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η_c-nucleus bound states

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ABSTRACT

η_c-nucleus bound state energies are calculated for various nuclei. Essential input for the calculations, namely the medium-modified D and D^* meson masses, as well as the density distributions in nuclei, are calculated within the quark-meson coupling (QMC) model. The attractive potentials for the η_c meson in the nuclear medium originate from the in-medium enhanced DD^* loops in the η_c self-energy. Our results suggest that the η_c meson should form bound states with all the nuclei considered.

1. Introduction

The study of the interactions of charmonium states, such as η_c and J/Ψ, with atomic nuclei offers opportunities to gain new insight into the properties of the strong force and strongly interacting matter [1,2]. Because charmonia and nucleons do not share light quarks, the Zweig rule suppresses interactions mediated by the exchange of mesons made of light quarks. It is therefore vital to explore other potential sources of attraction which could potentially lead to binding of charmonia to atomic nuclei.

A large body of work looking for alternatives to the meson-exchange paradigm has accumulated over the last three decades [3–6]. There are works based on the charmonium color polarizability [7,8], responsible for long-range van der Waals type of forces [9–16]. Others employ charmed meson loops, with light quarks created from the vacuum [13,17–20]. There are studies based on QCD sum rules [21–24] and phenomenological potentials [25,26]. More recently, there appeared lattice QCD simulations of the binding of charmonia to nuclear matter and finite nuclei [27], as well as light mesons and baryons [28]. The lattice QCD simulations of Ref. [27] have demonstrated that quarkonium-nucleus bound states exist for $A < 5$. Ref. [27] also infers a charmonium-nuclear matter binding energy $B^{NM} \sim 60$ MeV. However, these simulations have been performed at the flavor $SU(3)$-symmetric point, with unphysical pion masses, $m_\pi \sim 805$ MeV.

Model studies have suffered from scarce experimental data on the low-energy charmonium-nucleon interaction. However, this situation started to change for the J/Ψ case with the recent measurement, by the JLab GlueX Collaboration [29], of the $\gamma p \to J/\Psi p$ total cross section near threshold. It will further improve with the completion of other close-to-threshold J/Ψ photoproduction experiments at JLab [30,31]. Regarding production on nuclei, there is a JLab proposal to measure J/Ψ photoproduction off the deuteron [32]. On the other hand, unfortunately, there are not many experiments specially directed towards the η_c meson and its binding to nuclei, perhaps because it is more difficult to produce and detect. Recent studies on η_c production in heavy ion collisions (pp, pA, AA) at the LHC have been carried out in Refs. [33–37] towards the experimental study of its underlying production mechanisms. However, J/Ψ measurements are a first step towards to experimental study of binding of charmonia to nuclei. From the theory side, lattice QCD simulations of the free-space charmonium-nucleon interaction have become available within the last decade [38–42]. Unfortunately they have either been quenched or used large pion masses, which therefore require extrapolation to the physical mass [16].

Although crucial for constraining models, experimental knowledge of the free-space charmonium-nucleon interaction is not enough for assessing the likelihood of charmonium binding in nuclei. The overwhelming evidence that the internal structure of hadrons changes in medium [4,6,43,44] must be taken into account when addressing charmonium in nuclei. As shown in pre-

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vious studies [13,17–20], the effect of the nuclear mean fields on subthreshold $D\bar{D}$ states is of particular relevance. Those studies have revealed that modifications induced by the strong nuclear mean fields on the D mesons' light-quark content enhance the self-energy in such a way as to provide an attractive $J/\Psi-$nucleus effective potential. In the present paper we extend our previous study on the J/Ψ-nucleus bound states [17,19] to the case of η_c charmonium. η_c-nucleus bound states have also been predicted in other approaches [1,2,11,21,24], albeit with predictions for the binding energies varying over a wide range.

It is worth stressing that compared to the situation for the lighter ϕ meson [45–54], which couples strongly to above-threshold $K\bar{K}$ states, the charmonium states are expected to have a small width in medium and therefore the signal for the formation of such bound states may be experimentally cleaner.

This paper is organized as follows. In Sec. 2 we discuss the computation and present results for the mass shift of the η_c in symmetric nuclear matter. Using the results of Sec. 2, together with the density profiles of the nuclei calculated within the quark-meson coupling model, in Sec. 3 we present results for the scalar η_c-nucleus potentials, as well as the corresponding bound state energies. Finally, Sec. 4 is devoted to a summary and conclusions.

