

Selection Rules for Hadronic Transitions of XYZ Mesons

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Many of the XYZ mesons discovered in the last decade can be identified as bound states of a heavy quark and antiquark in Born-Oppenheimer (BO) potentials defined by the energy of gluon and light-quark fields in the presence of static color sources. The mesons include quarkonium hybrids, which are bound states in excited flavor-singlet BO potentials, and quarkonium tetraquarks, which are bound states in BO potentials with light-quark+antiquark flavor. The deepest hybrid potentials are known from lattice QCD calculations. The deepest tetraquark potentials can be inferred from lattice QCD calculations of static adjoint mesons. Selection rules for hadronic transitions are derived and used to identify XYZ mesons that are candidates for ground-state energy levels in the BO potentials for charmonium hybrids and tetraquarks.

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The XYZ mesons are unexpected mesons discovered during the last decade that contain a heavy quark-antiquark pair and are above the open-heavy-flavor threshold. Recent discoveries of charged XYZ mesons in both the $b\bar{b}$ sector and the $c\bar{c}$ sector have established unambiguously the existence of tetraquark mesons that contain two quarks and two antiquarks. In 2011, the Belle Collaboration discovered Z_b^+ (10610) and Z_b^+ (10650), which are revealed by their decays into $Y\pi^+$ to be tetraquark mesons with constituents $b\bar{b}u\bar{d}$ [1]. In 2013, the BESIII Collaboration discovered the Z_c^+ (3900), which is revealed by its decay into $J/\psi\pi^+$ to be a tetraquark meson with constituents $c\bar{c}u\bar{d}$ [2]. An updated list of the XYZ mesons as of August 2013 was given in Ref. [3]. The list included 15 neutral and 4 charged $c\bar{c}$ mesons. The BESIII Collaboration has recently observed an additional neutral state $Y(4220)$ [4] and additional charged states $Z_c^+(3885)$ [5] and $Z_c^+(4020)$ [6].

A full decade has elapsed since the discovery of the first XYZ meson, the $X(3872)$ [7], but no compelling explanation for the pattern of XYZ mesons has emerged. In simple constituent models, an XYZ meson consists of a heavy quark (Q) and antiquark (\bar{Q}) and possibly additional constituents that could be gluons (g) or light quarks (q) and light antiquarks (\bar{q}). The models that have been proposed can be classified according to how the constituents are clustered within the meson. They include (i) conventional quarkonium $(Q\bar{Q})_1$, (ii) quarkonium hybrid meson $(Q\bar{Q})_8 + g$, (iii) compact tetraquark $(Q\bar{Q}q\bar{q})_1$ [8], (iv) meson molecule $(Q\bar{q})_1 + (\bar{Q}q)_1$ [9], (v) diquarkonium $(Qq)_3 + (\bar{Q}\bar{q})_3$ [10], (vi) hadroquarkonium $(Q\bar{Q})_1 + (q\bar{q})_1$ [11], and (vii) quarkonium adjoint meson $(Q\bar{Q})_8 + (q\bar{q})_8$ [12]. The subscripts indicate the color charges of the clusters within the meson. All of these are possible models for neutral XYZ mesons. The last five are possible models for charged XYZ mesons. It would be desirable to have a theoretical framework firmly based on QCD that describes all the XYZ mesons, including their masses, widths, quantum numbers, and decay modes.

The Born-Oppenheimer (BO) approximation may provide such a theoretical framework.

