

Charmonium in nuclear matter and nuclei

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We present results for the η_c -nucleus bound state energies for various nuclei using an effective Lagrangians approach. The attractive potentials for the η_c in the nuclear medium originate from the medium-modified intermediate DD^* state in the η_c self energy, using the local density approximation. Our results suggest that the η_c should form bound states with all the nuclei considered.

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1. Introduction

Quantum chromodynamics (QCD) is the accepted fundamental theory of the strong interactions. However, a quantitative understanding of the strong force and strongly interacting matter in vacuum and extreme conditions of temperature and density from the underlying theory is still limited. Even though there are numerous reasons to study charmonium states, the study of the interactions of these states with atomic nuclei offers potential opportunities to gain an understanding of the workings of strongly interacting matter and QCD. Since charmonium and nucleons do not share light (u , d) quarks, the OZI rule suppresses the interactions mediated by the exchange of mesons made of only light quarks. If such states are indeed bound to nuclei, it is therefore important to search for other sources of attraction which could lead to the binding of charmonia to atomic nuclei. Since the suggestion of Brodsky [1], more than three decades ago, that charmonium states may be bound with nuclei, a large amount of research looking for alternatives to the meson exchange mechanism has accumulated over the years to investigate the possible existence of such exotic states [2]. The discovery of such bound states would represent an important step forward in our understanding of the nature of strongly interacting systems.

There is a great amount of evidence that the internal structure of hadrons changes in medium and this must be taken into account when addressing charmonium in nuclei. Studies carried by some of us [3, 4] have shown that the effect of the nuclear mean fields on the intermedium $D\bar{D}$ state is of particular importance in the J/Ψ case; the 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/Ψ -nucleus effective potential. We extended this approach to the case of η_c in Ref. [5] by considering an intermediate DD^* state in the η_c self-energy with the light quarks created from the vacuum. Here the effective scalar and vector meson mean fields in the nuclear medium couple to the light u and d quarks in the charmed mesons and this, as in the case of the J/Ψ , provides attraction to the η_c in the nuclear medium. Recently, we have extended this approach to the case of bottomonium [6].

This paper is organized as follows. In Sec. 2 we briefly 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 studied, in Sec. 3 we present results for the scalar η_c -nucleus potentials, as well as the corresponding bound state energies. Finally, in Sec. 4 we present our conclusions.

2. η_c in nuclear matter

In computing the η_c mass shift in nuclear matter we take into account only the intermediate DD^* state contribution to the η_c self-energy; see Refs [5, 6] and references therein for details. The effective interaction Lagrangian for the $\eta_c DD^*$ vertex is given by

$$\mathcal{L}_{\eta_c DD^*} = ig_{\eta_c DD^*} \left[(\partial^\mu \eta_c) \left(\bar{D}_\mu^* D - \bar{D} D_\mu^* \right) - \eta_c \left(\bar{D}_\mu^* (\partial^\mu D) - (\partial^\mu \bar{D}) D_\mu^* \right) \right], \quad (1)$$

where $D^{(*)}$ represents the $D^{(*)}$ -meson field isospin doublet, and $g_{\eta_c DD^*}$ is the coupling constant. Using Eq. (1), the Feynman diagram of Fig. 1, and considering an η_c at rest, the η_c self-energy is

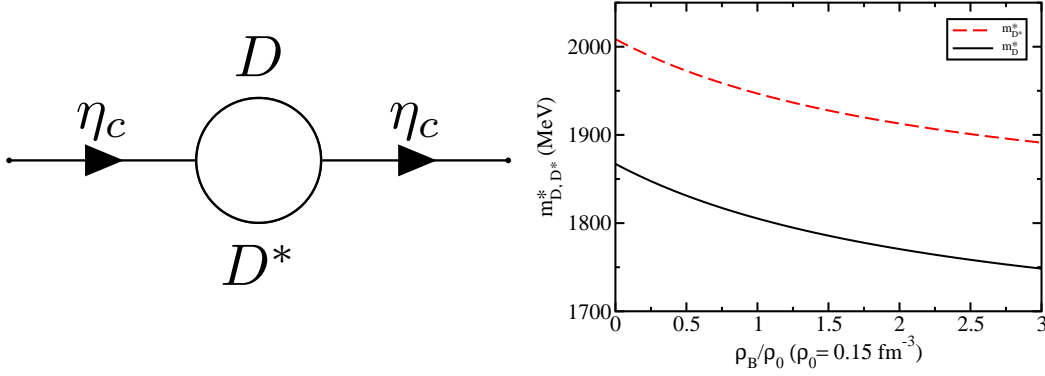


Figure 1: Left panel: Intermediate DD^* state contribution to the η_c self energy. Right panel: In-medium D and D^* meson masses in nuclear matter.

