ is the same in the continuum and resonant processes, and so drops out in their ratio. We find that, on the $\Upsilon(3S)$ the total continuum signal rate is only about 0.5% of the resonant rates through the $\chi_{b0}(2P)$ and $\chi_{b0}(1P)$, and on the $\Upsilon(2S)$ the total continuum signal rate is only about 3% of the resonant rate through the $\chi_{b0}(1P)$. Furthermore, this small continuum production is spread over a photon energy range of about 6 GeV, compared to the resonant photon peaks with widths of order an MeV or less (see the next subsection). We therefore neglect the continuum signal in what follows.

[1]We correct a misprint in Ref. [25] following Refs. [7,20] (see also Refs. [27–30]).

TABLE I. For the three processes considered we give the tagging photon energies E_γ in the Υ center-of-mass frame, the line width δE_γ of the photon peak, the differential cross section $d\sigma_B/dE_\gamma$ of the continuum $e^+e^- \to \gamma\tau^+\tau^-$ background evaluated at the photon peak, and the number N_B of continuum $e^+e^- \to \gamma\tau^+\tau^-$ events in a window of width $2\delta E_\gamma$ centered at the photon peak in 250 fb^{-1} of integrated luminosity.

Parent	Daughter	E_γ	δE_γ	$d\sigma_B/dE_\gamma$	N_B
$\Upsilon(3S)$	$\chi_{b0}(2P)$	122 MeV	0.24 MeV	36 fb/MeV	4320
$\Upsilon(3S)$	$\chi_{b0}(1P)$	484 MeV	1.3 MeV	8.8 fb/MeV	5720
$\Upsilon(2S)$	$\chi_{b0}(1P)$	163 MeV	1.3 MeV	30 fb/MeV	19500

C. Backgrounds

The resonant signal is a single photon, monoenergetic in the parent Υ rest frame, with the remainder of the collision energy taken up by the $\tau^+\tau^-$ pair. This must be discriminated from the reducible background $\Upsilon \to \gamma\chi_{b0}$ with χ_{b0} decaying to anything other than $\tau^+\tau^-$, as well as from the irreducible continuum background $e^+e^- \to \gamma\tau^+\tau^-$. We assume that the $\tau^+\tau^-$ identification purity will be good enough that the reducible backgrounds can be ignored. We then only have to worry about the irreducible background in a signal window around the characteristic photon energy.

In Table I we give the photon energies E_γ and natural line widths δE_γ (computed from the χ_{b0} total decay widths) for the three signal processes $\Upsilon \to \gamma\chi_{b0} \to \gamma\tau^+\tau^-$ evaluated in the Υ center-of-mass frame. We also give the differential cross section of the continuum $e^+e^- \to \gamma\tau^+\tau^-$ background at the corresponding photon energy, for running at the appropriate Υ resonance energy. The latter was evaluated using MadGraph5_aMC@NLO [31] with a generator-level cut on the photon rapidity of $|\eta_\gamma| < 5$ in the center-of-mass frame.

D. Sensitivity

To make a conservative first estimate of the sensitivity, we take as background the total number of $e^+e^- \to \gamma\tau^+\tau^-$ events with a photon energy within a window of width $2\delta E_\gamma$.[2] The resulting number of background events in the signal window is shown in the last column of Table I. We do not include the τ reconstruction and identification efficiencies. A more sophisticated event selection based on better modeling of the background, for example taking into account the angular distributions and τ polarizations, would improve the sensitivity.

In Figs. 1 and 2 we show the resulting 5σ discovery reach and 95% confidence level (C.L.) exclusion reach from 250 fb^{-1} of data on the $\Upsilon(3S)$ and $\Upsilon(2S)$, respectively. We plot the sensitivity reach as a function of $M_{\rm new}$ and $\tan\beta$,

[2]This choice of the photon energy window provides good signal efficiency while maintaining near-optimal signal significance in the presence of large backgrounds [32].

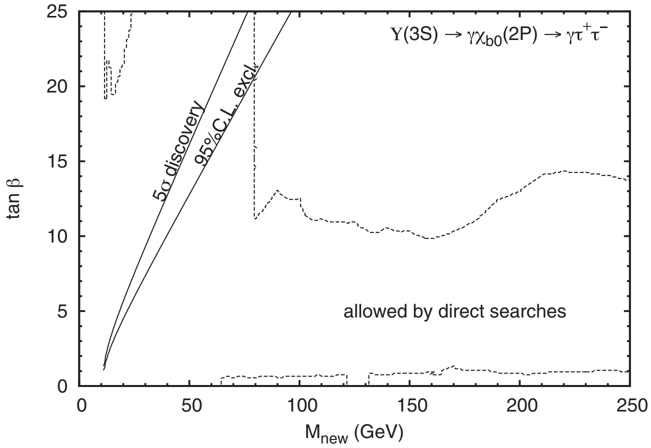

FIG. 1. 5σ discovery and 95% C.L. exclusion reach in the type-II 2HDM from 250 fb^{-1} of data on the $\Upsilon(3S)$. The sensitivity is to the regions to the left of the solid curves. We have set the couplings of the 125 GeV Higgs boson equal to their SM values. The dashed lines indicate the parameter regions still allowed by direct searches for $H_{\rm new}$, computed using HiggsBounds 4.2.0 [33].