2. Calculation of the η_c scalar potential in symmetric nuclear matter

For the computation of the η_c scalar potential in nuclear matter we use an effective Lagrangian approach at the hadronic level [55], which is an $SU(4)$-flavor extension of light-flavor chiral-symmetric Lagrangians of pseudoscalar and vector mesons [56].

The extracted interaction Lagrangian density for the $\eta_c DD^*$ vertex is given by

$$\mathcal{L}_{\eta_c DD^*} = \mathrm{i}g_{\eta_c DD^*}(\partial_\mu \eta_c)\left[\bar{D}^{*\mu}\cdot D - \bar{D}\cdot D^{*\mu}\right] - \mathrm{i}g_{\eta_c DD^*}(\eta_c)\left[\bar{D}^{*\mu}\cdot(\partial_\mu D) - (\partial_\mu \bar{D})\cdot D^{*\mu}\right], \quad (1)$$

where $D^{(*)}$ represents the $D^{(*)}$-meson field isospin doublet, and $g_{\eta_c DD^*}$ is the coupling constant to be specified below.

We employ the effective interaction Lagrangian Eq. (1) to compute the η_c and self energy in vacuum and symmetric nuclear matter, following our previous works [17–20,50–54], and considering only they would be dominant DD^* loop. The η_c self-energy is thus given by

$$\Sigma_{\eta_c}(k^2) = \frac{8g^2_{\eta_c DD^*}}{\pi^2}\int_0^\infty \mathrm{d}k\, k^2 I(k^2) \quad (2)$$

for an η_c at rest, where

$$I(k^2) = \left.\frac{m^2_{\eta_c}(-1+k^{0\,2}/m^2_{D^*})}{(k^0+\omega_{D^*})(k^0-\omega_{D^*})(k^0-m_{\eta_c}-\omega_D)}\right|_{k^0=m_{\eta_c}-\omega_{D^*}} + \left.\frac{m^2_{\eta_c}(-1+k^{0\,2}/m^2_{D^*})}{(k^0-\omega_{D^*})(k^0-m_{\eta_c}+\omega_D)(k^0-m_{\eta_c}-\omega_D)}\right|_{k^0=-\omega_{D^*}}, \quad (3)$$

and $\omega_{D^{(*)}} = (k^2+m^2_{D^{(*)}})^{1/2}$, with $k=|\vec{k}|$. The integral in Eq. (2) is divergent and thus needs regularization. For this purpose we employ a phenomenological vertex form factor

$$u_{D^{(*)}}(k^2) = \left(\frac{\Lambda^2_{D^{(*)}}+m^2_{\eta_c}}{\Lambda^2_{D^{(*)}}+4\omega^2_{D^{(*)}}(k^2)}\right)^2, \quad (4)$$

with cutoff parameter $\Lambda_{D^{(*)}}$, as in Refs. [17–20,50–54]. Thus, to regularize Eq. (2) we will introduce the factor $u_D(k^2)u_{D^*}(k^2)$ into the integrand.

The cutoff parameter Λ_D (we use $\Lambda_{D^*}=\Lambda_D$) is an unknown input to our calculation. However, it may be determined phenomenologically using, for example, a quark model. In fact, in Ref. [17] its value has been estimated to be $\Lambda \approx 2500$ MeV, and it serves us as a reasonable guide to quantify the sensitivity of our results to its value. Therefore we vary it over the interval 1500-3000 MeV.

Because $SU(4)$ flavor symmetry is strongly broken in Nature, we use experimental values for the meson masses [57] and empirically known values for the coupling constants, as explained below. For the D meson mass, we take the averaged masses of the neutral and charged states, and similarly for the D^*. Thus $m_D = 1867.2$ MeV and $m_{D^*} = 2008.6$ MeV. For the coupling constants, $g_{\eta_c DD^*} = 0.60\, g_{\psi DD}$ was obtained in Ref. [58] as the residue at the poles of suitable form factors using a dispersion formulation of the relativistic constituent quark model, where $g_{\psi DD} = 7.64$ was estimated in Ref. [59] using the vector meson dominance model and isospin symmetry. In this study we use the coupling constant, $g_{\eta_c DD^*} = (0.60/\sqrt{2})\, g_{\psi DD} \simeq 0.424\, g_{\psi DD}$, where the factor $(1/\sqrt{2})$ is introduced to give a larger $SU(4)$ symmetry breaking effect than Ref. [58]. In this connection we mention that recent investigations of $SU(4)$ flavor symmetry breaking in hadron couplings of charmed hadrons to light mesons are not conclusive; while two studies based on Schwinger-Dyson equations of QCD find large deviations from $SU(4)$ symmetry [60,61], studies using QCD sum rules [62,63], a constituent quark model [64] and a holographic QCD model [65] find moderate deviations.

