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Michael I. Eides
University of Kentucky, michael.eides@uky.edu

Victor Yu. Petrov
Petersburg Nuclear Physics Institute, Russia

Maxim V. Polyakov
Petersburg Nuclear Physics Institute, Russia

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Digital Object Identifier (DOI)
https://doi.org/10.1103/PhysRevD.93.054039

Notes/Citation Information
Published in *Physical Review D*, v. 93, issue 5, 054039, p. 1-6.

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PHYSICAL REVIEW D **93,** 054039 (2016)

Narrow nucleon-$\psi(2S)$ bound state and LHCb pentaquarks

Michael I. Eides,[1,2] Victor Yu. Petrov,[2] and Maxim V. Polyakov[2,3]

[1]*Department of Physics and Astronomy, University of Kentucky, Lexington, Kentucky 40506, USA*
[2]*Petersburg Nuclear Physics Institute, Gatchina 188300, St. Petersburg, Russia*
[3]*Institut für Theoretische Physik II, Ruhr-Universität Bochum, D-44780 Bochum, Germany*

(Received 2 December 2015; revised manuscript received 10 February 2016; published 28 March 2016)

We interpret the newly discovered pentaquark $P_c(4450)$ as a bound state of charmonium $\psi(2S)$ and the nucleon. The binding potential is due to the charmonium-nucleon interaction that in the heavy quark approximation is proportional to the product of the charmonium chromoelectric polarizability and the nucleon energy-momentum distribution. We use the large N_c expansion to estimate the quarkonium polarizability and calculate the nucleon properties in the framework of the mean-field picture of light baryons. Two almost degenerate states $J^P = (1/2)^-$ and $J^P = (3/2)^-$ are predicted at the position of the $P_c(4450)$ pentaquark. We find that the nucleon-$\psi(2S)$ bound state has a naturally narrow width in the range of tens of MeV. The unitary multiplet partners of the $P_c(4450)$ pentaquark and the generalization to $b\bar{b}$-nucleon pentaquark bound states are discussed.

DOI: 10.1103/PhysRevD.93.054039

The discovery of new pentaquark states by the LHCb Collaboration [1] opens the problem of their internal structure. A few interesting ideas were already proposed: the pentaquark as a loosely bound state of a charmed baryon and a meson [2], the pentaquark as a bound state of light and heavy diquarks with a c-quark [3], and even the pentaquark as a bound state of states with open color [4]. It was also suggested in [5] that the structures found by the LHCb Collaboration can be interpreted as threshold cusp effects.

In this paper, we explore another option: the pentaquark as a bound state of a charmonium state and a nucleon. A heavy quark-antiquark bound state is a small (compared to the size of a nucleon) heavy neutral object. Its interaction with a nucleon is relatively weak even when the distance between the quarkonium and the nucleon is small. A quarkonium can easily penetrate the nucleon and form a true pentaquark state. Strong interactions of a heavy quarkonium are naturally described in the framework of the nonrelativistic multipole expansion [6]. The quarkonium-nucleon interaction is dominated by virtual emission of two chromoelectric dipole gluons in a color singlet state. The effective heavy quarkonium-nucleon interaction potential is proportional to the product of the meson chromoelectric polarizability and the local gluon energy-momentum density inside the nucleon [7].

Chromoelectric polarizability of a very heavy quarkonium was calculated a long time ago [8–10].[1] Nondiagonal (transitional) polarizabilities also can be calculated in this approach. It is questionable how close the real heavy quark systems ($c\bar{c}$ or $b\bar{b}$ quarkonia) are to the pure Coulomb system. Phenomenological values of the transitional polarizabilities can be extracted, e.g., from the experimental data on the $\psi' \to J/\psi\pi\pi$ decays [7]. There is at least a qualitative agreement between the Coulombic and phenomenological values of nondiagonal polarizabilities.