The BO approximation is used in atomic and molecular physics to understand the binding of atoms into molecules. It exploits the large ratio of the time scale for the motion of the atomic nuclei to that for the electrons, which is a consequence of the large ratio of the nuclear and electron masses. The electrons respond almost instantaneously to the motion of the nuclei, which can be described by the Schrödinger equation in a BO potential defined by the energy of the electrons in the presence of static electric charges. The BO approximation for $Q\bar{Q}$ mesons in QCD was developed by Juge, Kuti, and Morningstar [13]. It exploits the large ratio of the time scale for the motion of the Q and \bar{Q} to that for the evolution of gluon fields, which is a consequence of the large ratio of the heavy-quark mass to the nonperturbative momentum scale Λ_{QCD} . The gluon field responds almost instantaneously to the motion of the $Q\bar{Q}$ pair, which can be described by the Schrödinger equation in a BO potential defined by the energy of the gluon field in the presence of static color sources. Conventional quarkonia are energy levels of a $Q\bar{Q}$ pair in the ground-state BO potential. The energy levels in the excited-state BO potentials are quarkonium hybrids. Juge, Kuti, and Morningstar calculated many of the BO potentials using quenched lattice QCD [13]. They calculated the spectrum of bottomonium hybrids by solving the Schrödinger equation in the BO potentials. They also calculated some of the bottomonium hybrid energies using lattice nonrelativistic QCD (NRQCD). The quantitative agreement between the predictions of the BO approximation and lattice NRQCD provided convincing evidence for the existence of quarkonium hybrids in the hadron spectrum of QCD. While direct lattice QCD calculations of the spectrum may be more accurate, an advantage of the BO approximation is that it reveals the pattern of the spectrum. As pointed out in Ref. [14], it is also possible to define BO

potentials with a light-quark+antiquark flavor by the energies of stationary configurations of light-quark and gluon fields with the specified flavor in the presence of static color sources. The energy levels of a $Q\bar{Q}$ pair in such a potential would be quarkonium tetraquarks.

In this Letter, we apply the BO approximation for quarkonium hybrids and tetraquarks to the $c\bar{c}XYZ$ mesons. The deepest flavor-singlet BO potentials have been determined using lattice QCD calculations. We infer the deepest BO potentials with light-quark+antiquark flavor from lattice QCD calculations of the static adjoint meson spectrum. We derive selection rules for hadronic transitions and use them to identify XYZ mesons that are candidates for ground-state energy levels of charmonium hybrids and tetraquarks.

The BO potentials for $Q\bar{Q}$ mesons can be labeled by quantum numbers for the gluon and light-quark fields that are conserved in the presence of static Q and \bar{Q} sources separated by a vector \mathbf{r} [13]: (i) the eigenvalue $\lambda = 0, \pm 1, \pm 2, \dots$ of $\hat{\mathbf{r}} \cdot \mathbf{J}_{\text{light}}$, where $\mathbf{J}_{\text{light}}$ is the total angular momentum vector for the light fields, (ii) the eigenvalue $\eta = \pm 1$ of $(CP)_{\text{light}}$, which is the product of the charge-conjugation operator and the parity operator that spatially inverts the light fields through the midpoint between the Q and \bar{Q} sources, (iii) for the case $\lambda = 0$, the eigenvalue $\epsilon = \pm 1$ of a reflection operator R_{light} that reflects the light fields through a plane containing the Q and \bar{Q} sources, (iv) quark flavors, which can be flavor singlet or $q_1\bar{q}_2$, where $q_1, q_2 = u, d, s$. For $q_1, q_2 = u, d$, the distinct BO potentials are specified by the isospin quantum number $I = 0, 1$. The value of $\Lambda = |\lambda|$ is traditionally specified by an upper-case Greek letter: $\Sigma, \Pi, \Delta, \dots$ for $\Lambda = 0, 1, 2, \dots$. The value $+1$ or -1 of η is traditionally specified by a subscript g or u on the upper-case Greek letter. In the case $\lambda = 0$, the value $+1$ or -1 of ϵ is traditionally specified by a superscript $+$ or $-$ on Σ . Thus the flavor-singlet BO potentials $V_\Gamma(r)$ are labeled by $\Gamma = \Sigma_\eta^\pm, \Pi_\eta^\pm, \Delta_\eta^\pm, \dots$, where η is g or u , which specify the eigenvalues of $|\hat{\mathbf{r}} \cdot \mathbf{J}_{\text{light}}|$, $(CP)_{\text{light}}$, and R_{light} . We will refer to the light-field configurations with those energies as Born-Oppenheimer configurations.

