PHYSICAL REVIEW D **84,** 031502(R) (2011)

Radiative transitions from $\Upsilon(5S)$ to molecular bottomonium

M. B. Voloshin

William I. Fine Theoretical Physics Institute, University of Minnesota, Minneapolis, Minnesota 55455, USA and Institute of Theoretical and Experimental Physics, Moscow, 117218, Russia

(Received 16 June 2011; published 29 August 2011)

The heavy quark spin symmetry implies that in addition to the recently observed $Z(10610)$ and $Z(10650)$ molecular resonances with $I^G = 1^+$, there should exist two or four molecular bottomonium-like states with $I^G = 1^-$. Properties of these G-odd states are considered, including their production in the radiative transitions from $\Upsilon(5S)$, by applying the same symmetry to the $\Upsilon(5S)$ resonance and the transition amplitudes. The considered radiative processes can provide a realistic option for observing the yet hypothetical states.

DOI: 10.1103/PhysRevD.84.031502 PACS numbers: 14.40.Pq, 12.39.Hg, 13.25.Gv

The masses and the decay properties of the recently discovered [1] isotriplet resonances $Z_b(10610)$ (Z_b) and $Z_b(10650)$ (Z_b') at the $B^*\bar{B}$ and $B^*\bar{B}^*$ thresholds strongly suggest that these are ‘‘molecular’’ type [2] threshold singularities in the S-wave channels for the corresponding heavy meson-antimeson pair. The properties of such threshold resonances are related by the heavy quark spin symmetry (HQSS) similarly to the relations between the properties of the B and B^* mesons. In particular, it has been argued [2] that the existence of the observed ‘‘twin’’ resonances Z_b and Z_b' resonances and the heavy quark spin symmetry imply an existence of a larger family of ‘‘molecular’’ peaks at the thresholds of the B and B^* meson-antimeson pairs, which family should consist of at least four, but more likely of six isovector resonances. The observed Z_b and Z_b' states have quantum numbers $I^G(J^P) = 1^+(1^+)$. The rest of the isovector states in the family have negative G parity and are denoted here as W_{bJ} with two states having $J = 0$ (W_{b0} and W_{b0}') and one each with $J = 1$ and $J = 2$: W_{b1} and W_{b2}. The W_{bJ} states cannot be produced in single-pion transitions from $\Upsilon(5S)$, but can be produced in hadronic transitions with emission of ρ meson from higher Υ-like bottomonium states with mass above approximately 11.4 GeV. At present, however, no data are available on such bottomonium states, and we have to consider other possibilities for a study of the expected new molecular resonances. In this paper, I discuss the production of the isotopically (and electrically) neutral components of the expected isovector multiplets in radiative transitions from the $\Upsilon(5S)$ resonance: $\Upsilon(5S) \to W_{bJ}\gamma$. Specifically, will be considered the relations, following from HQSS, between the rates of such transitions to all the W_{bJ} resonances. Although a prediction of the absolute rate is currently quite uncertain and is limited to the general expectation that $\Gamma[\Upsilon(5S) \to W_{bJ}\gamma]/\Gamma[\Upsilon(5S) \to Z_b\pi] \sim \alpha$, an understanding of the relative rates could serve as guidance in the searches for the expected molecular resonances.