The simplest but not too accurate estimate of the gluon energy-momentum density inside a nucleon is provided by the Skyrme soliton model [12]. We use the QCD-inspired chiral quark-soliton model (χQSM) [13] to calculate the gluon energy-momentum density inside a nucleon. The χQSM model was very successful in describing virtually all of the low-energy physics of the interacting nucleons and pseudoscalar mesons [14]. It arises in QCD in the large N_c limit and unambiguously leads to the mean-field picture of baryons [15] that we use in the calculations below. Let us mention that the Θ^+ pentaquark [16] and the charmed pentaquark [17] were earlier predicted in the χQSM model. However, the physical nature of those pentaquarks is completely different from the mechanism considered here. The two main ingredients of the present discussion, the small size of quarkonium and the quarkonium-nucleon interaction, played no role in those predictions. There is nothing special about the description of the nucleon in the χQSM mean-field picture for our present goals. Any model that guarantees that the quarkonium-nucleon binding energy is parametrically small in comparison with the nucleon mass can be used for the calculation of the nucleon energy-momentum distribution instead of the χQSM model.

The effective quarkonium interaction with light hadrons described above is attractive. It was used to discuss possible quarkonium bound states in light nuclear matter [18]. It was also applied to the interpretation of the exotic mesons with hidden charm [19]. A tentative interpretation of the LHCb pentaquarks as bound states of J/ψ and the nucleon resonances $N(1450)$ and $N(1520)$ was suggested in [20].

[1]See also the recent calculation [11] for the $1S$ state.

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Our estimates show that the quarkonium-nucleon interaction is not strong enough to bind together the charmonium ground state J/ψ and an individual nucleon. However, Coulombic chromoelectric polarizability increases like the cube of the quarkonium radius. We expect that the fast growth of polarizability with an increase in the radius of the heavy quark-antiquark bound state holds even for non-Coulombic systems. As a result, interaction of a nucleon with excited quarkonia is much stronger than interaction with J/ψ, and bound nucleon-excited quarkonia states should exist.

We obtain an attraction potential of about a few hundred MeV with the size of about 1 fm between the soliton and the excited $\psi(2S)$ state, just enough to form a bound state. We interpret this bound state as the $P_c(4450)$ pentaquark discovered by the LHCb Collaboration. We calculated the width of this bound state, which turned out to be rather small, about a few tens of MeV, consistent with the LHCb results [1]. We predict that the pentaquark $P_c(4450)$ is doubly degenerate. This degeneracy is due to a calculable spin-spin nucleon-quarkonium interaction. This interaction is suppressed by the heavy quark mass, which leads to degeneracy in the leading order of the heavy quark expansion. We also predict a rich spectrum of new pentaquark states that arise from the binding of the quarkonium states with the ordinary baryons. These new pentaquark states form flavor multiplets similar to the well-known baryon octets, decuplets, etc. The pattern of the masses and properties of these new pentaquarks can be calculated.

We use the multipole expansion to calculate the interaction of the heavy quarkonium with the light hadrons [21]. The role of the small parameter in this expansion plays the ratio of the quarkonium size over the effective gluon wavelength. The leading term in this expansion is due to two dipole gluons and can be parametrized in terms of chromoelectric polarizability α. The effective dipole Lagrangian has the form [21]

$$L_{\mathrm{eff}} = \frac{\alpha}{2}\boldsymbol{E}\cdot\boldsymbol{E}, \tag{1}$$

where $\boldsymbol{E}$ is the chromoelectric gluon field (with the coupling constant absorbed) and α is the chromoelectric polarizability.

The chromoelectric polarizabilities of the charmonium states are not known now, except in the case of very heavy quarks. For such quarks, the quarkonium is a Coulombic system, and polarizability admits a perturbative calculation [8–11] both with and without an additional expansion in the large N_c. The leading term of the large N_c expansion for the polarizability at $N_c = 3$ differs from the exact in N_c result by 5.5%. This difference is negligible for our goals, and we calculate the polarizability in the framework of the $1/N_c$ expansion. The polarizability for an arbitrary quarkonium nS energy level is

$$\alpha(nS) = \frac{16\pi n^2}{3g^2N_c^2}c_n a_0^3, \tag{2}$$

where $c_1 = 7/4$, $c_2 = 251/8$, and $c_n(n \geq 3) = (5/16)n^2(7n^2 - 3)$, $a_0 = 16\pi/(g^2 N_c m_q)$ is the Bohr radius of nonrelativistic quarkonium, and g is the coupling constant normalized at the size of the quarkonium. The nondiagonal $(2S \to 1S)$ chromoelectric polarizability is

$$\alpha(2S \to 1S) = -\frac{51200\sqrt{2}\pi}{1287g^2N_c^2}a_0^3. \tag{3}$$

Other transitional polarizabilities can be calculated in the same way.