given by,

$$\Sigma_{\eta_c}(k^2) = \frac{8g_{\eta_c DD^*}^2}{\pi^2} \int_0^\infty dk k^2 I(k^2) \quad (2)$$

where $I(k^2)$ is given by

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

and $\omega_{D^{(*)}} = (k^2 + m_{D^{(*)}}^2)^{1/2}$, with $k = |\vec{k}|$. The integral in Eq. (2) is divergent and will be regulated with a phenomenological vertex form factor, as in Refs [3–8],

$$u_{D^{(*)}}(k^2) = \left(\frac{\Lambda_D^2 + m_{\eta_c}^2}{\Lambda_D^2 + 4\omega_{D^{(*)}}^2(k^2)} \right)^2, \quad (4)$$

with cutoff parameter $\Lambda_D (= \Lambda_{D^*})$, by introducing the factor $u_D(k^2)u_{D^*}(k^2)$ into the integrand of Eq. (2). The cutoff parameter Λ_D is an unknown input to our calculation. Its value has been determined phenomenologically in Ref [3] to be $\Lambda_D \approx 2500$ MeV. However, in order to quantify the uncertainty of our results to its value we present results for Λ_D in the range 1500-3000 MeV.

The η_c mass shift in nuclear matter is computed from 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 computed by solving

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

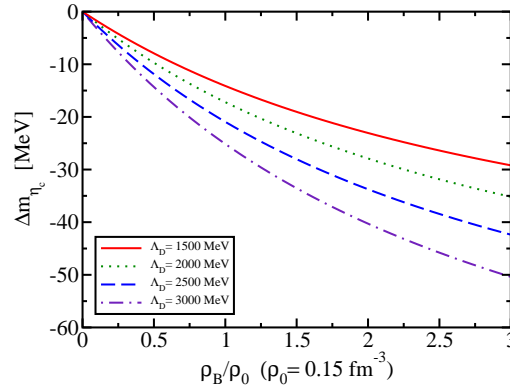


Figure 2: η_c mass shift as a function of the nuclear matter density for various values of the cutoff parameter.

where $m_{\eta_c}^0$ is the bare η_c mass and $\Sigma_{\eta_c}(k^2)$ is given by Eq. (2). The η_c meson bare mass is fixed using Eq. (6) by fitting the physical η_c mass, $m_{\eta_c} = 2983.9$ MeV. (For the other parameters in vacuum we use $g_{\eta_c DD^*} = (0.6/\sqrt{2})g_{J/\psi DD} = 3.24$ [9, 10], $m_D = 1867.2$ MeV, and $m_{D^*} = 2008.6$ MeV; see Ref. [5] for details). The η_c mass in the nuclear medium is similarly computed from Eq. (6) with the self-energy calculated with the medium-modified D and D^* masses. The nuclear density dependence of the η_c mass originates from the interactions of the intermediate DD^* state with the nuclear medium through their medium-modified masses, m_D^* and $m_{D^*}^*$, respectively. These in-medium masses are calculated within the quark-meson coupling (QMC) model [3, 4], in which effective scalar and vector meson mean fields couple to the light quarks in the charmed mesons [3, 4]. The resulting medium-modified masses for the D and D^* mesons are shown in the left panel of Fig. 1 as a function of ρ_B/ρ_0 , where ρ_B is the baryon density of nuclear matter and $\rho_0 = 0.15 \text{ fm}^{-3}$ is the saturation density of symmetric nuclear matter. As can be seen from Fig. 1 the masses of the D and D^* mesons are reduced in the nuclear medium. In Fig. 2, we present the η_c mass shift, Δm_{η_c} , as a function of ρ_B/ρ_0 for several values of Λ_D . As can be seen from Fig. 2, the effect of the nuclear medium is to shift the η_c mass downwards. This happens for all values of Λ_D .

