

Study of J/ψ mass shift and bound states: Impact of DD and DD^* meson loops

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We investigate the modification of the J/ψ meson mass in asymmetric nuclear matter at zero and finite temperatures employing an effective Lagrangian approach that considers the contributions of DD and DD^* meson loops. The medium dependence of D meson masses is determined using the hadronic chiral SU(3) model, where scalar condensates are calculated and subsequently utilized in the QCD sum rules approach. Our findings indicate that an increase in baryonic density results in a negative mass shift of the J/ψ meson. This suggests that the J/ψ meson is attracted to nuclear mean fields indicating the possibility of the formation of meson-nucleus bound states. Moreover, we have also determined the binding energy and absorption decay width of the J/ψ meson for both the ground and excited states of O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} nuclei. These results are expected to contribute to the understanding of experimental data from upcoming studies at Jefferson lab and the facility for antiproton and ion research, where low-momentum charmed mesons can be produced and examined within nuclei.

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I. INTRODUCTION

A quantitative description of strongly interacting matter is still an open challenge, especially for hadrons immersed in hot and dense matter, where they serve as probes of chiral symmetry. The study of heavy-light and heavy-heavy meson interactions with nucleons is essential for gaining insights into the confinement nature of quantum chromodynamics (QCD), both in a vacuum and under conditions of extreme temperature and density [1–6]. Unlike the deconfined quark-gluon plasma (QGP) where J/ψ suppression is associated with color debye screening as per the Matsui-Satz model [7–10], the present study entirely focused on the dense hadronic phase where the deconfinement phase transition has not yet taken place (quarks remain bound) and D , D^* , and J/ψ can be treated as hadronic quasiparticles. Thus, the QGP physics serves only as a broader motivation for charmonium studies as a contrast to the confined medium mass-shift problem addressed here, rather than as a dynamical framework for the present calculation [11–13]. Modern phenomenology understands that J/ψ suppression involves multiple stages influenced

by both partonic and hadronic dynamics. As detailed in quarkonium transport studies [13], the final charmonium yields depend not only on color screening in the deconfined QGP or on a static screened potential but also on statistical recombination [14–16], cold nuclear matter effects [17], late-stage hadronic dissociation through comover scattering in the dense, confined phase [18,19], and thermal inelastic dissociation [20,21]. Open heavy-flavor mesons, such as D mesons, exhibit a greater sensitivity to the surrounding medium than heavy quarkonia like J/ψ meson. This increased sensitivity is due to their interaction with light-quark condensates, which undergo significant changes with variations in temperature and density. As a result, these mesons are valuable tools for probing the dense hadronic medium prior to deconfinement and the partial restoration of chiral symmetry, which is crucial for understanding the conditions of the early universe and the nature of compact stars. At the facility for antiproton and ion research (FAIR), the compressed baryonic matter (CBM) collaboration seeks to explore the characteristics of hadrons within dense nuclear matter [22–24]. Concurrently, the PANDA (antiproton Annihilation at Darmstadt) collaboration is set to offer valuable insights into D meson and charmonium spectroscopy, alongside advancing the understanding of strangeness physics and electromagnetic processes [24–26]. Moreover, the PANDA experiment is designed to probe bound states extending to the charm sector, with the potential to explore heavier quark states, subject to the constraints of experimental capabilities [24,27]. The experiments, such as the continuous electron beam

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accelerator facility (CEBAF) at Jefferson Lab delve into low-momentum meson production within nuclei, offering insight into the existence and behavior of charm quarks [28]. The J-PARC-E29 collaboration intends to investigate the modification of meson masses by examining meson-nucleus bound states, which are produced through the process of antiproton annihilation [29]. The experimental determination of in-medium mass modifications remains a difficult challenge. Consequently, more comprehensive and detailed theoretical studies are essential to enhance our understanding of these effects.

Significant progress has been made in the theoretical study of heavy-flavor mesons and heavy-quarkonium interactions in hadronic media over the past several decades using numerous nonperturbative theoretical frameworks [12,30–61]. Early research relied on phenomenological potential models [12,30–32] to describe the binding and interaction dynamics between heavy mesons and nucleons. As the field progressed, more systematic approaches, such as QCD sum rules, were adopted to analyze the medium modifications of heavy-flavor meson properties by relating them to changes observed in quark and gluon condensates [33–42]. Coupled-channel approaches [43–46] and heavy meson effective theories [47–49] provide a more dynamical description by incorporating multichannel effects and heavy quark symmetry, respectively. Models based on quark-meson coupling (QMC) [50–56] have also been utilized to examine the interactions of heavy mesons with nucleons in dense matter. Moreover, other theoretical approaches for quarkonium–nucleon interactions include charmed meson loop contributions [53–57], color polarizability effects [61–63], and long-range van der Waals-type interactions [57–61]. Lattice QCD simulations have recently offered significant insights into charmonium–nucleon interactions and the potential for forming bound systems with nuclei, although many investigations still rely on heavier pion masses [64,65].

Previous theoretical studies have reported several estimates for the behavior of the J/ψ meson in nuclear medium. The QCD sum rules studies indicate a relatively small mass reduction of 4 to 7 MeV in the J/ψ meson [34,35,66]. On the other hand, analyses based on color polarizability predict a larger decrease in J/ψ meson mass, about 21 MeV [61]. Additionally, calculations carried out within the QMC model show that medium modifications of the DD , DD^* , and D^*D^* meson loops at normal nuclear matter density lead to a J/ψ mass shift ranging from -16 to -24 MeV [53]. The unexpected impact of the heavier D^*D^* meson loop on the J/ψ self-energy was observed, the authors have revised their prediction of the J/ψ mass shift by incorporating the contribution from the DD meson loop only, as detailed in Refs. [54,67]. The unexpected finding may be linked to two factors: the assumption of identical coupling strengths for the DD , DD^* , and D^*D^*

meson loops and uncertainties in the formulation of the effective Lagrangian [54,56]. The chiral effective model [68–71] has been widely used to explore the behavior of heavy quark systems in a dense hadronic environment, including the studies of the in-medium masses of charmonium and bottomonium states [72,73] and open heavy-flavor mesons [72–75]. In order to examine the decay mechanisms of quarkonium into open heavy-flavor mesons within hadronic matter, a combined approach of the chiral SU(3) model and the $^3P_0$ pair-creation model was utilized [76,77]. Investigating J/ψ photoproduction in the threshold kinematic region is crucial for providing input to theoretical models that calculate the gluon structure of the proton. This includes the generalized parton distribution (GPD), the mass radii of proton, and the contribution of the trace anomaly to the proton mass [78–80]. In Ref. [81], the in-medium mass shift of the J/ψ meson was analyzed in symmetric nuclear and strange matter, taking into account only the pseudoscalar DD meson loop contribution.

