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Heavy quark spin structure in Z_b resonances

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We discuss the heavy quark spin structure of the recently observed "twin" resonances $Z_b(10610)$ and $Z_b(10650)$, assuming that these are mostly of a "molecular" type, i.e. that their internal dynamics is dominated by the coupling to meson pairs $B^*\bar{B} - B\bar{B}^*$ and $B^*\bar{B}^*$. We find that the state of the $b\bar{b}$ pair within the $Z_b(10610)$ and $Z_b(10650)$ resonances is a mixture of a spin-triplet and a spin-singlet of equal amplitude and with the phase orthogonal between the two resonances. Such a structure gives rise to specific relations between observable amplitudes that are in agreement with the data obtained recently by Belle. We also briefly discuss possible properties of the isotopically singlet counterparts of the newly found resonances, and also of their C (G) parity opposites that likely exist in the same mass range near the open B flavor threshold.

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Very recently, the isotriplet resonances $Z_b(10610)$ and $Z_b(10650)$ were discovered in the processes $\Upsilon(5S) \to \pi\pi h_b(kP)$, and $\Upsilon(5S) \to \pi\pi\Upsilon(nS)$, (see Ref. [1]). Here and below, $n = 1, 2, 3$ and $k = 1, 2$. For simplicity, we refer to $Z_b(10610)$ and $Z_b(10650)$ as Z_b and Z_b'. Data analysis has shown that these processes go mainly as cascades, e.g., $\Upsilon(5S) \to Z_b\pi \to h_b\pi\pi$. The new bottomonium-type resonances apparently have quantum numbers $I^G(J^P) = 1^+(1^+)$, so that their electrically neutral isotopic states with $I_3 = 0$ should have $J^{PC} = 1^{+-}$. It turns out that the process with $\Upsilon(nS)$ in the final state has almost the same probability as those with h_b. At first glance, this fact looks quite astonishing. Assuming that b and $\bar{b}$ quarks are in a triplet spin state in $\Upsilon(5S)$ and $\Upsilon(nS)$, and they are in a singlet spin state in h_b, one may expect that a process with spin flip should be suppressed in comparison with that without spin flip because of the large mass of the b quark. In this Letter, we suggest an explanation of the puzzle. First of all, we note that the masses of the newly found states Z_b and Z_b' are close to the respective thresholds of the open B flavor channels $B^*\bar{B}$ and $B^*\bar{B}^*$. Therefore, it is natural to suggest that the new resonances have a "molecular" type structure of $B(B^*)$ meson pairs. Namely, the states with the quantum numbers of Z_b and Z_b' can be realized as S-wave $B^*\bar{B}$ and $B^*\bar{B}^*$ meson pairs, respectively. A possible "molecular" type structure in the charmonium family was initially discussed in Ref. [2].

The mass differences between the charged and neutral $B(B^*)$ mesons are negligible so that, unlike at the charm threshold, the isotopic symmetry should be well applicable to bottomoniumlike multiquark states. Therefore, we suggest that at long distances, $r \gg \Lambda_{\rm QCD}^{-1}$, the wave functions of the Z_b and Z_b' resonances are that of an S-wave meson pair in the $I^G(J^P) = 1^+(1^+)$ state, namely, $B^*\bar{B} - B\bar{B}^*$ for the Z_b and $B^*\bar{B}^*$ for the Z_b'. At shorter distances, $r \sim \Lambda_{\rm QCD}^{-1}$, the mesons overlap and form a system containing the heavy quark pair and a light component of quarks and gluons with the quantum numbers of an isotopic triplet.

In the limit $m_b \gg \Lambda_{\rm QCD}$, the spin degrees of freedom of b quark in the wave functions Ψ of B (B^*) mesons can be separated from other degrees of freedom. Thus, we treat a hyperfine interaction in B-meson as a perturbation. As a result, the wave function Ψ can be written as a direct product $\bar{\psi}_{\bar{q}} \otimes \chi_b$, where spinor χ_b describes the spin state of b-quark and $\bar{\psi}_{\bar{q}}$ describes the wave function of the bound light antiquark $\bar{q}$ and spinless b-quark. The total angular momentum j corresponding to the wave function ψ is fixed in the ground state $B(B^*)$: $j = 1/2$, and the rules of constructing the wave function Ψ are the same as in the nonrelativistic quark model. The precision of this picture is determined by the ratio $\Lambda_{QCD}/m_b = O(0.1)$ and the expected corrections should be at the level of about 10%.