We are interested in the difference between the in-medium, $m^*_{\eta_c}$, and vacuum, m_{η_c}, masses of the η_c,

$$\Delta m_{\eta_c} = m^*_{\eta_c} - m_{\eta_c}, \quad (5)$$

with the masses obtained self-consistently from

$$m^2_{\eta_c} = (m^0_{\eta_c})^2 + \Sigma_{\eta_c}(k^2 = m^2_{\eta_c}), \quad (6)$$

where $m^0_{\eta_c}$ is the bare η_c mass and $\Sigma_{\eta_c}(k^2)$ is given by Eq. (2). The Λ_D-dependent η_c-meson bare mass, $m^0_{\eta_c}$, is fixed by fitting the physical η_c-meson mass, $m_{\eta_c} = 2983.9$ MeV.

The in-medium η_c mass is obtained in a similar way, with the self-energy calculated with the medium-modified D and D^* masses. The nuclear density dependence of the η_c-meson mass is determined by the intermediate-state D and D^* meson interactions with the nuclear medium through their medium-modified masses. The in-medium masses m^*_D and $m^*_{D^*}$ are calculated within the quark-meson coupling (QMC) model [17,18], in which effective scalar and vector meson mean fields couple to the light u and d quarks in the charmed mesons [17,18]. The QMC model has proven to be very successful in studying the properties of hadrons in nuclear matter and finite nuclei [66–70]. This model considers infinitely large, uniformly symmetric, spin-isospin-saturated nuclear matter in its rest frame, where all the scalar and vector mean field potentials, which are responsible for the nuclear many-body interactions, become constant in the Hartree approximation [66,69,70].

In Fig. 1 we present the resulting medium-modified masses for the D and D^* mesons, calculated within the QMC model [17], as a function of ρ_B/ρ_0, where ρ_B is the baryon density of nuclear matter and $\rho_0 = 0.15$ fm^{-3} is the saturation density of symmetric nuclear matter. The net reductions in the masses of the D and D^* mesons are nearly the same as a function of density, with each decreasing by around 60 MeV at ρ_0.

The behavior of the D meson mass in medium (finite density and/or temperature) has been studied in a variety of approaches.

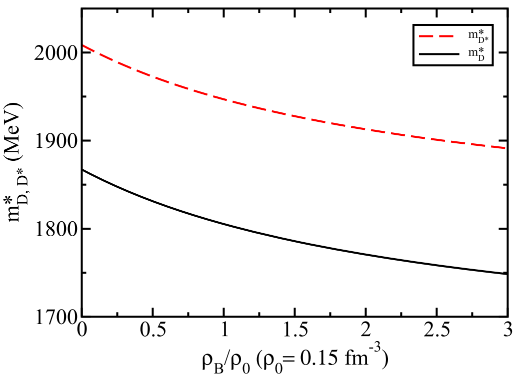

Fig. 1. In-medium D and D^* meson masses calculated within the QMC model. Adapted from Ref. [17].

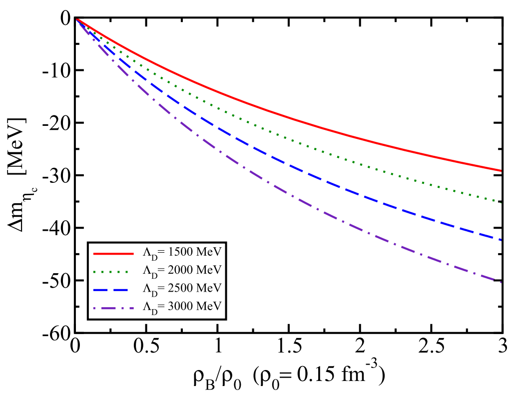

Fig. 2. η_c mass shift as a function of the nuclear matter density for various values of the cutoff parameter.

Some of these [71–73] find a decreasing D meson mass at finite baryon density, while others [74–78], interestingly, find the opposite behavior. However, it is important to note that none of the studies in nuclear matter are constrained by the saturation properties of nuclear matter, although it is constrained in the case of the present work. Furthermore, some of these works employ a non relativistic approach and relativistic effects might be important.

