New LHCb pentaquarks as Hadrocharmonium States

Michael I. Eides,1,2,∗ Victor Yu. Petrov,2† and Maxim V. Polyakov2,3,‡

1Department of Physics and Astronomy, University of Kentucky, Lexington, KY 40506, USA
2Petersburg Nuclear Physics Institute, Gatchina, 188300, St.Petersburg, Russia
3Ruhr-University Bochum, Faculty of Physics and Astronomy, Institute of Theoretical Physics II, D-44780 Bochum, Germany

Abstract

New LHCb Collaboration results on pentaquarks with hidden charm [1] are discussed. These results fit nicely in the hadrocharmonium pentaquark scenario [2, 3]. In the new data the old LHCb pentaquark \( P_c(4450) \) splits into two states \( P_c(4440) \) and \( P_c(4457) \). We interpret these two almost degenerate states as a result of hyperfine splitting between two color singlet hadrocharmonium states with \( J^P = 1/2^- \) and \( J^P = 3/2^- \) that was predicted in [2]. We improve the theoretical estimate of hyperfine splitting [2, 3] that is compatible with the experimental data. The new \( P_c(4312) \) state finds a natural explanation as a bound state of \( \chi_{c0} \) and the nucleon, with \( I = 1/2, \ J^P = 1/2^+ \) and binding energy 42 MeV. As a bound state of a spin zero meson and a nucleon hadrocharmonium pentaquark \( P_c(4312) \) does not experience hyperfine splitting. We find a series of hadrocharmonium states in the vicinity of the wide \( P_c(4380) \) pentaquark what can explain its apparently large decay width. We compare the hadrocharmonium and molecular pentaquark scenarios and discuss their relative advantages and drawbacks.

* Email address: meides@uky.edu
† Email address: Victor.Petrov@thd.pnpi.spb.ru
‡ Email address: maxim.polyakov@tp2.ruhr-uni-bochum.de
I. INTRODUCTION

Pentaquarks with hidden charm were discovered by the LHCb Collaboration about five
years ago [4]. According to [4] there are two pentaquarks with hidden charm, one with a
mass of $4380 \pm 8 \pm 29$ MeV and a width of $205 \pm 18 \pm 86$ MeV, and another with a mass
of $4449.8 \pm 1.7 \pm 2.5$ MeV and a width of $39 \pm 5 \pm 19$ MeV. The preferred $J^P$ assignments
are of opposite parity, with one state having spin $3/2$ and the other $5/2$. There is now an
extensive theoretical literature on the interpretation of the LHCb pentaquarks, see, e.g., the
recent review [5]. We will discuss the hadrocharmonium scenario, suggested in [6–8] (heavy
quarkonium interaction with nuclei was considered in [9, 10], see also references in [11]).
Hadrocharmonium approach to the LHCb pentaquarks was developed further in [2, 3, 12–
15]. In our previous works [2, 3] we discussed interpretation of the LHCb pentaquarks as
hadrocharmonium states, nonrelativistic bound states of $\psi(2S)$ and the nucleon. We de-
scribed $P_c(4450)$ as a hadrocharmonium state [2] with $I = 1/2$, $J^P = 3/2^-$, and calculated
its decay widths [15]. In the leading approximation the binding potential in hadrocharmo-
nium does not depend on spin, so we predicted existence of a degenerate state with $I = 1/2$, $J^P = 1/2^-$. Degeneracy between the two color singlet states with $J^P = 1/2^-$, $3/2^-$ is lifted
by a hyperfine interaction arising in the QCD multipole expansion, and the magnitude of
the hyperfine splitting was estimated in [2, 3].

New experimental data on the LHCb pentaquarks was presented recently in [1]. Two
narrow states $P_c(4440)$ and $P_c(4457)$ are seen at the position of the old LHCb pentaquark
$P_c(4450)$. Also a new narrow resonance $P_c(4312)$ shows up in the experimental data.
We consider discovery of two narrow states $P_c(4440)$ and $P_c(4457)$ as a confirmation of
the prediction in [2, 3] of two almost degenerate pentaquark states with $J^P = 1/2^-$ and
$J^P = 3/2^-$ and with the mass of the observed pentaquark 4450 MeV. We will improve the
estimate [2, 3] of the hyperfine splitting between $J^P = 1/2^-$ and $J^P = 3/2^-$ pentaquarks
below. We will use the approach developed in [15] to calculate and compare partial and
total decay widths of these pentaquarks. An interpretation of the new LHCb pentaquark
$P_c(4312)$ as a hadrocharmonium bound state will be also discussed. The analysis in [1]
was not sensitive to broad resonances, there were no new information on the status of the
LHCb pentaquark $P_c(4380)$. We will elaborate on the hadrocharmonium scenario for this
pentaquark [3] below.