In QCD without light quarks, the ground-state flavor-singlet BO potential $V_{\Sigma_g^+}(r)$ can be defined as the minimal energy of the gluon field in the presence of the static Q and \bar{Q} sources. An excited flavor-singlet BO (or hybrid) potential $V_\Gamma(r)$, can be defined as the minimal energy of the gluon field with quantum numbers Γ only if $V_\Gamma(r)$ does not exceed $V_{\Sigma_g^+}(r)$ by more than the mass of a glueball with the appropriate quantum numbers. In QCD with light quarks, the minimal-energy prescription breaks down if $V_\Gamma(r)$ exceeds $V_{\Sigma_g^+}(r)$ by more than 2 or 3 times the mass of a pion, depending on the quantum numbers Γ . It also breaks down if $V_\Gamma(r)$ exceeds twice the energy of a static meson, which is the minimal energy for light-quark and gluon fields with the flavor of a single light quark in the

presence of a static \bar{Q} source. Similar complications arise in the definition of a $q_1\bar{q}_2$ BO (or tetraquark) potential. In all these cases, if the BO potential exists, it must be defined by a more complicated prescription involving excited states of the light fields with the specified quantum numbers.

Many of the hybrid potentials were calculated by Juge, Kuti, and Morningstar using quenched lattice QCD [13,15], which does not include virtual quark-antiquark pairs. At large r , they approach linear functions of r . At small r , they approach the repulsive Coulomb potential between a Q and \bar{Q} in a color-octet state, which is approximately linear in $1/r$. The deepest hybrid potentials are Π_u^\pm and then Σ_u^- , which is equal to the Π_u^\pm potential at $r = 0$. In the limit $r \rightarrow 0$, the Q and \bar{Q} sources reduce to a local color-octet $Q\bar{Q}$ source, and the BO configuration reduces to a static hybrid meson or gluelump, which is a flavor-singlet state of the light fields bound to a static color-octet source. The gluelump spectrum was first calculated using quenched lattice QCD by Campbell, Jorysz, and Michael [16]. It was recently calculated by Marsh and Lewis using lattice QCD with dynamical light quarks [17]. The ground-state gluelump has J_{light}^{PC} quantum numbers 1^{+-} . In the limit $r \rightarrow 0$, the Π_u^\pm and Σ_u^- potentials differ from the repulsive Coulomb potential for a color-octet $Q\bar{Q}$ pair by an additive constant that can be interpreted as the energy of the ground-state 1^{+-} gluelump.

The tetraquark potentials can be specified by Λ_η^ϵ and quark flavors $q_1\bar{q}_2$. None of them have yet been calculated using lattice QCD. Some information about these potentials at small r can be inferred from lattice QCD calculations of static adjoint mesons, which are states with light-quark+antiquark flavor bound to a static color-octet source. Foster and Michael have calculated the adjoint meson spectrum using quenched lattice QCD with a light valence quark and antiquark [18]. The $q\bar{q}$ adjoint mesons with the lowest energies are a vector ($J_{\text{light}}^{PC} = 1^{--}$) and a pseudoscalar (0^{-+}). The central values of their energies were larger than that of the ground-state 1^{+-} gluelump by about 50 and 100 MeV, respectively, but the differences were within the statistical errors. Lattice QCD calculations with dynamical light quarks would be required to determine definitively the ordering of the three energies. For each adjoint meson, there must be tetraquark potentials that in the limit $r \rightarrow 0$ approach the repulsive Coulomb potential for a color-octet $Q\bar{Q}$ pair plus an additive constant that can be interpreted as the energy of the adjoint meson. If these potentials remain well defined at large r , it is possible that they increase linearly with r , like the flavor-singlet BO potentials. In this case, the tetraquark potentials would have the same qualitative behavior as the hybrid potentials.