In Fig. 2, we present the η_c-meson mass shift, Δm_{η_c}, as a function of the nuclear matter density, ρ_B, normalized to ρ_0, for four values of the cutoff parameter Λ_D. As can be seen from the figure, the effect of the in-medium D and D^* mass change is to shift the η_c mass downwards. This is because the reduction in the D and D^* masses enhances the DD^*-loop contribution in nuclear matter relative to that in vacuum. This effect increases the larger the cutoff mass Λ_D.

The results described above support a small downward mass shift for the η_c in nuclear matter and open the possibility to study the binding of η_c mesons to nuclei, to which we turn our attention in the next section.

3. η_c-nucleus bound states

We now discuss the situation where the η_c-meson is produced inside a nucleus A with baryon density distribution $\rho_B^A(r)$. The nuclei we consider here are ^{4}He, ^{12}C, ^{16}O, ^{40}Ca, ^{48}Ca, ^{90}Zr, ^{197}Au, and ^{208}Pb. Their nuclear density distributions are also calculated within the QMC model, except for ^{4}He, whose parametrization was obtained in Ref. [79]. Using a local density approximation, the η_c-meson potential within nucleus A is given by

$$V_{\eta_c A}(r) = \Delta m_{\eta_c}(\rho_B^A(r)), \tag{7}$$

Table 1
η_c-nucleus bound state energies for different values of the cutoff parameter Λ_D. All dimensionful quantities are in MeV.

	$n\ell$	Bound state energies			
		$\Lambda_D = 1500$	$\Lambda_D = 2000$	$\Lambda_D = 2500$	$\Lambda_D = 3000$
$^{4}_{\eta_c}$He	1s	−1.49	−3.11	−5.49	−8.55
$^{12}_{\eta_c}$C	1s	−5.91	−8.27	−11.28	−14.79
	1p	−0.28	−1.63	−3.69	−6.33
$^{16}_{\eta_c}$O	1s	−7.35	−9.92	−13.15	−16.87
	1p	−1.94	−3.87	−6.48	−9.63
$^{40}_{\eta_c}$Ca	1s	−11.26	−14.42	−18.31	−22.73
	1p	−7.19	−10.02	−13.59	−17.70
	1d	−2.82	−5.22	−8.36	−12.09
	2s	−2.36	−4.51	−7.44	−10.98
$^{48}_{\eta_c}$Ca	1s	−11.37	−14.46	−18.26	−22.58
	1p	−7.83	−10.68	−14.23	−18.32
	1d	−3.88	−6.40	−9.63	−13.41
	2s	−3.15	−5.47	−8.54	−12.17
$^{90}_{\eta_c}$Zr	1s	−12.26	−15.35	−19.14	−23.43
	1p	−9.88	−12.86	−16.53	−20.70
	1d	−7.05	−9.87	−13.38	−17.40
	2s	−6.14	−8.87	−12.29	−16.24
	1f	−3.90	−6.50	−9.81	−13.65
$^{197}_{\eta_c}$Au	1s	−12.57	−15.59	−19.26	−23.41
	1p	−11.17	−14.14	−17.77	−21.87
	1d	−9.42	−12.31	−15.87	−19.90
	2s	−8.69	−11.53	−15.04	−19.02
	1f	−7.39	−10.19	−13.70	−17.61
$^{208}_{\eta_c}$Pb	1s	−12.99	−16.09	−19.82	−24.12
	1p	−11.60	−14.64	−18.37	−22.59
	1d	−9.86	−12.83	−16.49	−20.63
	2s	−9.16	−12.09	−15.70	−19.80
	1f	−7.85	−10.74	−14.30	−18.37

where r is the distance from the center of the nucleus.

In Fig. 3 we present the η_c-meson potentials for a selection of the nuclei mentioned above and various values of the cutoff parameter Λ_D. From the figure one can see that the η_c potential in nuclei is attractive in all cases but its depth depends on the value of the cutoff parameter, being deeper the larger Λ_D is. For example, it varies, from −18 MeV to −32 MeV for ^{4}He and from −15 MeV to −26 MeV for ^{208}Pb, when the cutoff varies from 1500 MeV to 3000 MeV. This dependence is, indeed, an uncertainty in the results obtained in our approach.