A natural theoretical framework for discussion of exotic mesons and baryons is provided
by QCD at large $N_c$. It predicts a qualitative difference between exotic mesons and baryons.
Light tetraquarks do not exist in QCD at large $N_c$ since meson-meson interaction decreases
as $1/N_c$, see e.g., [16]. However, if quarks are so heavy that $m_Q \gg N_c\Lambda_{QCD}$ ($\Lambda_{QCD}$ is
the scale of strong interactions) then even a shallow potential $1/N_c$ can bind mesons with
heavy quarks. The binding energy of such molecular state would be small since the binding
potential is proportional to $1/N_c$. Therefore, only molecular type tetraquarks that are loosely
bound states of two mesons with heavy quarks can exist in QCD at large $N_c$.

The case of exotic baryons is radically different. At large $N_c$ baryon consists of $N_c$ quarks,
its mass is proportional to $N_c$, and its interactions with mesons does not depend on $N_c$ at
all. Hence, QCD at large $N_c$ does not ban existence neither of light nor heavy exotic baryons
with the binding energy of order $\Lambda_{QCD}$. Unlike exotic mesons large $N_c$ exotic baryons could
be bound states of uniformly packed quarks having no resemblance to molecules, could have
molecular structure, or have a more complicated hadrocharmonium structure. One cannot
say in advance what is the true structure of the LHCb pentaquarks, this problem can be
The new LHCb resonances \( P_c(4440) \) and \( P_c(4457) \) are a few MeV below the \( D^{*0}\Sigma^+_c \) threshold 4460 MeV and the \( P_c(4312) \) pentaquark is a few MeV below the \( D^0\Sigma^+_c \) threshold 4318 MeV. Due to proximity to the respective thresholds an interpretation of the three pentaquarks as loosely bound molecular states was suggested by the LHCb Collaboration [1] and developed further in the recent literature [19–24]. We will compare hadrocharmonium and molecular interpretations of the LHCb pentaquarks and will argue that the hadrocharmonium scenario is no less natural, makes unambiguous quantitative predictions, and describes some fine features of the experimental data. Experimentally verifiable predictions of the hadrocharmonium scenario will be presented.

II. NEW LHCb DATA AND HADROCHARMONIUM SCENARIO

The \( P_c(4450) \) LHCb pentaquark was interpreted in [2] as a hadrocharmonium bound state of \( \psi(2S) \) and the nucleon. Binding potential in the hadrocharmonium picture in the leading approximation is proportional to the chromoelectric polarizability of the small color singlet \( c\bar{c} \) pair and is spin-independent. This is why the hadrocharmonium \( P_c(4450) \) in [2, 3] is an almost degenerate doublet of states with spin-parities \( J^P = 1/2^- \) and \( J^P = 3/2^- \). As shown in [2, 3] degeneracy between these color-singlet bound states is lifted by the hyperfine splitting that arises due to interference of the chromoelectric dipole \( E_1 \) and the chromomagnetic quadrupole \( M_2 \) transitions in charmonium. Hyperfine interaction is described by the effective interaction Hamiltonian [11]

\[
H_{\text{eff}} = -\frac{\alpha}{2m_Q} S_j \langle N(p')|E_i^a(D_iB_j)^a|N(p)\rangle,
\]

where \( E_i^a \) and \( B_j^a \) are chromoelectric and chromomagnetic fields, and \( S_j, \alpha \) and \( m_Q \) are the \( \psi(2S) \) spin, chromoelectric polarizability and the heavy (c) quark mass, respectively.

The strength of the hyperfine interaction is determined by the chromoelectric polarizability and it is additionally suppressed by the heavy quark mass \( \sim 1/m_Q \) in comparison with the binding potential, see [2, 3] for more detail. Only the nucleon matrix element of the product of chromoelectric and chromomagnetic fields between the nucleon states with momenta \( p \) and \( p' \) in Eq. (1) requires calculation. This matrix element can be written as

\[
\langle N(p')|E_i^a(D_iB_k)^a|N(p)\rangle = i q_i \langle N(p')|E_i^aB_k^a|N(p)\rangle - \langle N(p')(D_iE_i)^aB_k^a|N(p)\rangle,
\]