3. η_c in nuclei

A negative mass shift for the η_c means that the nuclear medium provides attraction to the η_c and opens the possibility for the binding of η_c mesons to nuclei. Thus, we now consider the binding of the η_c to nuclei when the η_c is produced inside a nucleus A with baryon density distribution $\rho_B^A(r)$; see Ref. [5] for a more detailed discussion. Here we only consider the nuclei ${}^4\text{He}$, ${}^{12}\text{C}$, ${}^{16}\text{O}$, ${}^{197}\text{Au}$, and ${}^{208}\text{Pb}$, whose baryon density distributions were also calculated within the QMC model [5], except that for ${}^4\text{He}$ [11]. Using a local density approximation, the η_c -meson potential inside nucleus A is given by

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

where Δm_{η_c} mass shift shown in Fig. 2, and r is the distance from the center of the nucleus. In Fig. 3 we present the η_c potentials for the nuclei mentioned above and various values of Λ_D . From Fig. 3

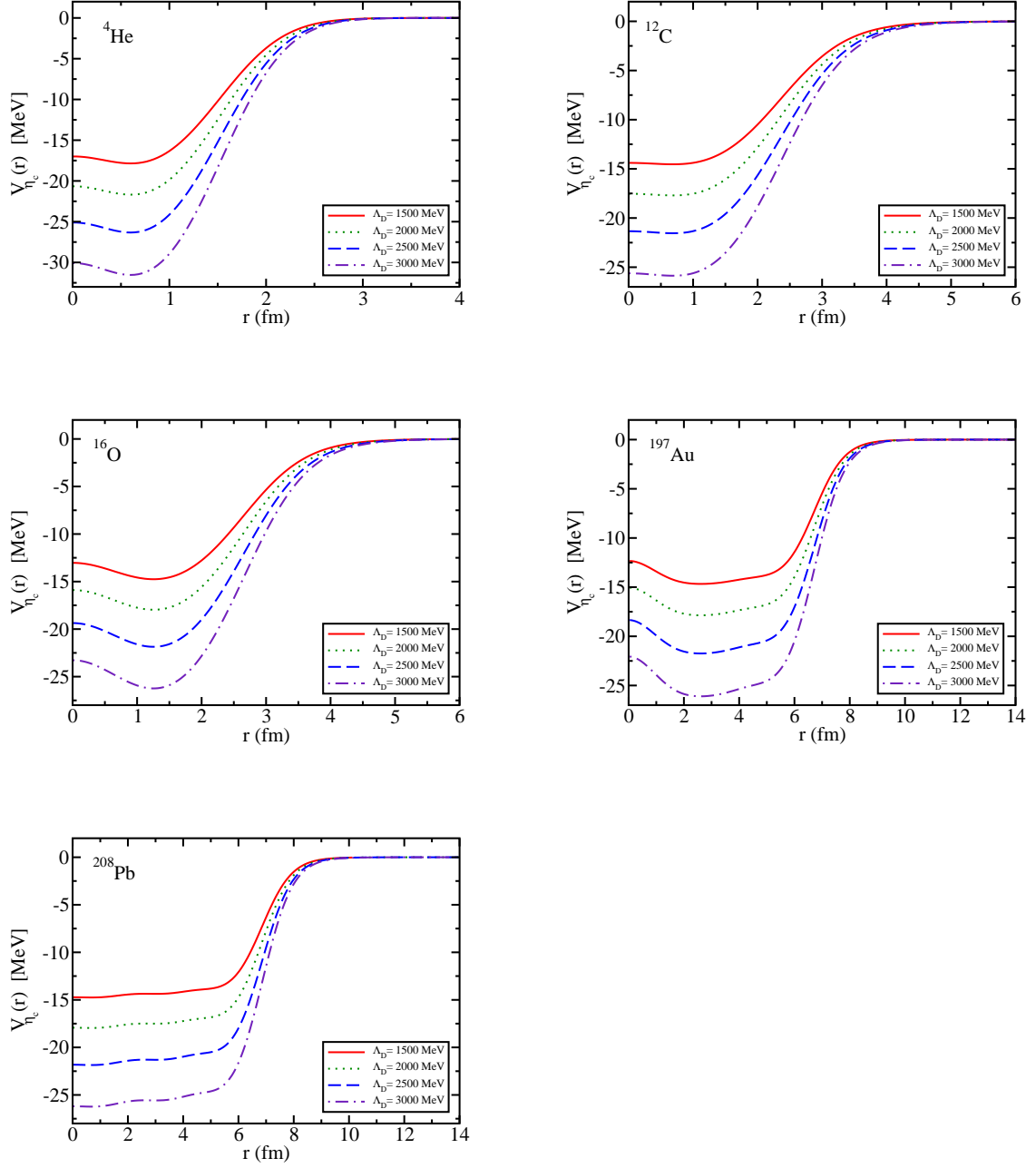


Figure 3: η_c -nucleus potentials for various nuclei and values of the cutoff parameter Λ_D .

one can see that all η_c potentials are attractive but their depths depend on Λ_D , being deeper for larger Λ_D .