Using the η_c-meson potentials obtained in this manner, we next calculate the η_c-meson–nucleus bound state energies for the nuclei listed above by solving the Klein-Gordon equation

$$\left(-\nabla^2 + \mu^2 + 2\mu V(\vec{r})\right)\phi_{\eta_c}(\vec{r}) = \mathcal{E}^2\phi_{\eta_c}(\vec{r}), \tag{8}$$

where $\mu = m_{\eta_c}m_A/(m_{\eta_c}+m_A)$ is the reduced mass of the η_c-meson-nucleus system with m_{η_c} (m_A) the mass of the η_c-meson (nucleus A) in vacuum, and $V(\vec{r})$ is the η_c-meson-nucleus potential given in Eq. (7).

The bound state energies (E) of the η_c-nucleus system, given by $E = \mathcal{E} - \mu$, where $\mathcal{E}$ is the energy eigenvalue in Eq. (8), are calculated for four values of the cutoff parameter Λ_D and are listed in Table 1. These results show that the η_c-meson is expected to form bound states with all the nuclei studied and this prediction is independent of the value of the cutoff parameter Λ_D. However, the particular values for the bound state energies are clearly dependent on Λ_D, namely, each of them increases in absolute value as Λ_D increases. This was expected from the behavior of the η_c potentials, since these are deeper for larger values of the cutoff parameter. Note also that the η_c binds more strongly to heavier nuclei. We have also solved the Schröedinger equation with the

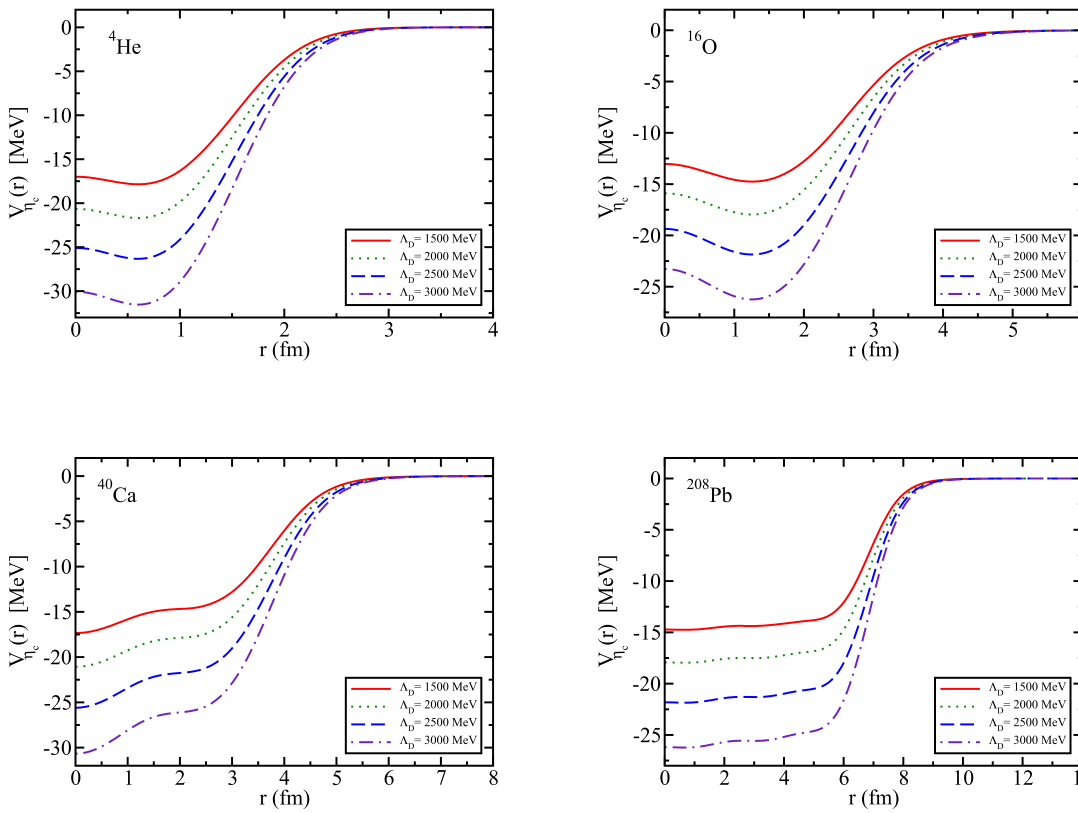

Fig. 3. η_c-nucleus potentials for various nuclei and values of the cutoff parameter Λ_D.

potential Eq. (7) to obtain the single-particle energies [6] and compared these with those given in Table 1. The results found in both cases are essentially the same.