For B meson, we have $\Psi_B = \bar{\psi}_{\bar{q}}\chi_b$, and for B^* meson we have $\vec{\Psi}_{B^*} = \bar{\psi}_{\bar{q}}\vec{\sigma}\chi_b$, where $\vec{\sigma}$ are the Pauli matrices. (It can be noticed that in this notation for the antimesons one has $\Psi_{\bar{B}} = \bar{\chi}_{\bar{b}}\psi_q$ and $\vec{\Psi}_{\bar{B}^*} = -\bar{\chi}_{\bar{b}}\vec{\sigma}\psi_q$.) Then, the S- state of the heavy meson pairs with the appropriate quantum numbers $I^G(J^P) = 1^+(1^+)$ is

$$\begin{aligned} &i\epsilon_{ijk}(\bar{\chi}_{\bar{b}}\sigma^j\psi_q)(\bar{\psi}_{\bar{Q}}\sigma^k\chi_b) \\ &\quad = (\bar{\chi}_{\bar{b}}\chi_b)(\bar{\psi}_{\bar{Q}}\sigma^i\psi_q) - (\bar{\chi}_{\bar{b}}\sigma^i\chi_b)(\bar{\psi}_{\bar{Q}}\psi_q) \\ &\quad \sim 0^-_{\bar{b}b} \otimes 1^-_{\bar{Q}q} - 1^-_{\bar{b}b} \otimes 0^-_{\bar{Q}q}, \end{aligned} \tag{1}$$

for $B^*\bar{B}^*$, and

$$\begin{aligned} &(\bar{\chi}_{\bar{b}}\sigma^i\psi_q)(\bar{\psi}_{\bar{Q}}\chi_b) + (\bar{\chi}_{\bar{b}}\psi_q)(\bar{\psi}_{\bar{Q}}\sigma^i\chi_b) \\ &\quad = -(\bar{\chi}_{\bar{b}}\chi_b)(\bar{\psi}_{\bar{Q}}\sigma^i\psi_q) - (\bar{\chi}_{\bar{b}}\sigma^i\chi_b)(\bar{\psi}_{\bar{Q}}\psi_q) \\ &\quad \sim 0^-_{\bar{b}b} \otimes 1^-_{\bar{Q}q} + 1^-_{\bar{b}b} \otimes 0^-_{\bar{Q}q}, \end{aligned} \tag{2}$$

for $B^*\bar{B} - B\bar{B}^*$ (The notation q and Q is used here to emphasize the fact that the light (anti)quarks are not necessarily of the same flavor.) Here we used the Fierz transforms, and 0^- and 1^- stand for para- and ortho- states with the negative parity. Clearly, the relations (1) and (2) refer only to the spin variables of the quarks. These relations describe the perfect mixtures of the two possible states corresponding to the para- and ortho-spin states of the $b\bar{b}$ pair. We thus conclude that, if the Z'_b and Z_b peaks are determined by a molecular dynamics of the meson pairs, their heavy quark spin structure should be the same as of the pairs, i.e.

$$|Z'_b\rangle = \frac{1}{\sqrt{2}}(0^-_{\bar{b}b} \otimes 1^-_{\bar{Q}q} - 1^-_{\bar{b}b} \otimes 0^-_{\bar{Q}q}),$$
$$|Z_b\rangle = \frac{1}{\sqrt{2}}(0^-_{\bar{b}b} \otimes 1^-_{\bar{Q}q} + 1^-_{\bar{b}b} \otimes 0^-_{\bar{Q}q}). \tag{3}$$

Such treatment of the heavy quark spin degrees of freedom in a molecular state is quite similar to that of the charmoniumlike peak X(3872) [3], whose spin structure is mandated to be [4] $(1^-_{c\bar{c}} \otimes 1^-_{u\bar{u}})_{J=1}$.