where \( q_i \) is the momentum transfer, \( q_i = p_i' - p_i \). Due to the QCD equations of motion the second term on the right hand side is a mixed quark-gluon operator suppressed by the QCD coupling constant \( g_s \) and we throw it away. As an additional argument in favor of its suppression let us mention that this term turns into zero in the instanton field. To estimate the first term we use an approximate relationship

\[
\langle N(p')|G_\mu^a\tilde{G}_\rho^a|N(p)\rangle \simeq \frac{1}{12} (g_{\mu\rho}g_{\nu\lambda} - g_{\nu\rho}g_{\mu\lambda}) \langle N(p')|G_{\alpha\beta}^a\tilde{G}_{\alpha\beta}^a|N(p)\rangle
\]

\[
+ \frac{1}{12} \epsilon_{\mu\nu\rho\lambda} \langle N(p')|G_{\alpha\beta}^aG_{\alpha\beta}^a|N(p)\rangle.
\]

To justify this approximation we notice that it is compatible with the general identities
\[ g^{\nu\lambda} \tilde{G}^a_{\mu\nu} \tilde{G}^a_{\rho\lambda} \equiv \frac{1}{4} g_{\mu\rho} G^a_{\alpha\beta} \tilde{G}^{a,\alpha\beta}_{\mu\nu} \quad \varepsilon^{\mu\nu\rho\lambda} G^a_{\mu\nu} \tilde{G}^a_{\rho\lambda} \equiv 2 G^a_{\alpha\beta} \tilde{G}^{a,\alpha\beta}_{\nu\lambda}, \] (4)

and holds in the instanton (antiinstanton) field. Effectively the approximation in Eq. (3) means that we omitted twist three and two gluon operators on the right hand side. The nucleon matrix elements of such operators are strongly suppressed \(^{25}\) in the instanton theory of QCD vacuum \(^{26,27}\). This suppression was confirmed experimentally \(^{28}\).

Returning to the matrix element in Eq. (2) we obtain with the help of Eq. (3)

\[ \langle N(p')|E_i^a (D_i B_k^a)|N(p)\rangle \approx \frac{i q_i}{12} \langle N(p')|G^a_{\alpha\beta} \tilde{G}^{a,\alpha\beta}|N(p)\rangle. \] (5)

The flavor singlet axial current in QCD is anomalous, what allows us to write the expression on the right hand side in terms of the singlet axial nucleon form factor \(g_A^{(0)}(q^2)\)

\[ \langle N(p')|G^a_{\alpha\beta} \tilde{G}^{a,\alpha\beta}|N(p)\rangle = \frac{32\pi^2}{N_f} g^{(0)}_A (q^2) m_N \bar{u}(p') i \gamma_5 u(p). \] (6)

For the nonrelativistic proton the momentum space hyperfine interaction Hamiltonian in Eq. (1) reduces to

\[ H_{\text{eff}} = \frac{\alpha}{m_Q} \frac{4\pi^2}{3N_f} g^{(0)}_A (-q^2) (S \cdot q) (s_N \cdot q), \] (7)

where \(S\) and \(s_N\) are the spin operators of \(\psi(2S)\) and the nucleon, respectively.

The coordinate space hyperfine potential for the \(S\)-wave hadrocharmonium bound state has the form

\[ V_{\text{hfs}}(r) = - (S \cdot s_N) \frac{\alpha}{m_Q} \frac{4\pi^2}{9N_f} \nabla^2 \int \frac{d^3q}{(2\pi)^3} e^{-iq \cdot r} g^{(0)}_A (-q^2). \] (8)

The dipole parameterization is usually used for the singlet axial nucleon form factor

\[ g^{(0)}_A (-q^2) = \frac{g^{(0)}_A}{(1 + q^2/M_A^2)^2}. \] (9)

The value of the form factor at zero momentum transfer \(g^{(0)}_A \approx 0.3\) can be obtained from the data on polarized deep inelastic scattering \(^{29}\). Various models, see, e.g., \(^{30}\) predict that the dipole mass parameter \(M_A\) is in the interval \(M_A \in [0.8, 1.1]\) GeV.