Note that we have ignored the natural width of 32 MeV in free space of the η_c but understand this could be an issue related to the observability of the predicted bound states. Furthermore, we have no reason to believe the width will be suppressed in medium. Thus, even though it could be difficult to resolve the individual states, it should be possible to see that there are bound states which is the main point of this work. It remains to be seen by how much the inclusion of a repulsive imaginary part will affect the predicted bound states. We believe this can be done in future work.

Another effect that could potentially have important consequences for the formation of the bound states presented here, since it is repulsive, is the Ericson-Ericson-Lorentz-Lorenz (EELL) double scattering correction. However, we estimate that this effect, even though it may play an important role for the light isoscalar η meson [80], is much reduced in the present case. This is because in the QMC model we work at the Hartree level and ignore the effect of nucleon correlations, or nonlocal interactions. Furthermore the EELL effect was aimed at low energy pion scattering with the assumption $qr \to 0$. For heavy mesons like the η_c the momenta are much higher and the inverse correlation length $< 1/r >$ that appears in the EELL effect will certainly be much reduced by cancellations associated with the oscillatory behavior of the exponential. This plus the fact that we only work at the Hartree level and ignore exchange corrections that appear in the Hartree-Fock-based treatment.

4. Summary and discussion

We have calculated the spectra of η_c-nucleus bound states for various finite nuclei. The meson-nucleus potentials were calculated using a local density approximation, with the inclusion of the DD^* meson loop in the η_c self-energy. The nuclear density distributions, as well as the in-medium D and D^* meson masses were consistently calculated by employing the quark-meson coupling model. Using the meson potentials in nuclei, we have solved the Klein-Gordon equation and obtained meson–nucleus bound state energies. The sensitivity of our results to the cutoff parameter Λ_D used in the vertex form factors appearing in the η_c self-energy has also been explored. Interestingly, the η_c-nucleus bound state energies calculated here are larger than the corresponding J/Ψ energies calculated in Ref. [18], by some of us, using the same approach. Needless to say, this deserves further investigation.

Our results show that one should expect the η_c to form bound states for all the nuclei studied, even though the precise values of the bound state energies are dependent on the cutoff mass values used in the form factors. The discovery of such bound states would represent an important step forward in our understanding of the nature of strongly interacting systems.

Declaration of competing interest

There are no financial interests/personal relationships which may be considered as potential competing interests.