We use the Coulombic values for the polarizabilities as order-of-magnitude estimates of their scale and characteristic features, but we do not rely on their numerical values. Fitting the J/ψ and ψ' masses, we extract the Bohr radius and the Coulomb values for the polarizabilities,[2]

$$\begin{aligned} \alpha(1S) &\approx 0.2\ \mathrm{GeV}^{-3}, \\ \alpha(2S) &\approx 12\ \mathrm{GeV}^{-3}, \\ \alpha(2S \to 1S) &\approx -0.6\ \mathrm{GeV}^{-3}. \end{aligned} \tag{4}$$

Transitional polarizability $|\alpha(2S \to 1S)| \approx 2\ \mathrm{GeV}^{-3}$ was extracted from the phenomenological analysis of the $\psi' \to J/\psi\pi\pi$ transitions [7]. There is a rather significant discrepancy between the perturbative result and this value. It could be explained by the non-Coulombic nature of the quarkonium. We expect that calculations with a more realistic potential lead to a better agreement with the phenomenological value of the polarizability.

The chromoelectric field squared in the Lagrangian in Eq. (1) can be easily connected with the gluon part of the QCD energy-momentum tensor T_{00}^G and, via the conformal anomaly, with the trace of the full energy-momentum tensor $T^\mu{}_\mu$,[3]

$$\boldsymbol{E}^2 = \frac{\boldsymbol{E}^2 - \boldsymbol{H}^2}{2} + \frac{\boldsymbol{E}^2 + \boldsymbol{H}^2}{2} = g^2\left(\frac{8\pi^2}{bg_s^2}T^\mu{}_\mu + T_{00}^G\right).$$

Here, $b = (11/3)N_c - (2/3)N_f$ is the leading coefficient of the Gell-Mann-Low function, and g_s is the strong coupling constant at a low normalization point. Notice that due to running of the coupling constant in QCD, $g \neq g_s$. The coupling constant g is defined at the scale of the quarkonium radius, while g_s is defined at the scale of the

[2] The result may vary slightly depending on how one treats a large N_c limit.

[3] We ignore the contribution of the light quarks' mass term. Simple estimates show that this term shifts the mass of the pentaquarks by only about 10 MeV upwards and hence can be safely neglected for all practical purposes.

nucleon radius. It seems that we can safely ignore this distinction in the case of charmonium, but it could become important for bottomonium.

Now, we are ready to adjust the effective Lagrangian in Eq. (1) for analysis of the quarkonium interaction with a light hadron. To this end, we average the operator in Eq. (1) over the hadron state and obtain

$$\mathcal{L}_{\rm eff} = \frac{\alpha}{2} g^2 \left(\frac{8\pi^2}{b g_s^2} T^\mu{}_\mu + T_{00}^G \right) = \frac{\alpha}{2} g^2 \left(\frac{8\pi^2}{b g_s^2} T^\mu{}_\mu + \xi T_{00} \right), \tag{5}$$

where $T^\mu{}_\mu$ and T_{00} are now expectation values of the respective operators in the light hadron state. At the last step, we also introduced a new parameter ξ that describes the fraction of the nucleon energy carried by the gluons at a low normalization point, $T_{00}^G = \xi T_{00}$.

We analyze the quarkonium-nucleon interaction with the help of the effective interaction Lagrangian in Eq. (5) using the χQSM model of the nucleon and the estimates of the chromoelectric polarizabilities above. Both the heavy quarkonium and the nucleon in the large N_c limit are nonrelativistic. In these conditions, the interaction Lagrangian in Eq. (5) describes a static interaction. The respective nonrelativistic potential can be written in terms of the local energy density $\rho_E(\boldsymbol{x})$ and pressure $p(\boldsymbol{x})$ [22],

$$V(\boldsymbol{x}) = -\alpha \frac{4\pi^2}{b} \left(\frac{g^2}{g_s^2} \right) \left[\rho_E(\boldsymbol{x}) \left(1 + \xi \frac{b g_s^2}{8\pi^2} \right) - 3 p(\boldsymbol{x}) \right]. \tag{6}$$

This effective potential has a simple interpretation. A pointlike quarkonium serves as a tool that scans the local energy density and local pressure inside the nucleon. It could happen that the size of quarkonium is not small enough in comparison with the size of the nucleon. In such a case, we need to consider the higher order terms in the QCD multipole expansion in order to improve the description of the quarkonium-nucleon interaction.