Since the masses of Z_b and Z'_b are very close to sum of the B and $\bar{B}^*$ masses and B^* and $\bar{B}^*$ masses, respectively, the mixture of states in Eq. (3) is small. In this picture, the mass splitting between the peaks should be equal to that between the B and B^* mesons: $M(Z'_b) - M(Z_b) = M(B^*) - M(B) \approx 46$ MeV with an expected correction of order $\Lambda^2_{\rm QCD}/m_b = O(1 \div 5)$ MeV. The spin structure described by Eq. (3) also leads to an important and experimentally testable conclusion that the resonances Z_b and Z'_b should have the same width. Indeed, in the large m_b limit all the ortho- and para- states of the $b\bar{b}$ pair are degenerate, so that the antisymmetric and the symmetric superposition of the spin states in Eq. (3) decay into degenerate (and orthogonal) states with lower mass, so that the total decay rates of the discussed resonances should be almost equal: $\Gamma(Z_b) = \Gamma(Z'_b)$. In particular, this also implies that in our approximation, the decays of the type $Z'_b \to B^*\bar{B}$ are forbidden by the heavy quark spin symmetry, in spite of being perfectly allowed by the overall quantum numbers and the kinematics. In other words, the heavy quark spin wave function in the Z'_b is orthogonal to that in the Z_b state.

The maximal ortho-para mixing of the heavy quarks in the Z_b and Z'_b resonances described by Eq. (3) immediately implies that these resonances have coupling of comparable strength to channels with states of ortho- and para-bottomonium. Furthermore, for each specific channel the absolute value of the coupling is the same for Z_b and Z'_b. However, the relative phase of the coupling of these resonances to the ortho- bottomonium is opposite to that for the para-bottomonium. In particular, the coupling of these resonances to the channels $\Upsilon(nS)\pi$ and $h_b(kP)\pi$ can readily be found (up to an overall normalization) as

$$E_\pi \vec{\Upsilon} \cdot (\vec{Z}_b - \vec{Z}'_b), \qquad (\vec{p}_\pi \times \vec{h}_b) \cdot (\vec{Z}_b + \vec{Z}'_b), \tag{4}$$

with $\vec{Z}'_b$, $\vec{\Upsilon}$ and $\vec{h}_b$ standing for the polarization amplitude of the corresponding spin one state, and E_π and $\vec{p}_\pi$ being the pion energy and momentum. The amplitudes described by Eq. (4) can be directly applied to the resonance part of the amplitudes of the observed transitions $\Upsilon(5S) \to \Upsilon(nS)\pi^+\pi^-$ and $\Upsilon(5S) \to h_b(kP)\pi^+\pi^-$. We have

$$\begin{aligned} &A(\Upsilon(5S) \to \Upsilon(nS)\pi^+\pi^-) \\ &\quad = A^{\rm nr}_\Upsilon + C_\Upsilon(\vec{\Upsilon}_5 \cdot \vec{\Upsilon})E_+E_-\bigg(\frac{1}{E - E_+ + \frac{i}{2}\Gamma} \\ &\qquad + \frac{1}{E' - E_+ + \frac{i}{2}\Gamma'} + \frac{1}{E - E_- + \frac{i}{2}\Gamma} + \frac{1}{E' - E_- + \frac{i}{2}\Gamma'}\bigg), \end{aligned} \tag{5}$$

and

$$\begin{aligned} &A(\Upsilon(5S) \to h_b(kP)\pi^+\pi^-) \\ &\quad = A^{\rm nr}_h + C_h\bigg\{[\vec{\Upsilon}_5 \cdot (\vec{p}_- \times \vec{h}_b)] \\ &\qquad \times E_+\bigg(\frac{1}{E - E_+ + \frac{i}{2}\Gamma} - \frac{1}{E' - E_+ + \frac{i}{2}\Gamma'}\bigg) \\ &\qquad + [\vec{\Upsilon}_5 \cdot (\vec{p}_+ \times \vec{h}_b)]E_- \\ &\qquad \times \bigg(\frac{1}{E - E_- + \frac{i}{2}\Gamma} - \frac{1}{E' - E_- + \frac{i}{2}\Gamma'}\bigg)\bigg\}, \end{aligned} \tag{6}$$