The study of attractive meson-nucleus interaction would provide evidence for the negative mass shift of mesons and the role of gluon forces in nuclei [1–4,6,82,83]. Facilities such as JLab and FAIR will be able to produce low-momentum J/ψ -nucleus bound states [24,27,28]. The second-order Stark effect [84], which results from chromo-electric polarizability, plays a crucial role in determining the mass shift of charmonium in nuclear matter. Assuming a charm quark of infinite mass, predictions indicated a binding energy of approximately 10 MeV for the J/ψ in nuclear matter and 70 MeV for the excited charmonium states [59,62]. According to Ref. [57], the binding energy for J/ψ in nuclear matter was found to be 8 MeV. In contrast, when the same theoretical method was applied with a finite charm quark mass and realistic charmonium bound state wave functions, the binding energy for the excited charmonium states exceeded 100 MeV. Coupled-channel approaches have been successful in identifying D meson bound states, which hold significant importance for the PANDA and CBM experiments at FAIR [85,86]. Similar investigations have been conducted for various mesons such as ρ , ω , η , ϕ , and J/ψ mesons [6,67,87–91]. In Refs. [88,92], the energies of heavy meson-nucleus bound states were determined by solving the Klein-Gordon equations, using real potentials that were calculated self-consistently within the quark meson field model. Most recently, Ref. [67] presented updated results on J/ψ -nucleus and η_c -nucleus bound states, incorporating an intermediate DD^* state in the calculation of the η_c self-energy.

In this study, we explore the modification of the J/ψ mass in isospin asymmetric nuclear matter at both zero and finite temperatures, employing an effective Lagrangian approach that includes DD and DD^* meson loops. Additionally, we examine the binding energy and

absorption decay width of the J/ψ meson for the ground and excited states of the O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} nuclei by solving the Klein-Gordon equation. The medium dependence of D meson masses is determined using the hadronic chiral SU(3) model, where scalar condensates are calculated and subsequently utilized in the QCD sum rules. Within the chiral SU(3) model, the dilaton field, χ , is introduced primarily to replicate the gluon dynamics of QCD and to account for the effects of QCD scale symmetry breaking at the hadronic level [68,93]. The finite isospin asymmetry of the medium is explained by the vector-isovector field ρ and the scalar isovector field δ .

This manuscript is structured as follows. In Sec. II, we introduce the chiral SU(3) hadronic model and present the relevant results. Section III describes the QCD sum rules approach and the obtained results for the D meson masses. Section IV focuses on the effective Lagrangian framework and the corresponding results of medium modifications of the J/ψ meson. We present an overview of mesic-nucleus potentials, along with a discussion on the binding energy and absorption decay width of mesons in nuclei, in Sec. V. The main findings are summarized in Sec. VI.

II. THE CHIRAL SU(3) HADRONIC MODEL

The chiral SU(3) hadronic mean-field model provides an effective theoretical framework for analyzing the nonperturbative dynamics of strong interactions, founded on a nonlinear realization of chiral symmetry and the effects of broken scale invariance [68–70]. The effective Lagrangian density associated with this model is represented as [68,93,94]

$$\mathcal{L}_{\text{chiral}} = \mathcal{L}_{\text{kin}} + \sum_{M=P,S,V,A} \mathcal{L}_{\text{BM}} + \mathcal{L}_{\text{Vec}} + \mathcal{L}_0 + \mathcal{L}_{\text{ESB}}, \quad (1)$$

where $\mathcal{L}_{\text{kin}}$ denotes the terms related to kinetic energy for different particles and is written as

$$\begin{aligned} \mathcal{L}_{\text{kin}} = & i\text{Tr}\bar{B}\gamma_\mu D^\mu B + \frac{1}{2}\text{Tr}D_\mu X D^\mu X \\ & + \text{Tr}(u_\mu X u^\mu X + X u_\mu u^\mu X) + \frac{1}{2}\text{Tr}D_\mu Y D^\mu Y \\ & + \frac{1}{2}D_\mu \chi D^\mu \chi - \frac{1}{4}\text{Tr}(V_{\mu\nu} V^{\mu\nu}) - \frac{1}{4}\text{Tr}(\mathcal{A}_{\mu\nu} \mathcal{A}^{\mu\nu}). \end{aligned} \quad (2)$$