The overall normalization of the effective potential,

$$\int d^3x V(\boldsymbol{x}) = -\alpha \frac{4\pi^2}{b} \left(\frac{g^2}{g_s^2} \right) M_N \left(1 + \xi \frac{b g_s^2}{8\pi^2} \right), \tag{7}$$

is determined by the total energy of the nucleon $\int d^3x \rho_E(\boldsymbol{x}) = M_N$ and the stability condition $\int d^3x p(\boldsymbol{x}) = 0$. The factor $\nu = 1 + \xi(b g_s^2/8\pi^2)$ is model dependent. An estimate of this factor for the pion in [23] produced $\nu \sim 1.45$–1.6. In the theory of instanton vacuum and the χQSM model, the strong coupling constant freezes at the size of the nucleon with a value of about $\alpha_s = g_s^2/4\pi \sim 0.5$. Using this coupling constant, we obtain $\nu \sim 1.5$ for the nucleon, which is close to the pion result in [23].

The local energy density $\rho_E(\boldsymbol{x})$ and pressure $p(\boldsymbol{x})$ were computed in the χQSM in [24]. Calculations involved the exact quark levels in the pion mean field (including the Dirac sea) and the solution of the self-consistent equations of motion for the mean field. In this approach, the normalization condition for the potential in Eq. (7) is satisfied automatically since the normalization condition for the energy density and the stability condition for the pressure hold in the self-consistent calculation due to the equations of motion.

The form of the nonrelativistic quarkonium-nucleon interaction potential in Eq. (6) is determined by the results of the self-consistent mean-field calculation in [24]; its overall strength is fixed by the values of the chromoelectric polarizabilities of the quarkonia. This potential is universal. The interaction of any quarkonium state with the nucleon is described by the same potential; only the scale of this interaction potential depends on the quarkonium energy levels. Explicitly, the quarkonium-nucleon potentials for the two lowest charmonium states have the form

$$\begin{aligned} V_{22}(r) &\equiv V(r), \\ V_{11}(r) &= \frac{\alpha(1S)}{\alpha(2S)} V(r), \\ V_{12}(r) &= \frac{\alpha(2S \to 1S)}{\alpha(2S)} V(r), \end{aligned} \tag{8}$$

where $V(r)$ is the potential in Eq. (6) with $\alpha = \alpha(2S)$. The nondiagonal potential $V_{12}(r)$ describes the transition $J/\psi \to \psi'$ off the nucleon. With the polarizabilities from Eq. (4), the potentials $V_{11}(r)$ are $V_{12}(r)$ are small in comparison with the potential $V(r)$.

Bound states in the channels $J/\psi + N$ and $\psi' + N$ are solutions of the eigenvalue problem for the Schrödinger equation

$$\left(-\frac{\boldsymbol{\nabla}^2}{2\mu} + V(r) - E \right) \Psi_b = 0, \tag{9}$$

where μ is the reduced mass in the respective channel, and the potentials are defined in Eq. (8). Due to the poor knowledge of the chromoelectric polarizability α, we can vary it in a relatively wide region.

We found the following:

(1) A bound state arises when the chromoelectric polarizability reaches the critical value $\alpha = 5.6\ \mathrm{GeV}^{-3}$. Comparing this polarizability with the Coulomb values in Eq. (4), we see that J/ψ does not form a bound state with the nucleon. For the excited charmonia states $\psi(2S)$, $\psi(3S)$, etc., the critical value of α is far below the expected chromoelectric polarizabilities of the excited charmonia. Therefore, they should form bound states with the mean-field nucleon. Here, we concentrate on the bound state(s)

of $\psi(2S)$, and higher excited charmonia will be considered elsewhere.

(2) A bound state with the orbital momentum $l = 0$ and with the binding energy $E_b = -176$ MeV [corresponding to the position of the $P_c^+(4450)$ pentaquark] is formed at $\alpha(2S) = 17.2\ \mathrm{GeV}^{-3}$. There is only one bound state with such a polarizability.