With the dipole parameterization the hyperfine potential in Eq. (8) has a simple analytic form

\[ V_{\text{hfs}}(r) = \frac{g^{(0)}_A \alpha}{m_Q} \frac{18\pi^4}{18N_f} \frac{e^{-rM_A}}{r^2} (2 - M_A r) (S \cdot s_N). \] (10)

We used \(\alpha = 17.2\) GeV\(^{-3}\) and hadrocharmonium wave functions from \(^{2,3}\) to calculate hyperfine splittings corresponding to different values of the dipole mass parameter \(M_A\) in the interval \([0.8, 1.1]\) GeV and collected the results in Table I. Taking into account approximations employed in the calculations the expected accuracy of the mass splitting estimate is
TABLE I. Hyperfine mass splitting between $J^P = 1/2^-$ and $3/2^-$ hadrocharmonium pentaquarks as function of the dipole mass parameter $M_A$

| $M_A$ [GeV] | 0.8 | 0.9 | 1.0 | 1.1 |
|------------|-----|-----|-----|-----|
| $\Delta E_{\text{hfs}}$ [MeV] | 21.1 | 27.7 | 34.9 | 42.5 |

around 30%. Comparison of the hyperfine splittings in Table I with the experimental splitting between pentaquarks $P_c(4457)$ and $P_c(4440)$ shows a satisfactory agreement between theory and experiment.

Let us turn to the $P_c(4440)$ and $P_c(4457)$ decay widths. Partial decay widths of the hadrocharmonium and molecular pentaquarks $P_c(4450)$ with $J^P = 3/2^-$ were calculated in [15]. We consider pentaquarks $P_c(4457)$ and $P_c(4440)$ as components of hadrocharmonium hyperfine doublet and use old results for the hadrocharmonium with $J^P = 3/2^-$. In the same formalism as in [15] we calculated now partial and total decay widths of the hadrocharmonium with $J^P = 1/2^-$. All partial and total widths of both components of the hyperfine hadrocharmonium doublet are collected in Table II. We see that decays to open charm of the $J^P = 1/2^-$ hadrocharmonium state are enhanced. This happens because the partial wave with $l = 0$ is allowed in these decays, to be compared with $l = 2$ allowed in decays of the $J^P = 3/2^-$ hadrocharmonium. The central potential that contributes to the $l = 0$ partial wave is stronger than the tensor potential that is responsible for the $l = 2$ partial wave, for more details see [15]. Additional accidental enhancement of $J^P = 1/2^-$ decays is due to the larger Clebsch-Gordon coefficients in this decay.

The theoretical uncertainties of the total widths in Table II are about 40%, they are compatible with the experimental widths in [1] at the level of two standard deviations.

Experimentally the total width of $P_c(4440)$ is roughly more than three times larger than the width of $P_c(4457)$. Comparing with the theoretical results in Table II we come to the conclusion that $P_c(4440)$ is a state with spin-parity $1/2^-$, while $P_c(4457)$ is a state with spin-parity $3/2^-$. The narrow LHCb pentaquark $P_c(4312)$ also finds a legitimate place in the hadrocharmonium scenario. We consider it as a bound state of the $\chi_{c0}(1P)$ charmonium state with $J^P = 0^+$ and the nucleon. The interaction potential between $\chi_{c0}(1P)$ and the nucleon is determined by the $\chi_{c0}(1P)$ chromoelectric polarizability. Polarizability is a symmetric

TABLE II. Decay widths of the spin doublet hadrocharmonium pentaquarks $P_c(4440)$ and $P_c(4457)$

| Decay mode | $\Gamma \left( \frac{1}{2}^- \right)^a_1$ | $\Gamma \left( \frac{3}{2}^- \right)^b_1$ |
|------------|-----------------|------------------|
| $P_c \rightarrow J/\psi N$ | 11 | 11 |
| $P_c \rightarrow \Lambda_c \bar{D}$ | 18.7 | 0.6 |
| $P_c \rightarrow \Sigma_c \bar{D}$ | 1.4 | 0.04 |
| $P_c \rightarrow \Lambda_c D^*$ | 13.7 | 4.2 |
| $P_c \rightarrow \Sigma_c^* \bar{D}$ | 0.004 | 0.4 |
| Total width | 44.8 | 16.2 |

$^a$ Decay width of $J = 1/2$ pentaquark in MeV.

$^b$ Decay width of $J = 3/2$ pentaquark in MeV.
two-index tensor $\alpha_{ik}$ and the effective $\chi_c(1P)N$ interaction is described by the effective Hamiltonian (see, e. g., [11])

$$H = -\frac{1}{2} \alpha_{ik} \langle N | E_i^a E_k^a | N \rangle,$$

(11)

where the nucleon matrix element in the coordinate space depends in the leading approximation on two structures $\delta_{ik}$ and $n_i n_k$

$$\langle N | E_i^a E_k^a | N \rangle = V_1(r) \delta_{ik} + V_2(r) n_i n_k,$$

(12)

where $r$ is the radius-vector from the center of the nucleon, and $n_r = r_i/r$.