The initial term in the equation above represents the kinetic energy for the baryon octet. Within this expression, the covariant derivative D_μ is expressed as $D_\mu B = \partial_\mu B + i[\Gamma_\mu, B]$, where $\Gamma_\mu = -\frac{i}{4}[u^\dagger \partial_\mu u - \partial_\mu u^\dagger u + u \partial_\mu u^\dagger - \partial_\mu u u^\dagger]$. Moreover, $u = \exp[\frac{i}{f_0} \pi^a \lambda^a \gamma_5] = \exp[iP/\sqrt{2}f_0]$, serves as the unitary transformation operator, which is defined using the pseudoscalar meson matrix P [68]. The second and third terms of Eq. (2) are the kinetic

energy for the scalar and pseudoscalar mesons, respectively, where X represents the scalar meson matrix and $u_\mu = -\frac{i}{2}[u^\dagger(\partial_\mu u) - u(\partial_\mu u^\dagger)]$. The fourth and fifth terms represent the kinetic energy terms for the pseudoscalar singlet Y and the dilaton field χ , respectively. The term having field tensors, $V_{\mu\nu}$, represents the kinetic term of spin-1 vector mesons whereas the last term with field tensor, $\mathcal{A}_{\mu\nu}$, is for the axial vector mesons. The second term, $\mathcal{L}_{\text{BM}}$, of Eq. (1) characterizes interaction between baryons and mesons, which can be described through the exchange of scalar fields σ , ζ , δ , as well as vector fields ω , and ρ , written as

$$\mathcal{L}_{\text{BM}} = -\sum_i \bar{\psi}_i [m_i^* + g_{\omega i} \gamma_0 \omega + g_{\rho i} \gamma_0 I_{3i} \rho] \psi_i, \quad (3)$$

here, the summation is taken over all nucleons, with ψ_i representing their respective fields. The baryon mass in the medium is indicated by $m_i^* = -(g_{\sigma i} \sigma + g_{\zeta i} \zeta + g_{\delta i} I_{3i} \delta) + m_{i0}$ with $g_{\sigma i}$, $g_{\zeta i}$, and $g_{\delta i}$ are couplings of scalar fields σ , ζ , and δ , respectively, with nucleons. The vacuum masses of baryons are fitted by the additional mass term m_{i0} . In Eq. (1), the third term, $\mathcal{L}_0$, represents the nonlinear interaction of scalar mesons with themselves and is given as

$$\begin{aligned} \mathcal{L}_0 = & -\frac{1}{2}k_0 \chi^2 (\sigma^2 + \delta^2 + \zeta^2) + k_1 (\sigma^2 + \delta^2 + \zeta^2)^2 \\ & + k_2 \left(\frac{\sigma^4}{2} + \frac{\delta^4}{2} + 3\sigma^2 \delta^2 + \zeta^4 \right) \\ & + k_3 \chi \zeta (\sigma^2 - \delta^2) - k_4 \chi^4 - \frac{1}{4} \chi^4 \ln \frac{\chi^4}{\chi_0^4} \\ & + \frac{d}{3} \chi^4 \ln \left(\left(\frac{(\sigma^2 - \delta^2) \zeta}{\sigma_0^2 \zeta_0} \right) \left(\frac{\chi}{\chi_0} \right)^3 \right), \end{aligned} \quad (4)$$

where the vacuum values of the σ , ζ , and χ fields are denoted by σ_0 , ζ_0 , and χ_0 , respectively. The parameters k_0 , k_2 , and k_4 serve to ensure the vacuum extrema for the respective σ , ζ , and χ field equations. For remaining coefficients, k_3 are adjusted to match the vacuum masses of the η and η' mesons, whereas k_1 is adjusted within a mean-field framework featuring quartic vector meson couplings to give an effective nucleon mass near $0.65m_N$ at saturation density and a σ -meson mass on the order of 500 MeV [68]. The fourth term in Eq. (1), describes the vector mesons' mass term and their self-interactions are incorporated through the following relation

$$\mathcal{L}_{\text{vec}} = \frac{1}{2}(m_\omega^2 \omega^2 + m_\rho^2 \rho^2) \frac{\chi^2}{\chi_0^2} + g_4(\omega^4 + 6\omega^2 \rho^2 + \rho^4). \quad (5)$$

The chiral SU(3) model's explicit chiral symmetry breaking, which removes the Goldstone bosons, is described by

TABLE I. The coupling constants and other parameters used in the current work.

Parameter	Value	Parameter	Value	Parameter	Value
k_0	2.53	σ_0 (MeV)	-93.3	$g_{\sigma N}$	10.6
k_1	1.35	ζ_0 (MeV)	-106.76	$g_{\zeta N}$	-0.46
k_2	-4.77	χ_0 (MeV)	409.76	$g_{\delta N}$	2.34
k_3	-2.77	d	0.064	$g_{\omega N}$	13.42
k_4	-0.21	g_4	79.9	$g_{\rho N}$	5.48
m_π (MeV)	139	m_K (MeV)	498	f_π (MeV)	93.3
f_K (MeV)	122.14	ρ_0 (fm ⁻³)	0.15	m_N (MeV)	939
m_ω (MeV)	783	m_ρ (MeV)	770	m_ϕ (MeV)	1020

the final term, $\mathcal{L}_{SB}$, in Eq. (1). Its expression in terms of mesonic fields is provided by

$$\mathcal{L}_{SB} = -\left(\frac{\chi}{\chi_0}\right)^2 \left[m_\pi^2 f_\pi \sigma + \left(\sqrt{2} m_K^2 f_K - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right) \zeta \right]. \quad (6)$$

The decay constants of pions and kaons in the equation above are denoted by f_π and f_K , respectively, while their masses in free space are denoted by m_π and m_K . The key parameters, including vacuum masses of mesons and coupling constants of mesonic fields with baryons within the nuclear medium used in this work, are listed in Table I, and the fitting of parameters is detailed in Refs. [68,93]. The mean-field approximation is utilized to examine the characteristics of hadronic matter by replacing the quantum field operators with their classical expectation values for scalar and vector fields, as these fields are anticipated to exhibit minimal fluctuations in dense environments [68,93]. Applying the Euler-Lagrange equation of motion, $\frac{\partial \mathcal{L}}{\partial \phi_i} - \partial_\mu \left(\frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi_i)} \right) = 0$, for given field ϕ_i , we obtain a set of coupled equations for the scalar fields σ , ζ , δ , the dilaton field χ , and the vector fields ω and ρ [40,72,93]. These are