(3) A bound state with the orbital momentum $l = 0$ and with the energy $E_b = -246$ MeV [corresponding to the position of the $P_c^+(4380)$ pentaquark] is formed at $\alpha = 20.2\ \mathrm{GeV}^{-3}$. Again, there is only one bound state with such a polarizability. Hence, interpretation of $P_c^+(4380)$ as a bound state with $E_b = -246$ MeV would mean that there are no heavier pentaquarks in the $J/\psi + N$ channel.

(4) An additional bound state with angular momentum $l = 1$ arises at a slightly larger value of polarizability $\alpha \approx 22.4$. One could try to identify the light pentaquark with the $l = 0$ bound state and the heavy pentaquark with the $l = 1$ bound state. The quantum numbers of such pentaquarks would be $(3/2)^-$ and $(5/2)^+$, respectively, which fits the experimental data nicely. But the mass difference of these states is about 300 MeV, not the observed 70 MeV. This large mass difference between the rotational excitation and the ground state is due to the relatively small size (around 0.8–0.9 fm) of the nucleon and its respectively relatively small moment of inertia. In the mean-field picture of the nucleon, the moment of inertia determines the energy of its rotational excitations, which is about a few hundred MeV, as can be seen from the $N - \Delta$ mass difference. In addition, the scenario with two pentaquarks as the $l = 0$ and $l = 1$ bound states cannot explain the widths of the observed pentaquarks. We consider this scenario to be absolutely excluded.

We see that charmonium $\psi(2S)$ can form bound states with the mean-field nucleon. Fitting the binding energies of the LHCb pentaquarks, we found the values of the chromoelectric polarizability that ensure the necessary strength of the binding potential. Compared with the theoretical predictions in Eq. (4), these polarizabilities are close to what is expected. However, only one bound state exists for each realistic value of polarizability, and only one of the LHCb pentaquarks can be described in our picture. Experimentally, the $P_c(4380)$ peak has a rather large width of $205 \pm 18 \pm 86$ MeV, whereas the $P_c(4450)$ peak is narrow with a width of $39 \pm 5 \pm 19$ MeV. We will see below that the nucleon-$\psi(2S)$ bound state has a naturally narrow width of about a dozen MeV. Therefore, we identify the nucleon-$\psi(2S)$ bound state with the LHCb $P_c(4450)$ pentaquark.

The nucleon-$\psi(2S)$ bound state is formed in the S wave; hence, its quantum numbers could be either $J^P = (1/2)^-$ or $J^P = (3/2)^-$. The spin-spin interaction between the color singlet states (quarkonium and the nucleon) arises due to interference of the chromoelectric dipole $E1$ and the chromomagnetic quadrupole $M2$ transitions. Its strength is determined by the chromoelectric polarizability, but it is additionally suppressed by a heavy quark mass as $\sim 1/m_q$. Hence, in the leading order of the heavy quark expansion the $(1/2)^-$ and $(3/2)^-$ states are degenerate. A semiquantitative estimate of hyperfine splitting produces a small value in the range of 5–10 MeV. Thus, we predict that there are two almost degenerate pentaquark states with $J^P = (1/2)^-$ and $J^P = (3/2)^-$ at the position of the observed pentaquark at $M_{pJ/\psi} = 4450$ MeV. It would be very interesting if the LHCb Collaboration could check this hypothesis in their partial wave analysis.

In the scenario above, the partial decay width of the pentaquark to $J/\psi + N$ can be calculated unambiguously. To this end, we consider J/ψ scattering off the nucleon as a nonrelativistic two-channel problem,

$$
\begin{aligned}
\left(-\frac{\boldsymbol{\nabla}^2}{2\mu_1} + V_{11}(r) - E\right)\Psi_1 + V_{12}(r)\Psi_2 = 0, \\
\left(-\frac{\boldsymbol{\nabla}^2}{2\mu_2} + V_{22}(r) - E + \Delta\right)\Psi_2 + V_{12}(r)\Psi_1 = 0.
\end{aligned} \tag{10}
$$

Here, μ_1 and μ_2 are the reduced masses of $J/\psi + N$ and $\psi' + N$, respectively, E is the energy in the center-of-mass frame ($E = \boldsymbol{p}^2/2\mu_1$, where $\boldsymbol{p}$ is the relative momentum), $\Delta = M_{\psi'} - M_{J/\psi}$, and the potentials $V_{11}(r)$, $V_{22}(r)$, and V_{12} are defined in Eq. (8).