The relations between the discussed resonances arise due to the known suppression by the inverse of the b quark mass of the strong interaction, depending on the spin of the b quark or antiquark. In particular, in the limit where this interaction is completely turned off, the spin variables of the heavy quark and antiquark are purely ‘‘classificational’’ in the sense that they define the quantum numbers of the states containing the $b\bar{b}$ pair, but the dynamics proceed as if the heavy quark had no spin at all. In particular, in this ‘‘spinless b’’ (SLB) limit, the B mesons behave as heavy particles with spin $1/2$, and an S wave $B\bar{B}$ pair can be in a state with total angular momentum of zero or one, and these states are denoted here as $0^-_{\rm SLB}$ and $1^-_{\rm SLB}$, respectively, where the superscript indicates the parity. Clearly, in the SLB limit, the dynamical properties, most importantly those arising from the interaction of the light components of the B mesons, are determined only by these quantum numbers so that there are only two independent channels for the $B\bar{B}$ meson pairs at a fixed total isospin. In particular, a threshold singularity in one or both channels gives rise to threshold ‘‘molecular’’ resonances for the physical meson-antimeson states. It can be seen [2] that for the case where the singularity at the threshold (a bound or a virtual state as appropriate for an S-wave scattering) exists only in the isovector $0^-_{\rm SLB}$ channel, the family of the physical resonances consists of two pairs of ‘‘twin’’ peaks: the Z_b and Z_b' resonances, and an additional pair of ‘‘twin’’ resonances with the quantum numbers $I^G(J^P) = 1^-(0^+)$: W_{b0} and W_{b0}', at, respectively, the $B\bar{B}$ and the $B^*\bar{B}^*$ thresholds. If, however, a threshold singularity is in the $1^-_{\rm SLB}$ channel, or in both channels, then, besides those four resonances, also arise two additional states: W_{b1} with $I^G(J^P) = 1^-(1^+)$ near the $B^*\bar{B}$ threshold and W_{b2} with $I^G(J^P) = 1^-(2^+)$ near the $B^*\bar{B}^*$ threshold. The resonances W_{b1} and W_{b2} contain the heavy $b\bar{b}$ quark pair in a pure spin-1 (ortho) state, while each of the pairs of ‘‘twin’’ resonances contains a mixture of the ortho and para (spin-0) states of the heavy quark pair. The Z_b and Z_b' resonances are two maximal (45°) mixtures, while in the W_{b0} and W_{b1} resonances the ortho-para mixing is with a 30° angle. Therefore, the HQSS requires that the W_{b1} and W_{b2} have an unsuppressed coupling to channels with orthobottomonium, while the $Z_b(Z_b')$ and $W_{b0}(W_{b0}')$ resonances couple to channels with ortho as well as with para bottomonium, with the relative coefficients in these couplings determined by the coefficients in the ortho-para mixing within these resonances.

It should be mentioned that if only one of the possible SLB channels possesses a near-threshold singularity, then it should be the $0^-_{\rm SLB}$ rather than $1^-_{\rm SLB}$, since it is known [3] from the general properties of the QCD that the energy of the ground state in a pseudoscalar flavor-non-singlet channel is not larger than that in the vector channel.

The described picture of the family of the threshold resonances can be established by considering the composition of the SLB spin states $0^-_{\rm SLB}$ and $1^-_{\rm SLB}$ with the spin states of the $b\bar b$ pair with the total spin 0^-_H and 1^-_H in terms of the physical meson-antimeson pairs containing a B or B^* meson and an antimeson. One can find the appropriate compositions with fixed overall quantum numbers by explicitly reconstructing the eigenstates of the spin-dependent Hamiltonian H_s that lifts the SLB degeneracy between those physical states. The Hamiltonian can be written in terms of the spin operators $\vec s_b$ ($\vec s_{\bar b}$) for the b ($\bar b$) quark and $\vec s_q$ ($\vec s_{\bar q}$) describing the B ($\bar B$) mesons in the SLB limit (the spin of the light (anti)quark):

$$H_s = \mu(\vec s_b \cdot \vec s_{\bar q}) + \mu(\vec s_{\bar b} \cdot \vec s_q) = \frac{\mu}{2}(\vec S_H \cdot \vec S_{\rm SLB}) - \frac{\mu}{2}(\vec \Delta_H \cdot \vec \Delta_{\rm SLB}), \tag{1}$$

where $\vec S_H = \vec s_b + \vec s_{\bar b}$, $\vec S_{\rm SLB} = \vec s_q + \vec s_{\bar q}$, $\vec \Delta_H = \vec s_b - \vec s_{\bar b}$,and $\vec \Delta_{\rm SLB} = \vec s_q - \vec s_{\bar q}$. The first expression in Eq. (1) is the standard phenomenological Hamiltonian for describing the masses of the B and B^* mesons: $M(B) = \bar M - 3\mu/4$, $M(B^*) = \bar M + \mu/4$, with $\bar M$ being the (common) mass of the B and B^* mesons in the SLB limit so that $\mu = M(B^*) - M(B) \approx 46$ MeV. The latter form of the expression in Eq. (1) is convenient for considering the states of the meson-antimeson pairs in terms of the total spin in the SLB limit and of the total spin of the $b\bar b$ pair. The convenience of the latter form of presenting the Hamiltonian H_s arises from the fact that the product $(\vec S_H \cdot \vec S_{\rm SLB})$ depends only on the overall total spin of the state $|\vec J| = |\vec S_H + \vec S_{\rm SLB}|$, while the operator $\vec \Delta$ has only nondiagonal matrix elements between the spin-singlet and spin-triplet states. Considering noninteracting meson-antimeson pairs, one can use the Hamiltonian H_s to find the eigenstates:

$$1^-(2^+):\ (1^-_H \otimes 1^-_{\rm SLB})|_{J=2}, \qquad \tfrac{1}{2}\mu; \tag{2}$$

$$1^-(1^+):\ (1^-_H \otimes 1^-_{\rm SLB})|_{J=1}, \qquad -\tfrac{1}{2}\mu; \tag{3}$$

$$1^-(0^+): \frac{\sqrt3}{2}(0^-_H \otimes 0^-_{\rm SLB}) + \frac12 (1^-_H \otimes 1^-_{\rm SLB})|_{J=0}, \quad \tfrac12\mu; \tag{4}$$

$$1^-(0^+): \frac12(0^-_H \otimes 0^-_{\rm SLB}) - \frac{\sqrt3}{2} (1^-_H \otimes 1^-_{\rm SLB})|_{J=0}, \quad -\tfrac32\mu; \tag{5}$$

$$1^+(1^-): \frac{1}{\sqrt2}(0^-_H \otimes 1^-_{\rm SLB}) - \frac{1}{\sqrt2}(1^-_H \otimes 0^-_{\rm SLB}), \qquad \tfrac12\mu; \tag{6}$$

$$1^+(1^-): \frac{1}{\sqrt2}(0^-_H \otimes 1^-_{\rm SLB}) + \frac{1}{\sqrt2}(1^-_H \otimes 0^-_{\rm SLB}), \quad -\tfrac12\mu. \tag{7}$$

In these formulas, the first expression indicates the quantum numbers $I^G(J^P)$, the second describes the composition in terms of the heavy and SLB spin states, and the third gives the energy of the state relative to the SLB $B\bar B$ threshold $2\bar M$. Considering these energy shifts, one readily identifies the states (2), (4), and (6) as being at the $B^*\bar B^*$ threshold, the states (3) and (7) at the $B^*\bar B$ threshold and, finally, the state (5) at the $B\bar B$ threshold.

The major effect of the interaction, depending on the spin of the heavy quark, is the splitting of the (otherwise degenerate) thresholds for the scattering channels described by Eqs. (2)–(7). The actual interaction between mesons takes place in a finite volume, where the effect of the forces related to the spin of the heavy quark is small in comparison with the interaction giving rise to near-threshold singularities in either one or both the $0^-_{\rm SLB}$ and $1^-_{\rm SLB}$ channels. Thus, one comes to the conclusion that the expected family of the near-threshold molecular resonances is as shown in Fig. 1. The existence of the $J^P = 0^+$ states W_{b0} and W'_{b0} follows from the existence of the $Z_b(Z'_b)$ resonances, while the existence of the W_{b1} and W_{b2} is contingent on the presence of a near-threshold singularity in the $1^-_{\rm SLB}$ channel. It can be also noted that the W_{b1} state is a pure isovector bottomonium-like analog of the charmonium-like resonance X(3872), which is a pure $(1^-_H \otimes 1^-_{\rm SLC})$ state [4].

Clearly, the $H \otimes$ SLB spin structure described by Eqs. (2)–(7) also implies relations between the total widths of the W_{bJ} states:

$$\Gamma(W_{b2}) = \Gamma(W_{b1}) = \frac32 \Gamma(W_{b0}) - \frac12 \Gamma(W'_{b0}), \tag{8}$$

as well as relations for the rates of decays of the resonances to specific channels, e.g.

$$\Gamma(W_{b0} \to \Upsilon\rho):\Gamma(W'_{b0} \to \Upsilon\rho):\Gamma(W_{b1} \to \Upsilon\rho):\Gamma(W_{b2} \to \Upsilon\rho) = \frac34 : \frac14 : 1 : 1, \tag{9}$$

with a possible slight modification due to the kinematical difference in the phase space. Such exclusive decays can be used for an experimental identification of the resonances. This paper, however, concentrates on the radiative transitions in which the W_{bJ} states can be produced.