The polarizability tensor $\alpha_{ik}$ is proportional to the Kronecker $\delta_{ik}$ for the $S$-state charmonium excitations, and is a linear combination of $\delta_{ik}$ and $L_i L_k$ ($L$ is the orbital momentum) for the $P$-states. It can be represented in the form

$$\alpha_{ik} = \alpha_1(J, S) \delta_{ik} + \alpha_2(J, S) J_i J_k,$$

(13)

where $S$ and $J$ are the charmonium state spin and total angular momentum, respectively.

Then the $\chi_c(1P)N$ interaction potential turns into a linear combination of a central and tensor potentials

$$H = V_c(r) + V_t \left[ (n \cdot J)(n \cdot J) - \frac{J^2}{3} \right].$$

(14)

For the estimates below we omit the potential $V_t(r)$, what can be justified by the instanton calculations. The potential $V_c(r)$ in Eq. (12) differs from the $\psi(2S)N$ interaction potential calculated in [2, 3] only by the value of the chromoelectric polarizability $\alpha = (1/3) \sum_i \alpha_{ii}$.

| $S$ | $J$ | $\alpha_1$ | $\alpha_2$ | $\alpha$ |
|-----|-----|----------|----------|--------|
| 0   | 1   | 105      | -78      | 53     |
| 1   | 2   | 79       | -13      | 53     |
| 1   | 1   | 27       | 39       | 53     |
| 1   | 0   | 53       | 0        | 53     |

- $S$ and $J$ are the charmonium state spin and total angular momentum, respectively.
- $\alpha = \Sigma_i \alpha_{ii}/3$ is one third of the trace of the polarizability tensor $\alpha_{ik}$.

Perturbative polarizabilities of the heavy Coulombic charmonium $P$-states can be calculated in QCD perturbation theory similarly to the $S$-state calculations in [17, 18]. The results of perturbative calculations are collected in Table III. Real charmonium is not a Coulombic bound state so results of the perturbative calculations should be taken with a grain of salt. We expect that ratios of perturbative polarizabilities are closer to the real world than their absolute values. The ratio of perturbative polarizabilities for $2S$ and $1P$ states is $\alpha(1P)/\alpha(2S) = 159/251 \approx 0.63$. The Schrödinger equation for $\chi_c(1P)$ and the nucleon has a bound state solution with the experimental mass of the LHCb pentaquark $P_c(4312)$ when the interaction potential is 0.58 times weaker than in [2, 3]. Taking into account that polarizabilities are not Coulombic we consider the substitution 0.63 → 0.58.
to be well inside the error bars of our calculations. Thus we identify the hadrocharmonium \( \chi_{c0}(1P)N \) bound state with the LHCb pentaquark \( P_c(4312) \), and predict that \( P_c(4312) \) has spin-parity \( 1/2^+ \). It does not have a hyperfine partner with approximately the same mass.

Let us discuss decays of the \( P_c(4312) \) hadrocharmonium. Its total width about 10 MeV can be easily explained as due to the decays of the weakly bound \( \chi_{c0}(1P) \) that has full width 10.8 MeV dominated by the decays into light hadrons. In addition \( P_c(4312) \) hadrocharmonium can decay into states with open charm. We expect that these decays are suppressed in comparison with such decays of the heavier pentaquarks (see Table IV) since the size of the hadrocharmonium \( P_c(4312) \) is larger due to the smaller binding energy about 42 MeV to be compared with about 170 MeV for heavier hadrocharmonium states. All this does not explain the decay \( P_c(4312) \to J/\psi + N \), where \( P_c(4312) \) was observed. The parities of \( \chi_{c0}(1P) \) and \( J/\psi \) are opposite so transitional polarizability \( \alpha(\chi_{c0}(1P) \to J/\psi) \) is zero and cannot explain this decay. The transition \( \chi_{c0}(1P) \to J/\psi \) could go through exchange by three gluons, at least it is allowed by quantum numbers. An estimate of the hadrocharmonium pentaquark decay \( P_c(4312) \to J/\psi + N \) is a challenging problem and we will not address it here.