Due to the nonzero transition potential V_{12}, the pentaquark arises as a resonance in the $J/\psi N$ scattering channel described by the standard Breit-Wigner formula. We find the width of the resonance from the resonance scattering amplitude.

The transition potential V_{12} is small, and we solve the scattering problem in Eq. (10) using perturbation theory. Due to coupling between the channels, the incoming plane wave $\Psi_1(\boldsymbol{x}) = e^{i\boldsymbol{q}\cdot\boldsymbol{x}}$ in the first channel leaks in the second channel where it induces the wave function

$$
\Psi_2(\boldsymbol{x}) = -\int d^3x' G_2(\boldsymbol{x}, \boldsymbol{x}') V_{12}(\boldsymbol{x}') e^{i\boldsymbol{q}\cdot\boldsymbol{x}'}. \tag{11}
$$

Here,

$$
G_2(\boldsymbol{x}, \boldsymbol{x}') = \left\langle \boldsymbol{x} \left| \frac{1}{-\frac{\nabla^2}{2\mu_2} - E + \Delta + V - i0} \right| \boldsymbol{x}' \right\rangle
$$

is the Green function of the Schrödinger equation for $\Psi_2(\boldsymbol{x})$ [see Eq. (9)]. Near the resonance

$$
G_2(\boldsymbol{x}, \boldsymbol{x}') = \frac{\psi_R(\boldsymbol{x})\psi_R^*(\boldsymbol{x}')}{E_R - E},
$$

where E_R is the resonance energy. The wave function $\Psi_2(\boldsymbol{x})$ in Eq. (11) in its turn generates a correction to $\Psi_1(\boldsymbol{x})$ [see the first line in Eq. (10)] that near the resonance has the form

$$\delta\Psi_1(x) = \int d^3x' G_1(\boldsymbol{x},\boldsymbol{x}')V_{12}(\boldsymbol{x}')\psi_R^*(\boldsymbol{x}') \times \frac{\int d^3x'' V_{12}(\boldsymbol{x}'')\psi_R(\boldsymbol{x}'')e^{i\boldsymbol{q}\cdot\boldsymbol{x}''}}{E_R - E}, \tag{13}$$

where $G_1(\boldsymbol{x},\boldsymbol{x}')$ is the free Green function and $q = |\boldsymbol{q}| = \sqrt{2\mu_1 E}$.

The wave function in the first channel at large x is a superposition of the incoming plane wave and the outgoing spherical wave

$$\Psi_1(\boldsymbol{x}) + \delta\Psi_1(x) = e^{i\boldsymbol{q}\cdot\boldsymbol{x}} + f(\theta)\frac{e^{iqr}}{r}, \tag{14}$$

where $f(\theta)$ is the scattering amplitude (θ is the scattering angle). The scattering amplitude determined by the wave function in Eq. (13) has a standard Breit-Wigner resonance form,

$$f(\theta) = -\frac{2l+1}{q}\frac{\Gamma/2}{E-E_R}P_l(\cos\theta), \tag{15}$$

where Γ is the resonance partial decay width into the $N + J/\psi$ channel. Calculating this width, we obtain

$$\Gamma = \left(\frac{\alpha(2S\to 1S)}{\alpha(2S)}\right)^2 (4\mu_1 q)\left|\int_0^\infty drr^2 R_l(r)V(r)j_l(qr)\right|^2, \tag{16}$$

where $R_l(r)$ is the resonance radial wave function normalized by the condition $\int drr^2 R_l(r) = 1$, and $j_l(z)$ is the spherical Bessel function.