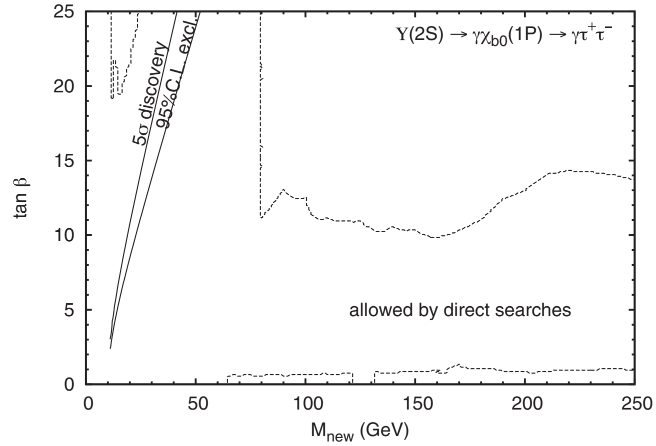

FIG. 2. As in Fig. 1 but for 250 fb^{-1} of data on the $\Upsilon(2S)$.

assuming that the couplings of the 125 GeV Higgs boson take their SM values.[3] We also show, using dotted lines, the parameter region allowed by direct searches, as computed using HiggsBounds 4.2.0 [33].[4] $H_{\rm new}$ masses below about 10 GeV are generally excluded by searches for $\Upsilon \to \gamma H_{\rm new}$ [25,26], which have not been included in HiggsBounds. We note that the 95% C.L. exclusion line for the $\Upsilon(3S)$-initiated process corresponds to 130 signal events on top of about 4300 background events, so that more sophisticated kinematic cuts could improve signal to background substantially and even more so for the 5σ discovery curves.

On the $\Upsilon(3S)$, the sensitivity comes almost entirely from decays to $\chi_{b0}(2P)$; the signal rate from $\chi_{b0}(1P)$ is more than one hundred times smaller but with comparable background. This process has the potential to probe a large region of the type-II 2HDM parameter space with $H_{\rm new}$ masses below 80 GeV and moderate to large $\tan\beta$ that is currently unconstrained by existing searches. On the $\Upsilon(2S)$ the sensitivity is not as good due to a combination of lower signal rate and a larger photon line width, resulting in more background. Nevertheless, this process can still probe a significant unexcluded region of type-II 2HDM parameter space for $H_{\rm new}$ masses up to 50 GeV at $\tan\beta = 25$.

[3]We further assume that the branching ratio for $H_{125} \to H_{\rm new}H_{\rm new}$ remains small in order to avoid constraints from modifications of the 125 GeV Higgs boson signal strengths.

[4]In Figs. 1–2 the large excluded region at large $\tan\beta$ and $M_{\rm new} > 80$ GeV is from a CMS $pp \to \phi \to \tau\tau$ search using 7 and 8 TeV data [34]. The small excluded region at large $\tan\beta$ and $M_{\rm new} \sim 10$–25 GeV is from a DELPHI $e^+e^- \to b\bar{b}\phi \to b\bar{b}b\bar{b}$ search for a CP-even ϕ [35]. The exclusion at very low $\tan\beta$ is from an ATLAS scalar diphoton resonance search using 8 TeV data [36].

V. SUMMARY

We have computed the scalar decay constants for $\chi_{c0}(1P)$, $\chi_{b0}(1P)$, and $\chi_{b0}(2P)$, which allowed us to predict the branching ratios for decays of these mesons into $\mu^+\mu^-$ or $\tau^+\tau^-$ via an s-channel Higgs boson. While the expected numbers of events from SM Higgs exchange are orders of magnitude too small to observe, the $\chi_{b0} \to \tau^+\tau^-$ branching ratio can be significantly enhanced in the type-II 2HDM. This leads to the first, to our knowledge, indirect probe of the second neutral CP-even scalar from scalar meson decays.

The most promising channel is $\Upsilon(3S) \to \gamma\chi_{b0}(2P) \to \gamma\tau^+\tau^-$. With 250 fb^{-1} of data collected on the $\Upsilon(3S)$ at SuperKEKB, this process has the potential to constrain a large region of yet-unexcluded type-II 2HDM parameter space in which the second CP-even neutral Higgs boson is lighter than about 80 GeV. A more sophisticated rejection of the continuum background should improve this reach further. Running instead on the $\Upsilon(2S)$ yields fewer signal events and larger backgrounds, but still allows a large region of 2HDM parameter space to be probed. We hope that this analysis provides further physics motivation for running at energies other than the $\Upsilon(4S)$ during the early stages of SuperKEKB data-taking.

ACKNOWLEDGMENTS

This work was supported by the Natural Sciences and Engineering Research Council of Canada.

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