$$k_0 \chi^2 \sigma - 4k_1 (\sigma^2 + \zeta^2 + \delta^2) \sigma - 2k_2 (\sigma^3 + 3\sigma \delta^2) - 2k_3 \chi \sigma \zeta - \frac{d}{3} \chi^4 \left(\frac{2\sigma}{\sigma^2 - \delta^2} \right) + \frac{\chi^2}{\chi_0^2} m_\pi^2 f_\pi = \sum_i g_{\sigma i} \rho_i^s, \quad (7)$$

$$k_0 \chi^2 \zeta - 4k_1 (\sigma^2 + \zeta^2 + \delta^2) \zeta - 4k_2 \zeta^3 - k_3 \chi (\sigma^2 - \delta^2) - \frac{d\chi^4}{3\zeta} + \left(\frac{\chi}{\chi_0} \right)^2 \left[\sqrt{2} m_K^2 f_K - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right] = \sum_i g_{\zeta i} \rho_i^s, \quad (8)$$

$$k_0 \chi^2 \delta - 4k_1 (\sigma^2 + \zeta^2 + \delta^2) \delta - 2k_2 (\delta^3 + 3\delta \sigma^2) + 2k_3 \chi \delta \zeta + \frac{d}{3} \chi^4 \left(\frac{2\delta}{\sigma^2 - \delta^2} \right) = \sum_i g_{\delta i} I_{3i} \rho_i^s, \quad (9)$$

$$k_0 \chi (\sigma^2 + \zeta^2 + \delta^2) - k_3 \chi (\sigma^2 - \delta^2) \zeta + \left[4k_4 - d + 1 + \ln \frac{\chi^4}{\chi_0^4} - \frac{4d}{3} \ln \left(\frac{(\sigma^2 - \delta^2) \zeta}{\sigma_0^2 \zeta_0} \left(\frac{\chi^3}{\chi_0^3} \right) \right) \right] \chi^3 + \frac{2\chi}{\chi_0^2} \left[m_\pi^2 f_\pi \sigma + \left(\sqrt{2} m_K^2 f_K - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right) \zeta \right] - \frac{\chi}{\chi_0^2} (m_\omega^2 \omega^2 + m_\rho^2 \rho^2) = 0, \quad (10)$$

$$\frac{\chi^2}{\chi_0^2} (m_\omega^2 \omega) + 4g_4 \omega^3 + 12g_4 \omega \rho^2 = \sum_i g_{\omega i} \rho_i^v, \quad (11)$$

$$\frac{\chi^2}{\chi_0^2} (m_\rho^2 \rho) + 4g_4 \rho^3 + 12g_4 \omega^2 \rho = \sum_i g_{\rho i} I_{3i} \rho_i^v. \quad (12)$$

For different values of baryon density (ρ_B), temperature (T), and isospin asymmetry (I_a), the system of nonlinear equations from Eqs. (7)–(12) are solved. In the above equations, ρ_i^v and ρ_i^s represent the vector and scalar densities of the i th baryon, respectively, and are given as

$$\rho_i^v = \gamma_i \int \frac{d^3 k}{(2\pi)^3} \left(\frac{1}{1 + \exp [\beta(E_i^*(k) - \mu_i^*)]} - \frac{1}{1 + \exp [\beta(E_i^*(k) + \mu_i^*)]} \right), \quad (13)$$

and

$$\rho_i^s = \gamma_i \int \frac{d^3 k}{(2\pi)^3} \frac{m_i^*}{E_i^*(k)} \left(\frac{1}{1 + \exp [\beta(E_i^*(k) - \mu_i^*)]} + \frac{1}{1 + \exp [\beta(E_i^*(k) + \mu_i^*)]} \right). \quad (14)$$

Here, the spin degeneracy factor for nucleons, denoted as γ_i , is set at 2. The in-medium energy for nucleons is calculated using the formula $E_i^*(k) = \sqrt{k^2 + m_i^{*2}}$ and the effective chemical potential, $\mu_i^* = \mu_i - g_{\omega i} \omega - g_{\rho i} I_{3i} \rho$,

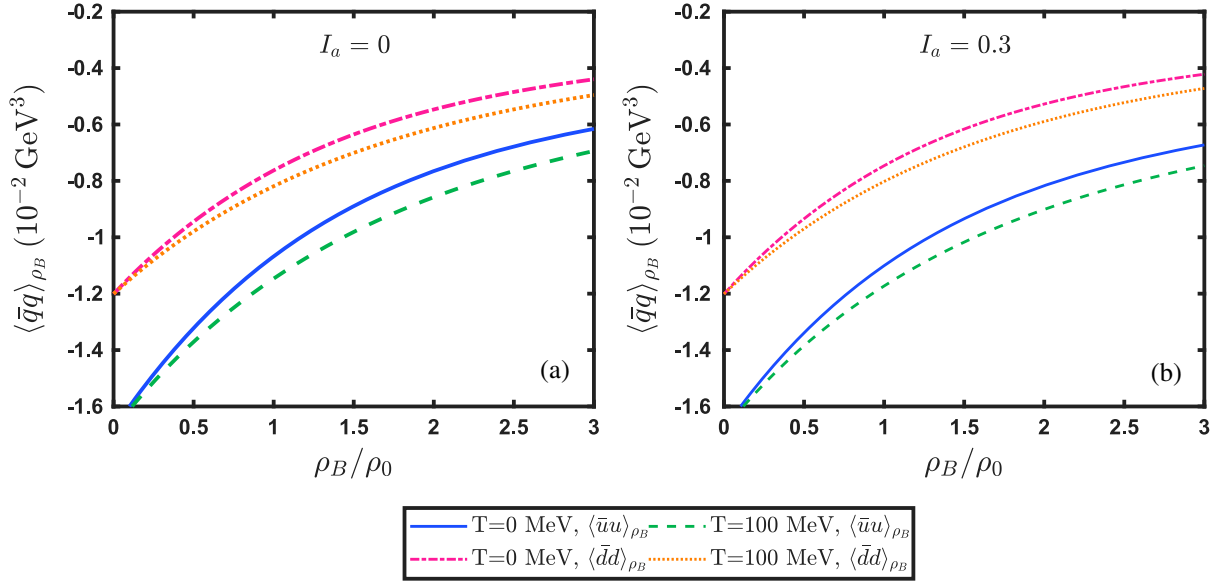


FIG. 1. The up $\langle \bar{u}u \rangle_{\rho_B}$ and down $\langle \bar{d}d \rangle_{\rho_B}$ condensates are analyzed as function of the baryon density ratio ρ_B/ρ_0 in isospin asymmetric nuclear matter, considering isospin asymmetric parameter $I_a = 0$ in (a) and 0.3 in (b), at temperatures $T = 0$ and 100 MeV.