The details of the transitions from $\Upsilon(5S)$ crucially depend on its structure in terms of its decomposition into heavy (H) and SLB spin states. The resonance $\Upsilon(5S)$ is produced in e^+e^- annihilation by the electromagnetic current, which creates a $b\bar b$ pair in an ortho state 1^-_H. In the standard notation, the heavy quark pair is created in either

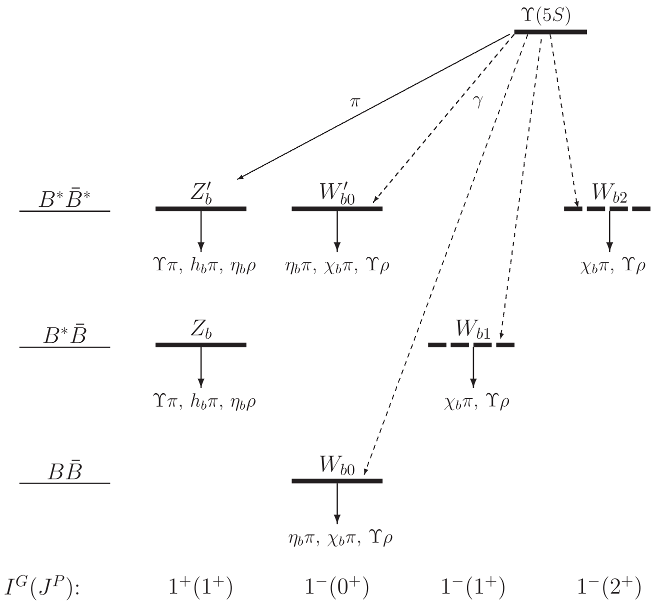

FIG. 1. The expected family of six isotriplet resonances at the $B\bar{B}$, $B^*\bar{B}$ and $B^*\bar{B}^*$ thresholds and their likely decay modes to bottomonium and a light meson. The excited bottomonium states can be present in the decays instead of the shown lower states (η_b, Υ, h_b, χ_b), where kinematically possible. The dashed arrowed lines show the discussed radiative transitions from $\Upsilon(5S)$. (The mass splitting to $\Upsilon(5S)$ is shown not to scale.)

a $3S_1$ state or in a $3D_1$ one. In terms of the $H \otimes SLB$ decomposition, the former is $1_H^- \otimes 0_{\rm SLB}^+$ and the latter is $1_H^- \otimes 2_{\rm SLB}^+$, since the angular momentum of the $b\bar{b}$ pair is relegated to the SLB system. However, the D-wave contribution in the production of the resonance is small inasmuch as the heavy quarks are nonrelativistic at the energy of the $\Upsilon(5S)$ resonance and can be neglected. It is thus reasonable to assume that the spin structure of $\Upsilon(5S)$ in terms of a $H \otimes SLB$ decomposition is dominated by $1_H^- \otimes 0_{\rm SLB}^+$.

To a certain extent, the assumed spin structure of the resonance can be tested against the available data on its decays. Namely, the relative yield of the meson-antimeson pairs $B^*\bar{B}^*$, $B^*\bar{B} + B\bar{B}^*$, and $B\bar{B}$ significantly depends on this structure. In the limit, where the interaction of the heavy quark spin is considered as small for a pure $1_H^- \otimes 0_{\rm SLB}^+$ state, the ratio of the yield in these channels is 7:3:1 [5] (see also the review in[6]). In the real world, there are finite effects due to the spin-dependent interaction both in the decay amplitudes and in the kinematical P-wave factors p^3 due to the mass splitting between the B^* and B mesons due to the same interaction. If only the phase space factors p^3 are taken into account, the ratio becomes 4.2:2.4:1. However, given the absence of a full calculation in the first order in the spin effects, it may be more reasonable to compare the lowest order theoretical result with the experimental data. Thus, the spread between the expected ratio with and without the kinematical factors illustrates the range of current theoretical uncertainty. The fraction for the yield in each of the three meson-antimeson channels at the $\Upsilon(5S)$ resonance as measured by Belle [7] corresponds to $f(B^*\bar{B}^*) = (37.5^{+2.1}_{-1.9} \pm 3.0)\%$, $f(B^*\bar{B} + B\bar{B}^*) = (13.7 \pm 1.3 \pm 1.1)\%$, $f(B\bar{B}) = (5.5^{+1.0}_{-0.9} \pm 0.4)\%$, which reasonably agrees with the 7:3:1 ratio, and given the errors, with the kinematically modified ratio. In either case, the suggested spin structure $1_H^- \otimes 0_{\rm SLB}^+$ of the $\Upsilon(5S)$ resonance appears to not contradict the data.