Given the quantum numbers J_{light}^{PC} of a $q\bar{q}$ adjoint meson, we can deduce the corresponding BO potentials. The component $\hat{\mathbf{r}} \cdot \mathbf{J}_{\text{light}}$ for an adjoint meson with spin J_{light} has $2J_{\text{light}} + 1$ integer values ranging from $-J_{\text{light}}$ to $+J_{\text{light}}$. There must, therefore, be a BO potential for each integer

value of Λ from 0 up to J_{light} . The quantum number η for the BO potentials is the value of $(CP)_{\text{light}}$ for the adjoint meson. One of the BO potentials is a Σ potential with reflection quantum number $\epsilon = (-1)^{J_{\text{light}}} P_{\text{light}}$. Thus the BO potentials associated with the 1^{--} adjoint meson are Π_g^\pm and Σ_g^\pm , while the BO potential associated with the 0^{-+} adjoint meson is Σ_u^- . Since the 1^{--} and the 0^{-+} adjoint mesons have the lowest energies, it is reasonable to expect the deepest tetraquark potentials to be Π_g^\pm , Σ_g^\pm , and Σ_u^- .

There are several angular momenta that contribute to the spin vector \mathbf{J} of a $Q\bar{Q}$ meson. In addition to $\mathbf{J}_{\text{light}}$, there is the orbital angular momentum $\mathbf{L}_{Q\bar{Q}}$ and the total spin \mathbf{S} of the $Q\bar{Q}$ pair. The spin vector of the meson can be expressed as $\mathbf{J} = \mathbf{L} + \mathbf{S}$, where $\mathbf{L} = \mathbf{L}_{Q\bar{Q}} + \mathbf{J}_{\text{light}}$. The condition $\mathbf{r} \cdot \mathbf{L}_{Q\bar{Q}} = 0$ implies $\hat{\mathbf{r}} \cdot \mathbf{L} = \lambda$, where λ is the quantum number for $\hat{\mathbf{r}} \cdot \mathbf{J}_{\text{light}}$. This puts a lower limit on the quantum number L for L^2 : $L \geq \Lambda$. For a flavor-singlet $Q\bar{Q}$ meson with BO configuration Λ_η^ϵ , the parity and charge-conjugation quantum numbers are

$$P = \epsilon(-1)^{\Lambda+L+1}, \quad (1a)$$

$$C = \eta\epsilon(-1)^{\Lambda+L+S}. \quad (1b)$$

The $Q\bar{Q}$ mesons are conveniently organized into heavy-quark spin-symmetry multiplets consisting of states with the same BO configuration Λ_η^ϵ , radial quantum number n , orbital-angular-momentum quantum number L , and $q_1\bar{q}_2$ flavor. Each multiplet consists of a spin-singlet ($S = 0$) state and either one or three spin-triplet ($S = 1$) states. Conventional quarkonia are energy levels in the flavor-singlet Σ_g^+ potential. The spin-symmetry multiplet for the ground state in this potential consists of a spin-singlet 0^{-+} state and a spin-triplet 1^{--} state. The lowest-energy quarkonium hybrids are energy levels in the flavor-singlet Π_u^+ , Π_u^- , and Σ_u^- potentials. The ground-state spin-symmetry multiplets in these potentials are given in Table I. Tetraquark $Q\bar{Q}$ mesons are energy levels in BO potentials labeled by Λ_η^ϵ and light-quark+antiquark flavor. The spin-symmetry multiplets for tetraquark $Q\bar{Q}$ mesons are most easily specified

TABLE I. Ground-state spin-symmetry multiplets for the two deepest hybrid potentials Π_u^\pm and Σ_u^- and for the Π_g^\pm , Σ_g^\pm , and Σ_u^- tetraquark potentials with flavor $q\bar{q}$. A bold \mathbf{J} indicates that J^{PC} is an exotic quantum number that is not possible if the constituents are only $Q\bar{Q}$.