where $g_{\omega i}$, and $g_{\rho i}$ are the coupling constants of the vector fields ω , and ρ with the nucleons, respectively. In this study, isospin asymmetry parameter $I_a = -\sum_i \frac{I_{3i} \rho_i^v}{2\rho_B}$, where ρ_B stands for the total baryonic density which can be described as $\rho_p^v + \rho_n^v$, and I_{3i} is the third component of the nucleons' isospin. To determine the in-medium masses of D and D^* mesons using QCD sum rules [94,95], it is essential to evaluate the light quark condensates $\langle \bar{u}u \rangle$ and $\langle \bar{d}d \rangle$, along with the scalar gluon condensate $\langle \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_a^{\mu\nu} \rangle$. In the chiral SU(3) model, the scalar light quark condensates are linked to explicit chiral symmetry breaking through the following relation [96],

$$\sum_i m_i \langle \bar{q}_i q_i \rangle_{\rho_B} = -\mathcal{L}_{\text{ESB}}. \quad (15)$$

The scalar light quark condensates are given by

$$\langle \bar{u}u \rangle_{\rho_B} = \frac{1}{m_u} \left(\frac{\chi}{\chi_0} \right)^2 \left[\frac{1}{2} m_\pi^2 f_\pi (\sigma + \delta) \right], \quad (16)$$

and

$$\langle \bar{d}d \rangle_{\rho_B} = \frac{1}{m_d} \left(\frac{\chi}{\chi_0} \right)^2 \left[\frac{1}{2} m_\pi^2 f_\pi (\sigma - \delta) \right], \quad (17)$$

here, m_u and m_d denote the masses of the up and down quarks, i.e., $m_u = 5$ MeV and $m_d = 7$ MeV.

In QCD, the breaking of scale invariance results in the energy-momentum tensor having a nonzero trace, which is directly linked to the scalar gluon condensate and given as [68,97]

$$\begin{aligned} & \left\langle \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_a^{\mu\nu} \right\rangle_{\rho_B} \\ &= \frac{8}{9} \left[(1-d) \chi^4 + \left(\frac{\chi}{\chi_0} \right)^2 \right. \\ & \quad \times \left. \left(m_\pi^2 f_\pi \sigma + \left(\sqrt{2} m_K^2 f_K - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right) \zeta \right) \right], \quad (18) \end{aligned}$$

where $d = \frac{2}{11}$ serves as a parameter associated with the QCD beta function [68]. In the present work, we examine the chiral model's predictions for determining the in-medium values of the chiral condensates, particularly emphasis on the light quark condensates $\langle \bar{u}u \rangle_{\rho_B}$, $\langle \bar{d}d \rangle_{\rho_B}$, and the gluon condensate $\langle \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_a^{\mu\nu} \rangle_{\rho_B}$. In Figs. 1 and 2, the results for these condensates are illustrated for isospin asymmetry values $I_a = 0$ and 0.3, and at temperatures $T = 0$ and 100 MeV. The condensate outputs are subsequently utilized as input in QCD sum rules to determine the effective masses of the pseudoscalar D (D^+ , D^0) and vector D^* (D^{*+} , D^{*0}) mesons within an isospin-asymmetric nuclear medium, as elaborated in Sec III. In Fig. 1(a), for $I_a = 0$, the results indicate that the magnitude of condensates $\langle \bar{u}u \rangle_{\rho_B}$ and $\langle \bar{d}d \rangle_{\rho_B}$ (of the order of 10^{-2} GeV^3) decrease as the baryonic density increases at a given temperature.

As the temperature rises from zero to 100 MeV, there is a less pronounced drop in the magnitude of quark condensates than the value observed at $T = 0$ MeV at a given ρ_B/ρ_0 as illustrated in Fig. 1(a). For instance, the value of $\langle \bar{u}u \rangle_{\rho_B}$ ($\langle \bar{d}d \rangle_{\rho_B}$) is found to be $-1.10(-0.74) \times 10^{-2} \text{ GeV}^3$ at a temperature $T = 0$ MeV. However, when the temperature reaches $T = 100$ MeV, these values change to