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References

[1] S.J. Brodsky, I. Schmidt, G. de Teramond, Phys. Rev. Lett. 64 (1990) 1011.
[2] D. Wasson, Phys. Rev. Lett. 67 (1991) 2237.
[3] A. Hosaka, T. Hyodo, K. Sudoh, Y. Yamaguchi, S. Yasui, Prog. Part. Nucl. Phys. 96 (2017) 88–153.
[4] G. Krein, AIP Conf. Proc. 1701 (1) (2016) 020012.
[5] V. Metag, M. Nanova, E.Y. Paryev, Prog. Part. Nucl. Phys. 97 (2017) 199–260.
[6] G. Krein, A.W. Thomas, K. Tsushima, Prog. Part. Nucl. Phys. 100 (2018) 161.
[7] M.E. Peskin, Nucl. Phys. B 156 (1979) 365–390.
[8] D. Kharzeev, Proc. Int. Sch. Phys. Enrico Fermi 130 (1996) 105.
[9] A. Kaidalov, P. Volkovitsky, Phys. Rev. Lett. 69 (1992) 3155–3156.
[10] M.E. Luke, A.V. Manohar, M.J. Savage, Phys. Lett. B 288 (1992) 355–359.
[11] G.F. de Teramond, R. Espinoza, M. Ortega-Rodriguez, Phys. Rev. D 58 (1998) 034012.
[12] S.J. Brodsky, G.A. Miller, Phys. Lett. B 412 (1997) 125–130.
[13] S.H. Lee, C. Ko, Phys. Rev. C 67 (2003) 038202.
[14] A. Sibirtsev, M. Voloshin, Phys. Rev. D 71 (2005) 076005.
[15] M. Voloshin, Prog. Part. Nucl. Phys. 61 (2008) 455–511.
[16] J. Tarrús Castellà, G. Krein, Phys. Rev. D 98 (1) (2018) 014029.
[17] G. Krein, A.W. Thomas, K. Tsushima, Phys. Lett. B 697 (2011) 136–141.
[18] K. Tsushima, D. Lu, G. Krein, A.W. Thomas, Phys. Rev. C 83 (2011) 065208.
[19] K. Tsushima, D. Lu, G. Krein, A.W. Thomas, AIP Conf. Proc. 1354 (1) (2011) 39–44.
[20] G. Krein, J. Phys. Conf. Ser. 422 (2013) 012012.
[21] F. Klingl, S.s. Kim, S.H. Lee, P. Morath, W. Weise, Phys. Rev. Lett. 82 (1999) 3396, Erratum: Phys. Rev. Lett. 83 (1999) 4224.
[22] A. Hayashigaki, Prog. Theor. Phys. 101 (1999) 923–935.
[23] S.s. Kim, S.H. Lee, Nucl. Phys. A 679 (2001) 517–548.
[24] A. Kumar, A. Mishra, Phys. Rev. C 82 (2010) 045207.
[25] V. Belyaev, N. Shevchenko, A. Fix, W. Sandhas, Nucl. Phys. A 780 (2006) 100–111.
[26] A. Yokota, E. Hiyama, M. Oka, PTEP 2013 (11) (2013) 113D01.
[27] S. Beane, E. Chang, S. Cohen, W. Detmold, H.W. Lin, K. Orginos, A. Parreño, M. Savage, Phys. Rev. D 91 (11) (2015) 114503.
[28] M. Alberti, G.S. Bali, S. Collins, F. Knechtli, G. Moir, W. Söldner, Phys. Rev. D 95 (7) (2017) 074501.
[29] A. Ali, et al., GlueX, Phys. Rev. Lett. 123 (7) (2019) 072001.
[30] S. Stepanyan, M. Battaglieri, A. Celentano, R. De Vita, V. Kubarovsky, Near Threshold J/ψ Photoproduction and Study of LHCb Pentaquarks with CLAS12, JLab E12-12-001A, Newport News, VA, USA, 2012.
[31] Z.E. Meziani, et al., A Search for the LHCb Charmed Pentaquark using Photo-Production of J/ψ at Threshold in Hall C at Jefferson Lab, arXiv:1609.00676 [hep-ex].
[32] Y. Ilieva, B. McKinnon, P. Nadel-Turonski, V. Kubarovsky, S. Stepanyan, Z.W. Zhao, Study of J/Ψ Photoproduction off Deuteron, JLab E12-11-003B, Newport News, VA, USA, 2011.
[33] R. Aaij, et al., LHCb, Eur. Phys. J. C 80 (3) (2020) 191, https://doi.org/10.1140/epjc/s10052-020-7733-0, arXiv:1911.03326 [hep-ex].
[34] Tichouk, H. Sun, X. Luo, Phys. Rev. D 101 (9) (2020) 094006.
[35] Tichouk, H. Sun, X. Luo, Phys. Rev. D 101 (5) (2020) 054035.
[36] V.P. Gonçalves, B.D. Moreira, Phys. Rev. D 97 (9) (2018) 094009.
[37] S.R. Klein, Phys. Rev. D 98 (11) (2018) 118501, arXiv:1808.08253 [hep-ph].
[38] K. Yokokawa, S. Sasaki, T. Hatsuda, A. Hayashigaki, Phys. Rev. D 74 (2006) 034504.