### Table IV. Expected hadrocharmonium pentaquarks

| Constituents     | Binding energy [MeV] | Mass [MeV] | Spin-parity |
|------------------|----------------------|------------|-------------|
| \( \eta_c(2S)N \) | 176.1                | 4401       | 1/2^-       |
| \( \chi_{c1}(1P)N \) | 44.2                | 4406       | 3/2^+, 1/2^+ |
| \( h_c(1P)N \)   | 43.9                | 4421       | 1/2^+, 3/2^+ |
| \( \chi_{c2}(1P)N \) | 43.7                | 4452       | 5/2^+, 3/2^+ |

Hadrocharmonium interpretation of \( P_c(4312) \) as a bound state of \( \chi_{c0}(1P) \) and the nucleon naturally leads to the discussion of bound states of other charmonia \( 1P \) excitations and the nucleon. Trace of the polarizability tensor is one and the same for all \( 1P \) states, so the states \( \chi_{c1}(1P) \), \( \chi_{c2}(1P) \), and \( h_c(1P) \) should also form bound states with the nucleon. In addition the spin zero \( S \)-state \( \eta_c(2S) \) should form a hadrocharmonium bound state with the nucleon because its polarizability coincides with the polarizability of \( \psi(2S) \). Solutions of the bound state Schrödinger equations for all these states and their characteristics are collected in Table IV. Minor differences between the binding energies of different \( P \) states exceed the accuracy of our calculations and should be ignored.

We expect that degeneracy of the states with the same spin will be lifted by hyperfine interaction, and the magnitude of this splitting will be roughly the same as the splitting between \( P_c(4440) \) and \( P_c(4457) \). All charmonium constituents in Table IV except \( \eta_c(2S) \) have positive parity and natural widths about or below 1-2 MeV. We expect that decays of the type \( (\chi_{c2}(1P)N) \to \chi_{c1}(1P) + N \) will go due to nonzero transitional polarizabilities \( \alpha_{ik}(\chi_{c2}(1P) \to \chi_{c1}(1P)) \) and have partial widths at the level of 10-20 MeV. Decays of the hadrocharmonium states in Table IV to the states with open charm are also allowed and could have partial widths comparable with the ones for the decays to the states with hidden charm. Thus we expect that the interval of masses 4380-4430 MeV will be populated by a grid of hadrocharmonium states with the step 10-15 MeV and widths of order 10-30 MeV. We speculate that this set of states was interpreted in [4] as a wide pentaquark \( P_c(4380) \) and further experiment would resolve this structure in a series of relatively narrow overlapping resonances. Let us mention the \( (\chi_{c2}(1P)N) \) hadrocharmonium state in the last line in
Table IV has the mass that almost coincides with the masses of the LHCb pentaquarks \( P_{c}(4440) \) and \( P_{c}(4457) \) what somehow makes this scenario less transparent.

### III. SUMMARY

We discussed above the hadrocharmonium interpretation of the new LHCb pentaquark results. The pentaquarks \( P_{c}(4440) \) and \( P_{c}(4457) \) nicely fit prediction of almost degenerate hadrocharmonium pentaquarks with \( J^P = 1/2^-, 3/2^- \). We improved the estimate of hyperfine splitting between the color singlet hadrocharmonium states due to interference of the \( E1 \) and \( M2 \) multipoles in the QCD multipole expansion (see, e.g., [11]) and obtained a satisfactory quantitative agreement with the experimental data [1], see Table IV.

We calculated partial and total widths of loosely bound hadrocharmonium (\( \psi(2S)N \)) states with \( J^P = 1/2^-, 3/2^- \) (see Table II), and found that the total widths are compatible with the experimental data for \( P_{c}(4440) \) and \( P_{c}(4457) \). Comparing the theoretical and experimental ratios of total widths we conclude that \( P_{c}(4440) \) has spin-parity \( 1/2^- \) and \( P_{c}(4457) \) has spin-parity \( 3/2^- \).

The narrow LHCb pentaquark \( P_{c}(4312) \) is naturally interpreted as a \( (\chi_{c0}(1P)N) \) hadrocharmonium bound state with the binding energy 42 MeV, isospin \( I = 1/2 \), and spin-parity \( J^P = 1/2^+ \). Unlike the case of of almost degenerate hadrocharmonium \( (\psi(2S)N) \) bound states with spin parities \( J^P = 1/2^-, 3/2^- \), hadrocharmonium \( \chi_{c0}(1P)N \) does not have a partner with another spin. This happens because \( \chi_{c0}(1P) \) is a spin zero state. We expect that the hadrocharmonium isodoublet pentaquark \( P_{c}(4312) \) with \( J^P = 1/2^+ \) has width about 10-20 MeV that arises due to natural decay width of the \( \chi_{c0}(1P) \) charmonium and also due to open channels for decays into states with open charm.

We found a series of hadrocharmonium bound states with masses from 4380 MeV to 4430 MeV, widths about 10-30 MeV and known spin-parities, see Table IV. We speculate that these overlapping states were observed as a wide resonance \( P_{c}(4380) \), and expect that future experiments will find the complicated structure in the vicinity of 4380 MeV.