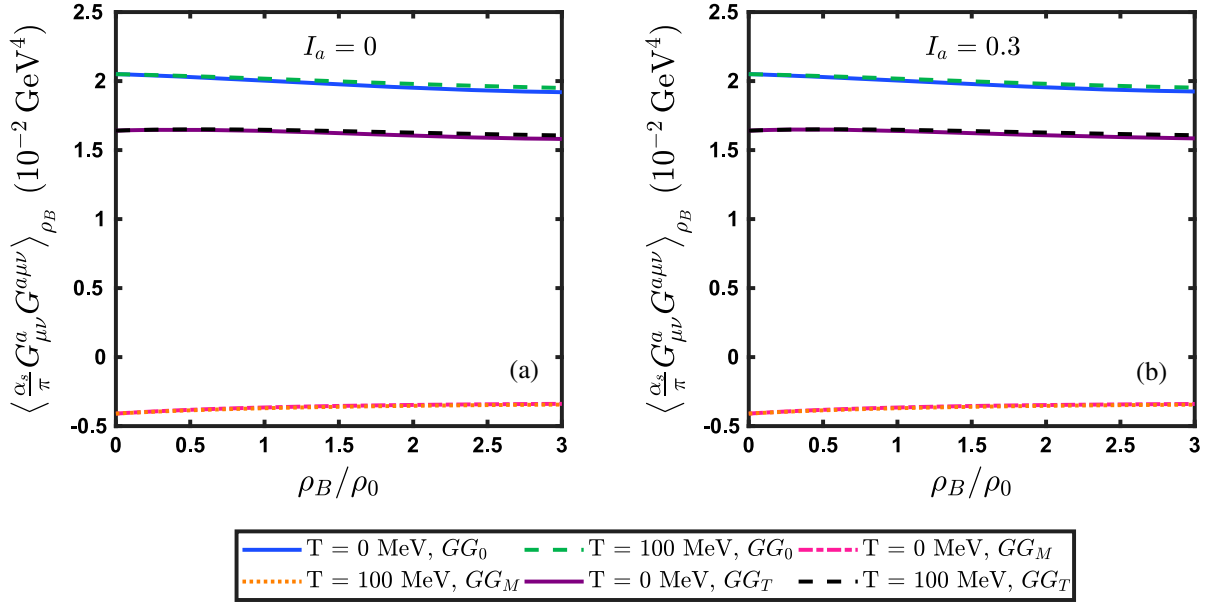


FIG. 2. The individual and total terms of gluon condensates (GG_0 , GG_M , and GG_T) are examined as function of the baryon density ratio ρ_B/ρ_0 within nuclear matter, considering $I_a = 0$ in (a) and 0.3 in (b), at temperatures $T = 0$ and 100 MeV.

$-1.17(-0.80) \times 10^{-2} \text{ GeV}^3$, at nuclear saturation density ($\rho_0 = 0.15 \text{ fm}^{-3}$) in symmetric nuclear matter. As we move from symmetric ($I_a = 0$) to an asymmetric medium with $I_a = 0.3$, the decrease in the magnitude of the quark condensates as a function of density becomes slightly less at higher baryon densities across all considered temperatures, as shown in Fig. 1(b). Next, the total contribution of gluon condensate is evaluated using Eq. (18), which consists of two terms. The first term depends solely on the dilaton field χ , while the second term, known as the mass term, depends on σ , ζ , and δ scalar fields. Figure 2 shows the variation of individual terms of $\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{a\mu\nu} \rangle_{\rho_B}$, denoted as GG_0 (first term), GG_M (mass term), and GG_T (total term) with respect to baryonic density ratio, at a fixed value of temperature. For both $I_a = 0$ and 0.3, the gluon condensate decreases in a manner similar to the quark condensates as a function of ρ_B , but with less change in their magnitude compared to quark condensates, as depicted in Figs. 2(a) and 2(b), respectively. We observed that the impact of temperature on the trend of gluon condensate contrasts with its effect on quark condensates. This variation in temperature behavior can be traced to their underlying field dependence. From Fig. 2, we noted that GG_0 term gives a dominant contribution compared to the GG_M term because the GG_0 term incorporates a χ^4 contribution of the dilaton field, as indicated in Eq. (18). Our study reveals that, at nuclear saturation density, the value of $\langle q\bar{q} \rangle_{\rho_B}$ drops to nearly 36% and $\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{a\mu\nu} \rangle_{\rho_B}$ experiences a reduction of only 6% of their vacuum values. In literature, similar results of condensates have been reported by utilizing the linear density approximation, the Dirac-Brueckner approach, and the quark meson coupling model-II [96,98–100].

In Ref. [101], the authors apply the Hellmann-Feynman theorem to derive quark condensates, leading to a 25%–50% reduction in their value, whereas, gluon condensates, estimated through the trace anomaly, show a decrease to merely 3%–6%. As reported in Ref. [99], these condensates are assessed using the QMC-II model, quark condensates diminish linearly at low densities, but this reduction becomes less pronounced at higher densities. For gluon condensates, the values fall to 4%–6% from their vacuum. The observed decreases in the quark and gluon condensates are caused by the partial restoration of chiral symmetry, which is crucial for examining the in-medium properties of mesons using different theoretical models.

III. QCD SUM RULES (QCDSR) FOR D AND D^* MESONS

The present section outlines the formulation of the QCD sum rules [95,96], which is integrated with the chiral SU(3) model to determine the in-medium properties of D (D^+ , D^0) and D^* (D^{*+} , D^{*0}) mesons. QCDSR offers a systematic framework for probing the nonperturbative QCD, relying on the Borel transformation as an essential tool to improve the convergence of the perturbative expansion. We begin with the two-point correlation function, $\Pi_{\mu\nu}(q)$, which can be separated into components corresponding to the vacuum, a static one-nucleon, and a pion bath [94,95,102]. Thus, we have

$$\Pi_{\mu\nu}(q) = \Pi_{\mu\nu}^{(0)}(q) + \frac{\rho_B}{2m_N} T_{\mu\nu}^N(q) + \Pi_{\mu\nu}^{P.B.}(q). \quad (19)$$

In Eq. (19), $T_{\mu\nu}^N$ represents the forward scattering amplitude, and the last term $\Pi_{\mu\nu}^{P.B.}(q)$, known as the pion