Another test of the suggested spin structure of $\Upsilon(5S)$ is provided by its decays into the channels $B^*\bar{B}^*\pi$, $(B^*\bar{B} + B\bar{B}^*)\pi$, and $B\bar{B}\pi$. The energy above the threshold for the heavy meson pair in these processes is small so that only the lowest possible partial wave amplitude can be retained when considering these decays. This is in agreement with the observed [7] suppression of the channel $B\bar{B}\pi$: $f(B\bar{B}\pi) = (0.0 \pm 1.2 \pm 0.3)\%$, since this process cannot go in the S wave unlike the other two. In the S wave, the states of the heavy meson pair in terms of the $H \otimes$ SLB decomposition can be read off the formulas (6) for $B^*\bar{B}^*$ and (7) for $(B^*\bar{B} + B\bar{B}^*)$. The heavy quark spin state is conserved so that the decays from $\Upsilon(5S)$ proceed only to the $1_H^- \otimes 0_{\rm SLB}^-$ component of the states of these heavy meson pairs. In other words, the underlying process can be viewed as factorized into the transition $(1_H^-)_{\Upsilon(5S)} \to (1_H^-)_{\rm final}$ for the heavy spin and $(0_{\rm SLB}^+)_{\Upsilon(5S)} \to (0_{\rm SLB}^-)_{\rm final} + \pi$ for the rest degrees of freedom. Clearly, since the states (6) and (7) contain the spin state $1_H^- \otimes 0_{\rm SLB}^-$ with the same amplitude (up to the sign), the ratio of the decay amplitudes of the observed processes can be found as

$$\left| \frac{A[\Upsilon(5S) \to B^*\bar{B}^*\pi(p_2)]}{A[\Upsilon(5S) \to (B^*\bar{B} + B\bar{B}^*)\pi(p_1)]} \right| = \frac{E_2}{E_1}, \tag{10}$$

where p_1 and p_2 are the momenta of the pion in these two decays, and E_1 and E_2 are the corresponding energies. (The proportionality of the S-wave amplitude to the pion energy is dictated by the chiral algebra. One can also note that the mechanism, where the discussed decays go through a pair with one virtual B or B^* meson with a subsequent emission of the pion by the heavy meson, $\Upsilon(5S) \to B^{(*)}\bar{B}^{(*)} \to B^{(*)}\bar{B}^{(*)}\pi$, results in higher partial waves, and its amplitude is proportional to the product of the pion spatial momentum and that of one of the heavy mesons. Such a contribution is expected to be suppressed at small available energy in the considered decays. This expectation is supported by the observed strong suppression of the channel $\Upsilon(5S) \to B\bar{B}\pi$, which is forbidden in the S wave and is generally allowed if higher partial waves are present.) Also, in Eq. (10), it is implied that state $(B^*\bar{B} + B\bar{B}^*)$ is the $G = +1$ state of the heavy meson pairs normalized to one. In these processes, the kinematical effect of the mass splitting between the B^* and B mesons is considerably enhanced by a very small released kinetic energy: about 75 MeV in the decay $\Upsilon(5S) \to B^*\bar{B}^*\pi$ and about 120 MeV in $\Upsilon(5S) \to (B^*\bar{B} + B\bar{B}^*)\pi$. It thus appears reasonable to take this kinematical effect into account. The estimate for the ratio of the total decay rates further depends on the distribution of the rate over the Dalitz plot. One can estimate the ratio of the relative yield $f(B^*\bar{B}^*\pi)/f(B^*\bar{B}\pi + B\bar{B}^*\pi)$ as approximately $1/2$ if the decay goes dominantly into the lowest invariant mass of the heavy meson pair (i.e,. if it is in fact dominated by the Z_b and Z_b' resonances), and as approximately $1/3.5$ if the spectrum of the invariant masses of the heavy pair is given by the phase space. Experimentally, [7] the total fractional rates are $f(B^*\bar{B}^*\pi) = (1.0^{+1.4}_{-1.3} \pm 0.4)\%$, $f(B^*\bar{B}\pi + B\bar{B}^*\pi) = (7.3^{+2.3}_{-2.1} \pm 0.8)\%$. To the best of my knowledge, a Dalitz analysis for these decays in not yet available. Given a visible relative suppression of the channel $B^*\bar{B}^*\pi$, one can expect, in the approach suggested here, that the Dalitz distribution in these decays should be spread over the physical region, rather than being dominated by the Z_b and Z_b' resonances.