We would like to make a few remarks on the molecular interpretation of the LHCb pentaquarks \( P_{c}(4440) \), \( P_{c}(4457) \), and \( P_{c}(4312) \). This interpretation was suggested in [1], and elaborated in a number of recent papers [19–23]. The most straightforward argument in favor of the molecular nature of the new pentaquarks is their proximity to two-particle thresholds of charmed particles. Two more massive pentaquarks \( P_{c}(4440) \) and \( P_{c}(4457) \) are just below the \( \Sigma_{c}^{+}\bar{D}^{*0} \) threshold 4460 MeV, and \( P_{c}(4312) \) is just below the \( \Sigma_{c}^{+}\bar{D}^{0} \) threshold 4318 MeV. In the molecular picture pentaquarks are loosely bound states made from hadrons with open charm, \( \Sigma_{c}^{+}\bar{D}^{0} \) in the case of \( P_{c}(4312) \), and \( \Sigma_{c}^{+}\bar{D}^{*0} \) in the case of \( P_{c}(4457) \). Binding in the molecular scenario is due to meson exchanges, the most long range potential arises due to the the pion exchange. Parity conservation bans an effective \( \pi D^{0}\bar{D}^{0} \) vertex and hence, the pion exchange does not give any contribution in the binding potential in the \( \Sigma_{c}^{+}\bar{D}^{0} \) bound state. This observation makes interpretation of \( P_{c}(4312) \) as a \( \Sigma_{c}^{+}\bar{D}^{0} \) bound state a bit suspicious because it is hard to understand how exchanges by heavier mesons that generate short range potentials could be responsible for the existence of a loosely bound state with the constituents at relatively large distances, see also [24].

1 Recently the GlueX Collaboration reported nonobservation of the \( P_{c}(4450) \) pentaquark in the photoproduction reaction \( \gamma + p \rightarrow J/\psi + p \), what creates a certain tension between the LHCb and GlueX results, see, e.g., [31]. We expect that this nascent disagreement will be resolved in the near future.
does not provide any natural explanation for the emergence of two narrow closely separated $P_c(4440)$, $P_c(4457)$ pentaquarks.

Both the hadrocharmonium and molecular scenarios have their advantages and drawbacks as we discussed above. The hadrocharmonium approach elegantly describes (really predicts, see [2, 3]) small mass splitting between $P_c(4440)$ and $P_c(4457)$ as due to the QCD effective Hamiltonian Eq. (1). It also uniquely predicts spin-parities of $P_c(4312)$, $P_c(4440)$, and $P_c(4457)$ and their decay widths. According to the hadrocharmonium scenario the mass interval 4380-4430 MeV is densely populated by hadrocharmonium resonances with widths of order 10-30 MeV. Predictions in the hadrocharmonium approach have a certain rigidity, they can be experimentally confirmed or falsified. For example, if it would turn out that the spin-parities of $P_c(4440)$ and $P_c(4457)$ are not $1/2^-$ and $3/2^-$, one will be compelled to abandon their interpretation as hadrocharmonium bound states ($\psi(2S)N$). Molecular interpretation of pentaquarks is very flexible, due to the freedom to choose magnitudes of different coupling constants and parameters of numerous form factors it can accommodate almost any experimental data. This flexibility, that is advantageous in fitting the experimental data, deprives to a large extent the molecular approach of predictive power. For example, the molecular scenario can describe observed by the LHCb Collaboration two closely separated narrow resonances $P_c(4440)$ and $P_c(4457)$ [23], but it fails to give a natural explanation to their proximity to each other.

There is a significant number of experimentally verifiable predictions that are different in the hadrocharmonium and molecular scenarios. The quantum number assignments for the LHCb states do not coincide. For example, parity of $P_c(4312)$ is negative in the molecular picture [23] and it is positive in the hadrocharmonium picture. One more way to test both models is to consider the decay patterns. We have calculated partial and full decay widths in both pictures in [15] and obtained an intuitively appealing result that decays into states with hidden charm dominate for hadrocharmonium, while the molecule dominantly decays into states with open charm. Thus the decay patterns of molecular and hadrocharmonium pentaquarks are vastly different. We think that at the present stage the molecular and hadrocharmonium scenarios need further theoretical development, and hope that the dichotomy between them could be resolved by future experimental data.

ACKNOWLEDGMENTS

We are grateful to Valery Kubarovsky for very useful discussions. M. I. Eides and V. Yu. Petrov have been supported by the NSF grant PHY-1724638. M. V. Polyakov has been supported by the BMBF grant 05P2018.