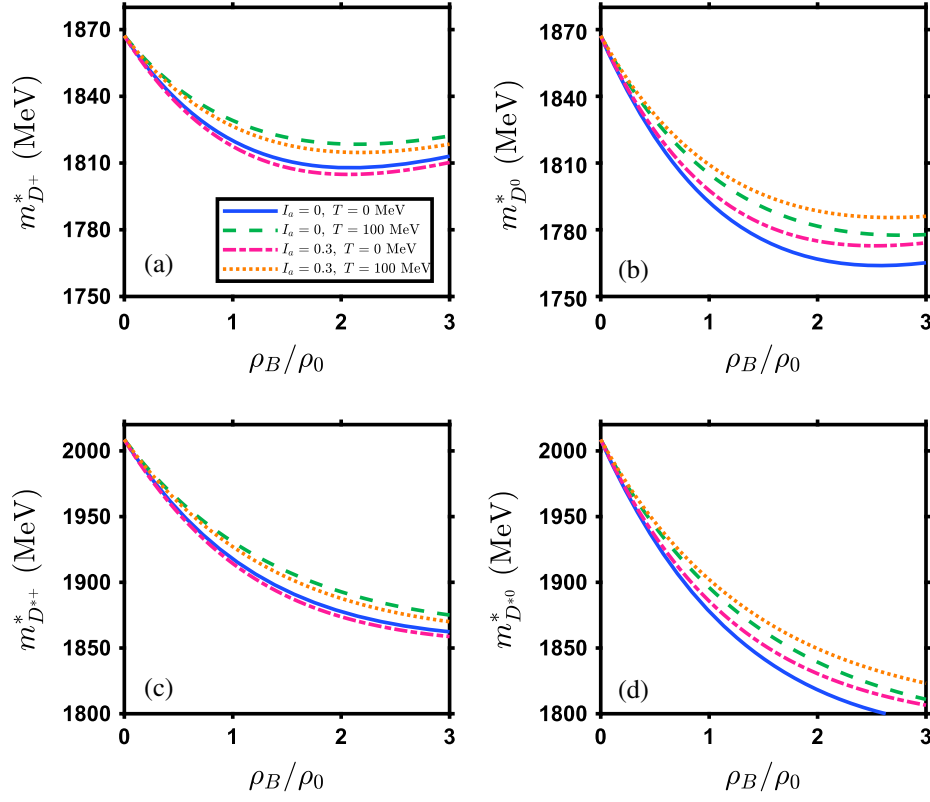


FIG. 3. In medium masses of D^+ , D^0 , D^{*+} , and D^{*0} mesons as function of the baryon density ratio ρ_B/ρ_0 in (a), (b), (c), and (d), respectively, within isospin asymmetric nuclear matter, $I_a = 0$ and 0.3 , and temperatures, $T = 0$ and 100 MeV.

bath term, has been utilized in previous studies to explore how temperature influences meson properties [95,103]. In this study, we omit the pion bath term and instead address temperature dependence on meson properties through the chiral model, which allows the grand thermodynamic potential to explicitly incorporate temperature effects via scalar and vector meson fields [72,74,104,105]. The unknown parameters a , b , and c involved in the parametrization of the phenomenological spectral densities are computed through the Borel transformation [95,103], leading to the evaluation of the mass shift of pseudoscalar D and vector D^* mesons, which can be defined as,

$$\Delta m_{D/D^*} = 2\pi \frac{m_N + m_{D/D^*}}{m_N m_{D/D^*}} \rho_B a_{D/D^*}. \quad (20)$$

The scattering length for D/D^* mesons, symbolized by a_{D/D^*} , is represented as follows [33]

$$a_D = \frac{am_c^2}{f_D^2 m_D^4 [-8\pi(m_N + m_D)]}, \quad (21)$$

and

$$a_{D^*} = \frac{a}{f_{D^*}^2 m_{D^*}^2 [-8\pi(m_N + m_{D^*})]}. \quad (22)$$

The in-medium masses of D/D^* mesons can be expressed as

$$m_{D/D^*}^* = m_{D/D^*} + \Delta m_{D/D^*}. \quad (23)$$

The vacuum masses and decay constants of D mesons are specified as $m_D = 1.867$ GeV, $m_{D^*} = 2.008$ GeV, $m_N = 0.939$ GeV, $f_D = 0.210$ GeV, $f_{D^*} = 0.270$ GeV, with the charm quark mass $m_c = 1.3$ GeV. The work reported in Refs. [98,106] involves the study of mass splitting of oppositely charged D mesons by dividing the current correlation function into even and odd terms, while this study adopts an approach based on the isospin-average meson current for D and D^* mesons [33,40,96,103]. Figure 3 illustrates how the masses of D mesons vary with baryonic density for isospin asymmetry, $I_a = 0$ and 0.3 , and temperatures $T = 0$ and 100 MeV. The masses of D mesons are evaluated using Eqs. (20)–(23), where the parameter a depends on the quark and gluon condensates, which are taken as inputs from the chiral SU(3) model as discussed in detail in Refs. [96,107]. Initially, the Borel window is chosen in a way that ensures the uncertainties in the masses of D (D^*) mesons are negligible with respect to the Borel mass parameter M^2 . In this work, Borel windows are chosen to be approximately $4.2 - 5.3$ GeV³ for D mesons and $4.7 - 5.6$ GeV³ for D^* mesons. We observed

TABLE II. In-medium values of m_D^* are presented at specific constant values of ρ_B , I_a , and T (in MeV).

m_D^* (MeV)	T = 0 MeV				T = 100 MeV			
	$I_a = 0$		$I_a = 0.3$		$I_a = 0$		$I_a = 0.3$	
	ρ_0	$3\rho_0$	ρ_0	$3\rho_0$	ρ_0	$3\rho_0$	ρ_0	$3\rho_0$
D^+	1820	1813	1817	1810	1830	1822	1826	1818
D^0	1792	1765	1797	1774	1805	1777	1809	1786
D^{*+}	1918	1862	1914	1858	1931	1875	1927	1870
D^{*0}	1878	1793	1885	1806	1896	1811	1902	1823