where E_+ and $\vec{p}_+$ (E_- and $\vec{p}_-$) stand for the energy and momentum of the positive (negative) pion, the parameters E and Γ (E' and Γ') are those of the Z_b (Z'_b) resonance with $E = M[\Upsilon(5S)] - M[Z_b]$ and $E' = M[\Upsilon(5S)] - M[Z'_b]$, the vector $\vec{\Upsilon}_5$ is the polarization amplitude of the initial $\Upsilon(5S)$ resonance, and $\vec{\Upsilon}$ and $\vec{h}_b$ are the same for, respectively, the final $\Upsilon(nS)$ and $h_b(kP)$ resonances. Furthermore, the coefficients C_Υ and C_h are constants, and, finally, $A^{\rm nr}_{\Upsilon,h}$ are the corresponding nonresonant amplitudes. The latter amplitudes generally depend on the polarizations and the kinematical variables [5,6] and can be studied in much the same way as for other similar two-pion transitions between heavy quarkonium states. It can be stated,however, that the nonresonant amplitude $A^{\rm nr}_h$ should be heavily suppressed due to the heavy quark spin symmetry and an absence of enhancing factors [7]. In Eqs. (5) and (6) we take into account two isotopic resonant branches, through the Z^+_b (Z'^+_b) and through Z^-_b (Z'^-_b).

Clearly, the Eqs. (5) and (6) describe two different patterns of the interference between the Z_b and Z'_b resonances in the two considered transitions. In the process $\Upsilon(5S) \to \Upsilon(nS)\pi^+\pi^-$, the interference is destructive when the energy of one of the pions lies between the positions of the resonances, E' and E, and the interference is constructive when both energies lie outside of the "twin" resonance band on the Dalitz plot. In the transition $\Upsilon(5S) \to h_b(kP)\pi^+\pi^-$, the pattern is exactly the opposite: the interference is constructive inside the resonance band

and is destructive outside the band. In fact, the probability of the latter transition outside the "twin" resonance band on the Dalitz plot should be very small due to the mentioned suppression of the nonresonant amplitude $A_h^{\rm nr}$. Such picture is fully supported by the experimental results of Ref. [1].

Based on the considerations presented here, one can expect the existence of hadronic transitions from the Z_b and Z_b' resonances to other ortho- and para-bottomonium states. In particular, the transitions $Z_b \to \eta_b \rho$ are of an S-wave type and a significant rate is possible for this process. Following the same consideration as presented above, the resonance amplitudes of the cascade $\Upsilon(5S) \to \eta_b \rho \pi$ should have the opposite sign between the Z_b and Z_b' Breit-Wigner factors. The processes of the type $Z_b(Z_b') \to \chi_b(1P)\pi\pi$ are also kinematically possible, but could be suppressed because two pions have to be in the $I^G = 1^+$ state and the ρ peak is beyond the kinematical region. The finding of the Z_b and Z_b' resonances may call for revisiting the analyses of the previously known processes, such as the transitions $\Upsilon(3S, 4S) \to \Upsilon(1S, 2S)\pi\pi$ as well as the decay $\Upsilon(3S) \to h_b(1P)\pi\pi$, for which a significant upper bound has become available recently [8]. A contribution of an isovector bottomoniumlike resonance in the decay $\Upsilon(3S) \to \Upsilon(1S)\pi\pi$ was, in fact, discussed some time ago [9,10].

The existence or nonexistence of "molecular" bottomoniumlike resonances depends on details of a yet unknown dynamics. However, the very existence of the Z_b and Z_b' resonances necessarily implies that additional isovector peaks also exist. Indeed, the resonance properties are determined by the interaction of the quasiparticles, which are the bound states of light quark and spinless b antiquark, while the spin of the $b\bar{b}$ pair plays only a "classificational" role. In particular, the emergence of the $Z_b(Z_b')$ resonances can be due to a near threshold singularity in either the $0^-_{\bar{Q}q}$ or the $1^-_{\bar{Q}q}$ state (or both) in the $I = 1$ channel. The total width of the Z_b and Z_b' in the range $15 \div 20$ MeV makes the distinction between the specific type of the singularity, a shallow bound state, a virtual state, or a resonance, rather moot. In either of these cases, there should exist a pair of isotriplet singularities at the respective $B^*\bar{B}^*$ and $B\bar{B}$ thresholds with spin 0 and the G parity opposite to that of $Z_b(Z_b')$, i.e. with $I^G(J^P) = 1^-(0^+)$. This is because the threshold S-wave states of $B^*\bar{B}^*$ and $B\bar{B}$ pairs with such overall quantum numbers are expressed as mixed combinations of the light and heavy quark spin states:

$$\begin{aligned}(B^*\bar{B}^*)|_{1^-(0^+)} &\sim \frac{(\bar{\chi}_{\bar{b}}\boldsymbol{\sigma}\psi_q)(\bar{\psi}_{\bar{Q}}\boldsymbol{\sigma}\chi_b)}{\sqrt{3}} \\ &= \frac{\sqrt{3}}{2}(\bar{\chi}_{\bar{b}}\chi_b)(\bar{\psi}_{\bar{Q}}\psi_q) - \frac{1}{2}\frac{(\bar{\chi}_{\bar{b}}\boldsymbol{\sigma}\chi_b)(\bar{\psi}_{\bar{Q}}\boldsymbol{\sigma}\psi_q)}{\sqrt{3}} \\ &\sim \frac{\sqrt{3}}{2}0^-_{\bar{b}b} \otimes 0^-_{\bar{Q}q} - \frac{1}{2}1^-_{\bar{b}b} \otimes 1^-_{\bar{Q}q}, \end{aligned} \tag{7}$$

and

$$\begin{aligned}(B\bar{B})|_{1^-(0^+)} &\sim (\bar{\chi}_{\bar{b}}\psi_q)(\bar{\psi}_{\bar{Q}}\chi_b) \\ &= \frac{1}{2}(\bar{\chi}_{\bar{b}}\chi_b)(\bar{\psi}_{\bar{Q}}\psi_q) + \frac{\sqrt{3}}{2}\frac{(\bar{\chi}_{\bar{b}}\boldsymbol{\sigma}\chi_b)(\bar{\psi}_{\bar{Q}}\boldsymbol{\sigma}\psi_q)}{\sqrt{3}} \\ &\sim \frac{1}{2}0^-_{\bar{b}b} \otimes 0^-_{\bar{Q}q} + \frac{\sqrt{3}}{2}1^-_{\bar{b}b} \otimes 1^-_{\bar{Q}q}. \end{aligned} \tag{8}$$

Thus, the mixing angle between para- and ortho-spin states of the $b\bar{b}$ pair in this case is $\pi/6$, which can be readily checked experimentally. Clearly, such resonances form another 'twin' pair and couple to both ortho- and para-bottomonium and can thus decay to e.g. $\Upsilon\rho$ as well as to $\eta_b\pi$.

Additionally, if a threshold singularity $1^-_{\bar{Q}q}$ state contributes to the $Z_b(Z_b')$ resonances, its combination with $1^-_{\bar{b}b}$ state of the heavy quark pair should also produce an $I^G(J^P) = 1^-(1^+)$ state at the $B^*\bar{B}$ threshold (an isovector bottomonium analog of X(3872) [3]), and isospin triplet at the $B^*\bar{B}^*$ threshold with spin 2: $I^G(J^P) = 1^-(2^+)$. These latter states couple to ortho-bottomonium, e.g., to the channel $\Upsilon\rho$. Because of the negative G parity, all these expected resonances are not accessible in single pion transitions from $\Upsilon(5S)$, but their production can be sought for at a somewhat higher energy above the B flavor threshold. One can also notice that the isotriplet states cannot mix with pure bottomonium states, so that they are unlikely to be produced at Tevatron and/or LHC at a detectable rate.