[1] R. Aaij et al. [LHCb Collaboration], arXiv:1904.03947 [hep-ex].
[2] M. I. Eides, V. Y. Petrov, and M. V. Polyakov, Phys. Rev. D 93, no. 5, 054039 (2016), arXiv:1512.00426 [hep-ph].
[3] M. I. Eides, V. Y. Petrov, and M. V. Polyakov, Eur. Phys. J. C 78, no. 1, 36 (2018), arXiv:1709.09523 [hep-ph].
[4] R. Aaij et al. [LHCb Collaboration], Phys. Rev. Lett. 115, 072001 (2015), arXiv:1507.03414 [hep-ex].
[5] Y. R. Liu, H. X. Chen, W. Chen, X. Liu, and S. L. Zhu, arXiv:1903.11976 [hep-ph].
[6] S. Dubynskiy and M. B. Voloshin, Phys. Lett. B 666, 344 (2008), arXiv:0803.2224 [hep-ph].
[7] A. Sibirtsev and M. B. Voloshin, Phys. Rev. D 71, 076005 (2005), hep-ph/0502068.
[8] X. Li and M. B. Voloshin, Mod. Phys. Lett. A 29, no. 12, 1450060 (2014), arXiv:1309.1681 [hep-ph].
[9] S. J. Brodsky, I. A. Schmidt, and G. F. de Teramond, Phys. Rev. Lett. 64, 1011 (1990).
[10] M. E. Luke, A. V. Manohar, and M. J. Savage, Phys. Lett. B 288, 355 (1992), hep-ph/9204219.
[11] M. B. Voloshin, Prog. Part. Nucl. Phys. 61, 455 (2008), arXiv:0711.4556 [hep-ph].
[12] V. Kubarovsky and M. B. Voloshin, arXiv:1609.00050 [hep-ph].
[13] J. Ferretti, E. Santopinto, M. Naeem Anwar, and M. A. Bedolla, Phys. Lett. B 789, 562 (2019), arXiv:1807.01207 [hep-ph].
[14] I. A. Perevalova, M. V. Polyakov, and P. Schweitzer, Phys. Rev. D 94 (2016) no.5, 054024, arXiv:1607.07008 [hep-ph].
[15] M. I. Eides and V. Y. Petrov, Phys. Rev. D 98, no. 11, 114037 (2018), arXiv:1811.01691 [hep-ph].
[16] S. Coleman, Aspects of Symmetry, Cambridge University Press, Cambridge, 1985.
[17] M. E. Peskin, Nucl. Phys. B 156, 365 (1979).
[18] G. Bhanot and M. E. Peskin, Nucl. Phys. B 156, 391 (1979).
[19] F. K. Guo, H. J. Jing, U. G. Meißner, and S. Sakai, arXiv:1903.11503 [hep-ph].
[20] M. Z. Liu, Y. W. Pan, F. Z. Peng, M. Sánchez Sánchez, L. S. Geng, A. Hosaka, and M. Pavon Valderrama, arXiv:1903.11560 [hep-ph].
[21] H. X. Chen, W. Chen, and S. L. Zhu, arXiv:1903.11001 [hep-ph].
[22] J. He, arXiv:1903.11872 [hep-ph].
[23] R. Chen, Z. F. Sun, X. Liu and S. L. Zhu, arXiv:1903.11013 [hep-ph].
[24] C. Fernandez-Ramirez et al. [JPAC Collaboration], arXiv:1904.10021 [hep-ph].
[25] J. Balla, M. V. Polyakov, and C. Weiss, Nucl. Phys. B 510 (1998) 327 hep-ph/9707515.
[26] D. Diakonov and V. Y. Petrov, Nucl. Phys. B 245, 259 (1984). doi:10.1016/0550-3213(84)90432-2.
[27] D. Diakonov and V. Y. Petrov, Nucl. Phys. B 272, 457 (1986).
[28] P. L. Anthony et al. [E155 Collaboration], Phys. Lett. B 553, 18 (2003), hep-ex/0204028.
[29] D. de Florian, R. Sassot, M. Stratmann and W. Vogelsang, Phys. Rev. D 80, 034030 (2009), arXiv:0904.3821 [hep-ph].
[30] A. Silva, H. C. Kim, D. Urbano and K. Goeke, Phys. Rev. D 72, 094011 (2005), hep-ph/0509281.
[31] E. Chudakov on behalf of the GlueX Collaboration, talk at the 23d International Spin Symposium, Ferrara, September 2018.
[32] X. Cao and J. p. Dai, arXiv:1904.06015 [hep-ph].