Quarkonium hybrids			$Q\bar{Q}q\bar{q}$ tetraquarks		
$\Gamma(nL)$	$S = 0$	$S = 1$	$\Gamma(nL)$	$S = 0$	$S = 1$
$\Pi_u^+(1P)$	1^{--}	$(0, \mathbf{1}, 2)^{--}$	$\Pi_g^-(1P)$	1^{+-}	$(0, \mathbf{1}, 2)^{++}$
$\Pi_u^-(1P)$	1^{++}	$(\mathbf{0}, \mathbf{1}, 2)^{+-}$	$\Pi_g^+(1P)$	$\mathbf{1}^{+-}$	$(\mathbf{0}, \mathbf{1}, 2)^{--}$
$\Sigma_u^-(1S)$	0^{++}	1^{+-}	$\Sigma_g^+(1S)$	0^{-+}	1^{--}
			$\Sigma_u^-(1S)$	0^{++}	1^{+-}

by giving the J^{PC} quantum numbers for tetraquark mesons with flavor $q\bar{q}$. The ground-state spin-symmetry multiplets in the Π_g^+ , Π_g^- , Σ_g^+ , and Σ_u^- potentials are given in Table I. The J^{PC} quantum numbers are those for $I = 0$ and $s\bar{s}$ tetraquarks and for the neutral member of the $I = 1$ isospin triplet. The charged members of the $I = 1$ triplet have the same J^P and G -parity $G = -C$.

Most of the observed decay modes of the XYZ mesons are hadronic transitions to a quarkonium. Selection rules for the hadronic transitions provide constraints on the quarkonium hybrids or tetraquarks that can be considered as candidates for specific XYZ mesons. The spin selection rule $S = S'$, where S and S' are the total spin quantum numbers for the $Q\bar{Q}$ pair before and after the transition, follows from the approximate heavy-quark spin symmetry, which is a consequence of the large mass of the heavy quark. The Born-Oppenheimer selection rules also require the BO approximation, in which the hadronic transition between $Q\bar{Q}$ mesons proceeds through a transition between BO configurations with fixed separation vector \mathbf{r} for the Q and \bar{Q} sources. For simplicity, we will deduce these selection rules for transitions between neutral $Q\bar{Q}$ mesons with quantum numbers J^{PC} and $J'^{P'C'}$. The corresponding selection rules involving charged tetraquark mesons can then be deduced from isospin symmetry. We consider a transition via the emission of a single hadron h with quantum numbers $J_h^{P_h C_h}$ and orbital-angular-momentum quantum number L_h . The conservation of the component of $\mathbf{J}_{\text{light}}$ along the $Q\bar{Q}$ axis can be expressed as $\lambda = \lambda' + \hat{\mathbf{r}} \cdot (\mathbf{J}_h + \mathbf{L}_h)$, where \mathbf{J}_h and \mathbf{L}_h are the spin and orbital-angular-momentum vectors of h . This constraint implies the selection rule

$$|\lambda - \lambda'| \leq J_h + L_h. \quad (2)$$

Conservation of $(CP)_{\text{light}}$ implies the selection rule

$$\eta = \eta' \cdot C_h P_h (-1)^{L_h}. \quad (3)$$

If $\lambda = \lambda' = 0$, there is an additional constraint from invariance under reflection through a plane containing the $Q\bar{Q}$ axis:

$$\epsilon = \epsilon' \cdot P_h (-1)^{L_h} \quad (\lambda = \lambda' = 0). \quad (4)$$

The amplitude for a transition that violates the selection rules is generally suppressed by a factor of $\Lambda_{\text{QCD}}/M_Q v^2$, where v is the typical velocity of the heavy quark. The suppression may be avoided if the wave function has significant support near the avoided crossing between two BO potentials.