At this point, the "molecular" interaction in the isoscalar channel is not known. However, based on the existence of the $X(3872)$ state in the charmonium family, one may expect an existence of $I = 0$ counterparts of the isotriplet Z_b and Z_b' in the same mass region near the $B^*\bar{B}$ and $B^*\bar{B}^*$ thresholds. Such states, Y_b and Y_b', have the quantum numbers $J^{\rm PC} = 1^{+-}$ and can mix with $^1P_1$ states of bottomonium, $h_b(kP)$. Such mixing can generally tilt the completely mixed orto- para-heavy quark spin structure in the $Y_b(Y_b')$ resonances, and it would be interesting if this behavior could be studied experimentally. For the reasons of isospin, these resonances are not accessible from the $\Upsilon(5S)$ by single pion transitions, but could be studied in the future at higher initial energies in e^+e^- annihilation. Moreover, the likely presence of the bottomonium $^1P_1$ "core" in the $Y_b(Y_b')$ states makes it possible that, unlike the $Z_b(Z_b')$, these states can be produced in hard processes such as the high-energy $p\bar{p}$ or pp collisions at the Tevatron and LHC, similarly to the production of $X(3872)$ at the Tevatron [11]. The discussed bottomoniumlike resonances can be identified, e.g., by their decay into $\Upsilon(2S)\eta$, or $\Upsilon(1S)\eta$, or $\Upsilon(1S)\eta'$ which all are S-wave processes and which one would not expect to be suppressed. Other possibly identifiable in a collider setting decay channels of $Y_b(Y_b')$ are $\Upsilon(1S)\pi\pi$ and $\Upsilon(1S)K\bar{K}$, including those through the $f_0(980)$ resonance: $Y_b(Y_b') \to \Upsilon(1S)f_0(980)$, although an expectation for the rate of these latter processes is a somewhat subtle due to the required orbital momentum of the light mesons. A comparison of the decay rates to states with and without hidden strangeness could also shed light on the significance of an admixture in these resonances of the states of the type $b\bar{b}s\bar{s}$. The C-even states X_b of the same type can mix with the ${}^3P_J$ bottomonium and can similarly be produced in hard collisions. These resonances can be sought for at the colliders by their decay, e.g., into $\Upsilon(1S)\omega$.

In summary: We argue that, if the newly found Z_b and Z_b' isovector resonances are states of a "molecular" type in, respectively, the channels $B^*\bar{B} - B\bar{B}^*$ and $B^*\bar{B}^*$ with the quantum numbers $I^G(J^P) = 1^+(1^+)$, each of them has to contain an (almost) complete mixture of spin-triple and spin-singlet states of the heavy $b\bar{b}$ pair. The heavy quark spin wave functions in the two resonances have to be orthogonal to each other, as described by Eq. (3). In our approach, using a separation of the b quark spin degrees of freedom in the wave function of B mesons, based on the large value of the b quark mass m_b, the mass splitting between Z_b' and Z_b should be the same as between B^* and B mesons: $M(Z_b') - M(Z_b) = M(B^*) - M(B) \approx$ 46 MeV, and their total decay widths should be equal to one another: $\Gamma(Z_b') = \Gamma(Z_b)$. Any deviations from these relations are due to the finite mass of b quark and should be small. In particular, a kinematically allowed process $Z_b' \to B^*\bar{B}$ should be strongly suppressed. Furthermore, the resonances Z_b and Z_b' should have equal coupling to specific decay channels with the states of bottomonium. The relative sign between the couplings of Z_b and Z_b' to such channels depends on the spin state of the bottomonium, this relative sign of the coupling to ortho-states is opposite to that in the coupling to the para- states. Such behavior leads to a specific interference pattern in the contribution of the discussed resonances to the observed processes $\Upsilon(5S) \to \Upsilon(nS)\pi^+\pi^-$ and $\Upsilon(5S) \to h_b(kP)\pi^+\pi^-$. Finally, we point out that a similar behavior can be tested in the yet unobserved, but expected, processes $\Upsilon(5S) \to \eta_b\rho\pi$ and $\Upsilon(5S) \to \chi_b(1P)\pi\pi\pi$. The coupling of the bottomonium states to the Z_b and Z_b' resonances can also affect the rates and the spectra in hadronic transitions in bottomonium, such as $\Upsilon(3S) \to \Upsilon(1S)\pi\pi$ and/or $\Upsilon(3S) \to h_b(1P)\pi\pi$. We also suggest that isospin-singlet resonances Y_b and Y_b' with the quantum numbers $J^{PC} = 1^{+-}$ of a similar "molecular" structure can mix with the ${}^1P_1$ states of bottomonium and can thus be produced in "hard" collisions at the Tevatron and LHC. These states can be sought for in the high-energy data by their decay channels $\Upsilon(2S, 1S)\eta$, or $\Upsilon(1S)\eta'$, or $\Upsilon(1S)\pi\pi(\text{or}K\bar{K})$, and a possible isoscalar resonance X_{b0} with $J^{PC} = 1^{++}$ can be sought for by its decay into $\Upsilon(1S)\omega$.

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