[39] L. Liu, H.W. Lin, K. Orginos, PoS LATTICE2008 (2008) 112.
[40] T. Kawanai, S. Sasaki, Phys. Rev. D 82 (2010) 091501.
[41] T. Kawanai, S. Sasaki, PoS LATTICE2010 (2010) 156.
[42] U. Skerbis, S. Prelovsek, Phys. Rev. D 99 (9) (2019) 094505.
[43] R.S. Hayano, T. Hatsuda, Rev. Mod. Phys. 82 (2010) 2949.
[44] S. Leupold, V. Metag, U. Mosel, Int. J. Mod. Phys. E 19 (2010) 147.
[45] C.M. Ko, P. Levai, X.J. Qiu, C.T. Li, Phys. Rev. C 45 (1992) 1400–1402.
[46] F. Klingl, T. Waas, W. Weise, Phys. Lett. B 431 (1998) 254–262.
[47] E. Oset, A. Ramos, Nucl. Phys. A 679 (2001) 616–628.
[48] D. Cabrera, M.J. Vicente Vacas, Phys. Rev. C 67 (2003) 045203.
[49] P. Gubler, W. Weise, Phys. Lett. B 751 (2015) 396–401.
[50] J.J. Cobos-Martínez, K. Tsushima, G. Krein, A.W. Thomas, Phys. Lett. B 771 (2017) 113–118.
[51] J.J. Cobos-Martínez, K. Tsushima, G. Krein, A.W. Thomas, Phys. Rev. C 96 (3) (2017) 035201.
[52] J.J. Cobos-Martínez, K. Tsushima, G. Krein, A.W. Thomas, J. Phys. Conf. Ser. 912 (1) (2017) 012009.
[53] J.J. Cobos-Martínez, K. Tsushima, G. Krein, A.W. Thomas, PoS Hadron2017 (2018) 209.
[54] J.J. Cobos-Martínez, K. Tsushima, G. Krein, A.W. Thomas, JPS Conf. Proc. 26 (2019) 024033.
[55] F. Klingl, N. Kaiser, W. Weise, Z. Phys. A 356 (1996) 193.
[56] Z.w. Lin, C.M. Ko, B. Zhang, Phys. Rev. C 61 (2000) 024904.
[57] M. Tanabashi, et al., Particle Data Group, Phys. Rev. D 98 (3) (2018) 030001.
[58] W. Lucha, D. Melikhov, H. Sazdjian, S. Simula, Phys. Rev. D 93 (1) (2016) 016004, Addendum: Phys. Rev. D 93 (1) (2016) 019902.
[59] Z.w. Lin, C.M. Ko, Phys. Rev. C 62 (2000) 034903.
[60] B. El-Bennich, G. Krein, L. Chang, C.D. Roberts, D.J. Wilson, Phys. Rev. D 85 (2012) 031502.
[61] B. El-Bennich, M.A. Paracha, C.D. Roberts, E. Rojas, Phys. Rev. D 95 (3) (2017) 034037.
[62] F.S. Navarra, M. Nielsen, Phys. Lett. B 443 (1998) 285.
[63] A. Khodjamirian, C. Klein, T. Mannel, Y.-M. Wang, J. High Energy Phys. 1109 (2011) 106.
[64] C.E. Fontoura, J. Haidenbauer, G. Krein, Eur. Phys. J. A 53 (5) (2017) 92.
[65] A. Ballon-Bayona, G. Krein, C. Miller, Phys. Rev. D 96 (1) (2017) 014017.
[66] K. Saito, K. Tsushima, A.W. Thomas, Prog. Part. Nucl. Phys. 58 (2007) 1.
[67] P.A.M. Guichon, H.H. Matevosyan, N. Sandulescu, A.W. Thomas, Nucl. Phys. A 772 (2006) 1–19.
[68] J.R. Stone, P.A.M. Guichon, P.G. Reinhard, A.W. Thomas, Phys. Rev. Lett. 116 (9) (2016) 092501.
[69] K. Tsushima, K. Saito, A.W. Thomas, S.V. Wright, Phys. Lett. B 429 (1998) 239, Erratum: Phys. Lett. B 436 (1998) 453.
[70] P.A.M. Guichon, Nucl. Phys. A 497 (1989) 265C.
[71] A. Hayashigaki, Phys. Lett. B 487 (2000) 96–103.
[72] K. Azizi, N. Er, H. Sundu, Eur. Phys. J. C 74 (2014) 3021.
[73] Z.G. Wang, Phys. Rev. C 92 (6) (2015) 065205.
[74] K. Suzuki, P. Gubler, M. Oka, Phys. Rev. C 93 (4) (2016) 045209.
[75] A. Park, P. Gubler, M. Harada, S.H. Lee, C. Nonaka, W. Park, Phys. Rev. D 93 (5) (2016) 054035.
[76] P. Gubler, T. Song, S.H. Lee, Phys. Rev. D 101 (11) (2020) 114029.
[77] T. Hilger, R. Thomas, B. Kampfer, Phys. Rev. C 79 (2009) 025202.
[78] T.F. Caramés, C.E. Fontoura, G. Krein, K. Tsushima, J. Vijande, A. Valcarce, Phys. Rev. D 94 (3) (2016) 034009.
[79] K. Saito, K. Tsushima, A.W. Thomas, Phys. Rev. C 56 (1997) 566.
[80] S.D. Bass, A.W. Thomas, Phys. Lett. B 634 (2006) 368–373.