We proceed to apply the selection rules to the $c\bar{c}$ XYZ mesons. The spin selection rule implies that XYZ mesons with transitions to the spin-singlet $\Sigma_g^+(1P)$ charmonium state h_c must be spin-singlet states while those with transitions to the spin-triplet $\Sigma_g^+(1S)$ charmonium state

J/ψ or the spin-triplet $\Sigma_g^+(1P)$ charmonium state χ_{c1} must be spin-triplet states. This puts strong constraints on XYZ mesons with quantum numbers 1^{--} . In the hybrid multiplets in Table I, the only 1^{--} state is the spin-singlet state in $\Pi_u^+(1P)$. In the tetraquark multiplets in Table I, the only 1^{--} states are spin-triplet states in $\Pi_g^+(1P)$ and $\Sigma_g^+(1S)$. The 1^{--} meson $Y(4220)$ decays into $h_c\pi^+\pi^-$ [4], so it must be a spin singlet. It can only be identified with the 1^{--} state in the $\Pi_u^+(1P)$ multiplet of the charmonium hybrid. The 1^{--} meson $Y(4260)$ decays into $J/\psi\pi^+\pi^-$ [19]. In Ref. [14], it was identified as the spin-singlet member of the $\Pi_u^+(1P)$ multiplet of the charmonium hybrid. However, the decay into $J/\psi\pi^+\pi^-$ would then violate the spin selection rule. We instead identify it as the spin-triplet 1^{--} state in either the $\Pi_g^+(1P)$ or $\Sigma_g^+(1S)$ multiplet of the isospin-0 charmonium tetraquark.

Several of the hadronic transitions for the neutral $c\bar{c}XYZ$ mesons are the emission of a single vector meson ω or ϕ with $J_h^{P_h C_h} = 1^{--}$. Since the kinetic energy of the vector meson is small compared to its mass, we assume it is emitted in an S -wave state. The BO selection rules in Eqs. (2)–(4) reduce to $|\lambda - \lambda'| \leq 1$, $\eta = \eta'$, and also $\epsilon = -\epsilon'$ if $\lambda = \lambda' = 0$. If the final-state configuration is Σ_g^+ corresponding to quarkonium, the selection rules reduce further to $\Lambda \leq 1$, $\eta = +1$, and also $\epsilon = -1$ if $\Lambda = 0$. The only possible initial-state configurations are Π_g^-, Π_g^+ , and Σ_g^- . This excludes all the hybrid multiplets in Table I. The only tetraquark multiplets in Table I that are allowed are $\Pi_g^-(1P)$ and $\Pi_g^+(1P)$. The $X(3915)$ decays into $J/\psi\omega$ and has quantum numbers 0^{++} [20]. It can be identified with the 0^{++} state in the $\Pi_g^-(1P)$ multiplet of isospin-0 charmonium tetraquarks. The $Y(4140)$ decays into $J/\psi\phi$ [21] and therefore has $C = +$. The ϕ in the final state suggests that the meson is an $s\bar{s}$ charmonium tetraquark. The only spin-triplet $C = +$ states in Table I are the $0^{++}, 1^{++}$, and 2^{++} states in $\Pi_g^-(1P)$. The mass difference of about 230 MeV between the $Y(4140)$ and the $X(3915)$ is approximately twice the difference between the constituent masses of the s and lighter quarks. This is compatible with the identifications of $Y(4140)$ and $X(3915)$ as states in $s\bar{s}$ and isospin-0 $\Pi_g^-(1P)$ multiplets, respectively.