The radiative decays $\Upsilon(5S) \to W_{bJ}\gamma$ can be considered in terms of the $H \otimes \text{SLB}$ decomposition in the same way as the decays into heavy meson pairs. Indeed, an emission of the photon by the spin of the heavy quarks is negligible so that in these decays, as before, the 1_H^- spin state "goes through" without a change in the orientation of the spin, while the photon emission occurs in the process $(0^+)_{\Upsilon(5S)} \to (1^-_{\text{SLB}})_{\text{final}} + \gamma$, whose polarization structure is described by only one amplitude in terms of the polarization amplitudes of the 1^-_{SLB} state ($\vec{\psi}$) and of the photon ($\vec{a}$):

$$A[(0^+)_{\Upsilon(5S)} \to (1^-_{\text{SLB}})_{\text{final}} + \gamma] = C\omega(\vec{\psi} \cdot \vec{a}), \quad (11)$$

where ω is the photon energy, and C is a constant, currently unknown, except its obvious dependence on fine structure constant: $C^2 \propto \alpha$. Using this structure of the amplitude and the amplitudes of the state $1_H^- \otimes 1^-_{\text{SLB}}$ in the W_{bJ} resonances, described by the relations (2)–(5), one can readily find the ratio of the rates of the discussed radiative transitions:

$$f(W_{b0}\gamma):f(W'_{b0}\gamma):f(W_{b1}\gamma):f(W_{b2}\gamma) = \frac{3}{4}\omega_0^3:\frac{1}{4}\omega_2^3:3\omega_1^3:5\omega_2^3, \quad (12)$$

where $\omega_{0,1,2}$ are the photon energies in the corresponding transitions: $\omega_0 \approx 305$ MeV, $\omega_1 \approx 260$ MeV, and $\omega_2 \approx 215$ MeV. If the kinematical ω^3 factors are taken into account in Eq. (12), then the ratio of the rates is estimated approximately as 8.5:1:21:20, rather than as 3:1:12:20 in the case where one ignores these factors on the grounds that their difference is just another effect of the heavy quark spin interaction. This, in fact, illustrates the range of uncertainty in the predictions based on the heavy quark limit for the discussed transitions. However, in spite of such an uncertainty, the presented estimates clearly indicate the relative feasibility of observing the yet hypothetical W_{bJ} resonances.

It can be noted that the isoscalar counterparts of the W_{bJ} resonances can in principle be also sought for in the radiative transitions from $\Upsilon(5S)$. Such C-even molecular resonances X_{bJ} were considered in Ref. [2]. One can expect, however, that the rates for the radiative transitions are small as compared to those for the isovectors. Indeed, the b quark and the antiquark are slow in both the initial and the final state in the transition. Thus, an emission of the photon by the heavy quarks can be neglected, and the photon is radiated by the current of the light u and d quarks. In the latter current, the isoscalar part is only $1/3$ of the isovector one in the amplitude, and one estimates

$$\frac{\Gamma[\Upsilon(5S) \to X_{bJ}\gamma]}{\Gamma[\Upsilon(5S) \to W_{bJ}\gamma]} \approx 1/9. \quad (13)$$

However, unlike the W_{bJ} states, the isoscalar resonances X_{bJ} can contain an admixture of χ_{bJ} bottomonium states, and thus can be possibly be produced with an observable rate in high-energy $p\bar{p}$ and pp collisions at the Tevatron and LHC [2], which provides a viable option for their discovery.

In summary, the relations between the rates of radiative transitions from $\Upsilon(5S)$ to the (yet hypothetical) isovector molecular bottomonium resonances W_{bJ} with $G = -1$ are considered in the HQSS limit using the decomposition of the states in terms of the heavy and SLB spin structure. It is argued that the $\Upsilon(5S)$ resonance is dominantly a $1_H^- \otimes 0^+_{\text{SLB}}$ state and that the existing data do not contradict such an assignment. The relations between the rates of the radiative transitions are then found by applying the HQSS to the radiative processes.

I am grateful to Alexander Bondar for many enlightening discussions. This work is supported, in part, by the DOE Grant No. DE-FG02-94ER40823.

[1] I. Adachi *et al.* (Belle Collaboration), arXiv:1105.4583.
[2] A. E. Bondar *et al.*, arXiv:1105.4473.
[3] D. Weingarten, Phys. Rev. Lett. **51**, 1830 (1983).
[4] M. B. Voloshin, Phys. Lett. B **604**, 69 (2004).
[5] A. De Rujula, H. Georgi, and S. L. Glashow, Phys. Rev. Lett. **38**, 317 (1977).
[6] M. B. Voloshin, Prog. Part. Nucl. Phys. **61**, 455 (2008).
[7] A. Drutskoy *et al.* (Belle Collaboration), Phys. Rev. D **81**, 112003 (2010).