Numerically, we obtain $\Gamma(P_c(4450) \to N + J/\psi) \approx 11$ MeV for the phenomenological value of polarizability $\alpha(2S \to 1S) = 2\ \mathrm{GeV}^{-3}$ [7]. We also made a rough estimate of the partial width $P_c \to J/\psi + N + \pi$, and it turned out to be even smaller than the partial width into the $J/\psi + N$ channel. The decays of the pentaquark into an (anti)charmed meson + charmed baryon are strongly suppressed in the scenario above since decays of the pentaquark into open charm channels can go only via t-channel exchange by a heavy D meson. Therefore, the total width of the P_c pentaquark in our picture is small, in the range of tens of MeV, which is in excellent agreement with the experimentally observed width $\Gamma_{\mathrm{exp}} = 39 \pm 5 \pm 19$ MeV of the $P_c(4450)$ pentaquark.

To summarize, we have calculated the effective potential of the heavy quarkonium-nucleon interaction using the QCD multipole expansion and the mean-field description of the nucleon in the χQSM model. This potential depends on the quarkonium polarizability and the energy-momentum distribution inside the nucleon [24]. We found a $\psi(2S)$-nucleon bound state in this potential that arises at reasonable values of the chromoelectric polarizability $\alpha(2S)$. The polarizability can be adjusted is such a way that the mass of the bound state coincides with the position of either $P_c(4380)$ or $P_c(4450)$. Only one $\psi(2S)$-nucleon bound state arises in our approach,[4] and we cannot describe both resonances by the same mechanism. We have calculated the width of the nucleon-$\psi(2S)$ bound state and obtained a value in the range of tens of MeV. This width fits nicely with the width of the LHCb pentaquark $P_c(4450)$. Therefore, we identify the $\psi(2S)$-nucleon bound state with the narrow $P_c(4450)$ pentaquark. The wide $P_c(4380)$ pentaquark does not fit our picture; it should be explained in some other way, perhaps as some kind of a threshold enhancement. We predict that the $P_c(4450)$ peak consists of two almost degenerate pentaquark states with $J^P = (1/2)^-$ and $J^P = (3/2)^-$. This is at variance with the most favorable quantum number of $J^P = (5/2)^+$ obtained for this pentaquark in the analysis by the LHCb Collaboration [1].

The ability of the $c\bar{c}$ resonances to form bound states with baryons opens a new perspective on the world of pentaquarks. A compact weakly interacting quarkonium state bound inside a baryon does not change its properties in a significant way. Then, the spectrum of pentaquark states should duplicate all already known baryon multiplets. For example, the $P_c(4450)$ pentaquark should be a member of a baryon octet. Masses of other particles in this octet can be read off the table of baryons: we expect analogues of N, Σ, Ξ, and Λ. The next multiplet of pentaquarks is similar to the baryon decuplet and should consist of pentaquarks with properties similar to Δ, Σ, Ξ, and Ω. This is also not the end of the story—we see no reason why $\psi(2S)$ cannot form a bound state with the Roper resonance or any other known baryon with positive or negative parity.

The other opportunity to proliferate the number of pentaquark states even more is to consider possible bound states of baryons with other excited states of $c\bar{c}$ systems. It is also worth noticing that the spin-spin interaction between $c\bar{c}$ mesons and nucleons is very weak. This means that every pentaquark state should be accompanied by a nearly degenerate state with a different spin and the same parity.

In the scenario discussed above, the bottomonium states also should form bound states with the light baryons. Moreover, our considerations should become more reliable for systems with the b quarks as they are heavier and closer to the pure Coulomb systems. On the other hand, the $b\bar{b}$

[4]Possible bound states of the nucleon and higher excited states of charmonia [$\psi(3S)$, etc.] will be considered elsewhere.

mesons are more compact, and therefore, the respective chromoelectric polarizabilities are smaller. Very naively, the polarizabilities in bottomonia are suppressed by the factor $\left(\frac{\alpha_s(m_c)m_c}{\alpha_s(m_b)m_b}\right)^3$ (ratio of the Bohr radii cubed) in comparison with the polarizabilities in the charmonia. This estimate shows that the chromoelectric polarizability in bottomonia is close to the value that corresponds to the formation of a nucleon-$\Upsilon(2S)$ bound state. More accurate calculations are required. A more detailed study of the interaction of higher excited quarkonia with the nucleon is also warranted.

ACKNOWLEDGMENTS

This paper was supported by the NSF, Grant No. PHY-1402593. The work of V. P. is supported by the Russian Science Foundation, Grant No. 14-22-00281. M. V. P. is grateful to Hyun-Chul Kim and Michal Praszalowicz for fruitful discussions.

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