The only hadronic transitions that have been observed for charged XYZ mesons are the emission of a single pion with $J_h^{P_h C_h} = 0^{-+}$. The Goldstone nature of the pion requires that it be emitted predominantly in a P -wave state if its momentum is much larger than its mass. The BO selection rules in Eqs. (2)–(4) reduce to $|\lambda - \lambda'| \leq 1$, $\eta = \eta'$, and also $\epsilon = \epsilon'$ if $\lambda = \lambda' = 0$. If the final-state configuration is Σ_g^+ corresponding to quarkonium, the selection rules reduce further to $\Lambda \leq 1$, $\eta = +1$, and also $\epsilon = +1$ if $\Lambda = 0$. The only possible initial-state configurations are Π_g^-, Π_g^+ , and Σ_g^+ . Isospin symmetry requires the initial configuration to have isospin 1. The only tetraquark multiplets in Table I that are allowed are $\Pi_g^-(1P)$, $\Pi_g^+(1P)$, and $\Sigma_g^+(1S)$. The $Z_c^+(3900)$ decays into $J/\psi\pi^+$ [2]. Its neutral isospin partner

$Z_c^0(3900)$ has $C = -$. The spin-triplet $C = -$ tetraquark states in Table I are the $(\mathbf{0}, 1, 2)^{--}$ states in $\Pi_g^+(1P)$ and the 1^{--} state in $\Sigma_g^+(1S)$. The $Z_c^+(4020)$ decays into $h_c\pi^+$ [6]. Its neutral isospin partner $Z_c^0(4020)$ has $C = -$. The only spin-singlet $C = -$ tetraquark state in Table I is the 1^{+-} state in $\Pi_g^-(1P)$. The $Z_1^+(4050)$ decays into $\chi_{c1}\pi^+$ [22]. Its neutral isospin partner $Z_1^0(4050)$ has $C = +$. The spin-triplet $C = +$ tetraquark states in Table I are the $0^{++}, 1^{++}$, and 2^{++} states in $\Pi_g^-(1P)$. The small mass difference between $Z_c(4050)$ and $Z_1(4020)$ is compatible with them being in the same $\Pi_g^-(1P)$ multiplet.

We used the spin selection rule to identify the $Y(4260)$ as a spin-triplet 1^{--} state in either the $\Pi_g^+(1P)$ or $\Sigma_g^+(1S)$ multiplet of isospin-0 charmonium tetraquarks. We used the spin and BO selection rules to identify the $Z_c(3900)$ as a spin-triplet state in either $\Pi_g^+(1P)$ or $\Sigma_g^+(1S)$ multiplet of isospin-1 charmonium tetraquarks. The Π_u^\pm hybrid potential is deeper than the Σ_u^- hybrid potential. If the Π_g^\pm tetraquark potential is similarly deeper than the Σ_g^+ tetraquark potential, the most plausible identifications would be $Z_c(3900)$ as a $\Pi_g^+(1P)$ state and $Y(4260)$ as a $\Sigma_g^+(1S)$ state.

Our selection rules for hadronic transitions do not provide very useful constraints on the few $b\bar{b}XYZ$ mesons that have been observed. The $Z_b^+(10610)$ and $Z_b^+(10650)$ have transitions by emission of a single pion into both the spin-triplet bottomonium states $\Upsilon(nS)$ and the spin-singlet bottomonium states $h_b(nS)$ [1]. This violation of the spin selection rule can be explained by the $Z_b^+(10610)$ having a large $B^*\bar{B}$ molecular component and the $Z_b^+(10650)$ having a large $B^*\bar{B}^*$ molecular component [23–25]. Within the BO approach, large molecular components can arise if an energy level is fortuitously close to the $B^*\bar{B}$ and $B^*\bar{B}^*$ thresholds, so that the wave function has significant support in the region of the avoided crossing between the tetraquark BO potential and the BO potential for two static mesons.

We have used the Born-Oppenheimer approximation to derive selection rules for hadronic transitions between $Q\bar{Q}$ mesons. They strongly constrain the $c\bar{c}XYZ$ mesons that can be candidates for ground-state energy levels in the BO potentials for charmonium hybrids and tetraquarks. The selection rules should provide valuable guidance in the search for additional XYZ states through their hadronic transitions. Lattice QCD calculations of the tetraquark potentials are needed to confirm that the deepest potentials have been correctly identified. They would also allow the BO approximation to be developed into a quantitative theoretical framework for understanding the XYZ mesons that is based firmly on QCD.

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