		Bound state energies			
$n\ell$		$\Lambda_D = 1500$	$\Lambda_D = 2000$	$\Lambda_D = 2500$	$\Lambda_D = 3000$
${}^4_{\eta_c}\text{He}$	1s	-1.49	-3.11	-5.49	-8.55
${}^{12}_{\eta_c}\text{C}$	1s	-5.91	-8.27	-11.28	-14.79
	1p	-0.28	-1.63	-3.69	-6.33
${}^{16}_{\eta_c}\text{O}$	1s	-7.35	-9.92	-13.15	-16.87
	1p	-1.94	-3.87	-6.48	-9.63
${}^{197}_{\eta_c}\text{Au}$	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
${}^{208}_{\eta_c}\text{Pb}$	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

Table 1: η_c -nucleus bound state energies for various nuclei. All dimensionful quantities are in MeV.

Now we calculate the η_c -nucleus bound state energies for the nuclei listed above by solving the Klein-Gordon equation

$$\left(-\nabla^2 + m^2 + 2mV(\vec{r})\right) \phi_{\eta_c}(\vec{r}) = \mathcal{E}^2 \phi_{\eta_c}(\vec{r}), \quad (8)$$

where m is the reduced mass of the η_c -nucleus system in vacuum and $V(\vec{r})$ is the η_c -nucleus potential given in Eq. (7).

The computed bound state energies (E) of the η_c -nucleus system, given by $E = \mathcal{E} - m$, where \mathcal{E} is the energy eigenvalue in Eq. (8), are listed in Table 1 for four values of Λ_D . The results in Table 1 show that the η_c -meson is expected to form bound states with all the nuclei considered, independently of the value of Λ_D . Clearly, however, particular values for the bound state energies are dependent on Λ_D . This dependence was expected since the η_c potentials are also dependent on Λ_D , and it is therefore an uncertainty in the results obtained in our approach. Finally, note that the binding is stronger for larger values of Λ_D and that the η_c is predicted to be bound more strongly to heavier nuclei.

4. Summary and Conclusions

We have calculated the η_c -nucleus bound states energies for various nuclei. The η_c -nucleus potentials were calculated in the a local density approximation from the η_c mass shift as a function of the nuclear matter density and the baryon density profiles of the nuclei studied. Using these potentials, we have solved the Klein-Gordon equation to obtain η_c -nucleus bound state energies. 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

Λ_D 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. Note that we have ignored the natural width of 32 MeV in free space of the η_c and this could be an issue related to the observability of the predicted bound states. However, it should still be possible to see that there are bound states, which is the main result of this work. Addition of an imaginary part to the η_c potentials is underway and will be reported elsewhere. Furthermore, we have done an initial study [6] using a different form factor since this may impact the results reported here.

5. Acknowledgments

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References

- [1] S. J. Brodsky, I. Schmidt and G. de Teramond, *Phys. Rev. Lett.* **64** (1990), 1011
- [2] G. Krein, A. W. Thomas and K. Tsushima, *Prog. Part. Nucl. Phys.* **100**, 161 (2018)
- [3] G. Krein, A. W. Thomas and K. Tsushima, *Phys. Lett. B* **697** (2011), 136-141
- [4] K. Tsushima, D. Lu, G. Krein and A. W. Thomas, *Phys. Rev. C* **83** (2011), 065208
- [5] J. J. Cobos-Martínez, K. Tsushima, G. Krein and A. W. Thomas, *Phys. Lett. B* **811** (2020), 135882
- [6] G. N. Zeminiani, J. J. Cobos-Martínez and K. Tsushima, *Eur. Phys. J. A* **57** (2021) no.8, 259
- [7] J. J. Cobos-Martínez, K. Tsushima, G. Krein and A. W. Thomas, *Phys. Lett. B* **771** (2017), 113-118
- [8] J. J. Cobos-Martínez, K. Tsushima, G. Krein and A. W. Thomas, *Phys. Rev. C* **96** (2017) no.3, 035201
- [9] Z. w. Lin and C. M. Ko, *Phys. Rev. C* **62**, 034903 (2000)
- [10] W. Lucha, D. Melikhov, H. Sazdjian and S. Simula, *Phys. Rev. D* **93**, no. 1, 016004 (2016); Addendum: [*Phys. Rev. D* **93**, no. 1, 019902 (2016)]
- [11] K. Saito, K. Tsushima and A. W. Thomas, *Phys. Rev. C* **56**, 566 (1997)