that masses of D^+ , D^0 , D^{*+} , and D^{*0} mesons decrease as the density of the medium increases at a given value of isospin asymmetry parameter I_a and temperature T as depicted in Figs. 3(a)–3(d), respectively. The in-medium values of m_D^* at different fixed values of ρ_B , I_a , and T are tabulated in Table II. In Fig. 3(a), with a rise in temperature from zero to 100 MeV, it is noticed that there is less drop in the magnitude of $m_{D^+}^*$ as a function of density for both symmetric and asymmetric nuclear medium. We noted that the influence of the isospin asymmetric parameter I_a on the behavior of $m_{D^+}^*$ is opposite to its influence on temperature as illustrated in Fig. 3(a). Moreover, within the same environment, the rate of drop in $m_{D^0}^*$ is more pronounced compared to $m_{D^+}^*$, as one can see in Fig. 3(b). The different masses of the u and d quarks utilized in the calculations of quark condensates [see Eqs. (16) and (17)] are responsible for the observed mass difference between $m_{D^+}^*$ and $m_{D^0}^*$ in symmetric nuclear matter, as one can see from Table II. For example, the masses of the D^+ , D^0 , D^{*+} , and D^{*0} mesons experience downward shifts of nearly 47, 75, 90, and 130 MeV, respectively, from their vacuum values (1867.2 MeV for D and 2008.6 MeV for D^*) at a temperature of $T = 0$ MeV within isospin symmetric nuclear matter. As the temperature rises to $T = 100$ MeV, these shifts are modified to ~ 37 , 62, 77, and 120 MeV at nuclear saturation density, respectively. These temperature-dependent changes in the mass shift of D mesons reflect the medium-modified behavior of the quark and gluon condensates. The inclusion of isospin asymmetry into the medium differentiates the behavior of $m_{D^+}^*$ and $m_{D^0}^*$, specifically, for a given value of ρ_B , the rate of drop in $m_{D^+}^*$ increases and while for $m_{D^0}^*$, it decreases as I_a rises as one can see in Figs. 3(a) and 3(b). The mass of vector D^* mesons decreases more significantly with an increase in baryonic density ratio at a specific temperature than that of pseudoscalar D mesons, which is attributed to the stronger attractive interaction of the D^* meson with the nuclear medium, as illustrated in Figs. 3(c) and 3(d). In literature, the QCD sum rules method has been utilized in various studies to investigate the in-medium modifications of D meson masses at nuclear saturation density ρ_0 . For instance, in Ref. [95], the authors incorporated both leading-order and next-to-leading-order corrections in symmetric nuclear

matter at zero temperature. The NLO analysis predicts a mass shift of D mesons to be -72 MeV, while considering only the LO terms results in a smaller shift of -47 MeV in D meson mass at ρ_0 . Also, the study presented in Ref. [106], the Lehmann representation for the correlation function was employed, and the QCDSR was divided into odd and even components, leading to a splitting of approximately 60 MeV between D^+ and D^- mesons. According to Ref. [50], the mass shifts of D^+ and D^- mesons were determined to be around -62 MeV at nuclear matter density within the quark-meson coupling model. The in-medium characteristics of pseudoscalar D mesons have been analyzed within a nuclear environment using an effective chiral model, which explicitly includes D meson-nucleon interactions via the dispersion relation. At ρ_0 , the mass shift for D^+ and D^0 mesons is estimated to lie between -72 and -62 MeV [72]. Moreover, Ref. [108] forecasts opposite mass shift for the D mesons, with nearly $+32$ for D^- and -27 MeV for D^+ , at $\rho_0 = 0.17 \text{ fm}^{-3}$ using a self-consistent coupled-channel approach.

IV. THE EFFECTIVE LAGRANGIAN APPROACH

In this section, we predict the results of the mass shift of the J/ψ meson in isospin asymmetric nuclear matter, utilizing an effective Lagrangian approach under the flavor SU(4) symmetry, combined with QCD sum rules. We account for DD and DD^* meson loop contributions to J/ψ self-energy. It is important to clarify the status of the approximation applied in this section. The present calculation does not represent a complete real-time finite-temperature calculation of the J/ψ spectral function in the Schwinger-Keldysh formalism [109–111]. Instead, it involves computing the self-energy of hadronic quasiparticles, where the vacuum J/ψ mass is initially set, and obtaining the in-medium mass shift of the J/ψ from the change in the open-charm meson-loop contribution. The medium dependence enters through the masses of the virtual D , and D^* mesons, which are calculated independently from QCD sum rules [95,96] using the in-medium light-quark and gluon condensates obtained in the chiral SU(3) model [68]. In the hadronic regime, the primary influence of the medium on the J/ψ loop correction is attributed to changes in the D and D^* intermediate states.

This assumption is justified because the J/ψ is a compact hidden-charm state with no light valence quark, whereas the D and $\bar{D}$ mesons contain a light quark and are therefore much more sensitive to the partial restoration of chiral symmetry and to changes in the scalar condensates. In this sense, the J/ψ does not acquire its temperature dependence from a direct thermal bath contribution in the loop integral. However, the temperature and density dependence are transmitted through the medium-modified open-charm thresholds. A full real-time finite-temperature treatment including thermal scattering cuts, explicit thermal occupation factors, and finite in-medium widths of the intermediate D and D^* mesons will be considered in the future work. For the moderate hadronic temperatures considered here, thermally populated open-charm mesons are strongly suppressed, so the dominant medium effect retained in the present approach is the shift of the virtual D and D^* thresholds rather than explicit thermal scattering from real open-charm mesons. We utilize the effective Lagrangian densities to describe the interaction of J/ψ with D and D^* mesons (the field corresponding to J/ψ will be denoted as ψ) [53,112,113], and given as

$$\mathcal{L}_{\psi DD} = ig_{\psi DD} \psi^\mu [\bar{D}(\partial_\mu D) - (\partial_\mu \bar{D})D], \quad (24)$$

$$\mathcal{L}_{\psi DD^*} = \frac{g_{\psi DD^*}}{m_\psi} \varepsilon_{\alpha\beta\mu\nu} (\partial^\alpha \psi^\beta) [(\partial^\mu \bar{D}^{*\nu})D + \bar{D}(\partial^\mu D^{*\nu})], \quad (25)$$

$$\begin{aligned} \mathcal{L}_{\psi D^* D^*} = & ig_{\psi D^* D^*} [\psi^\mu ((\partial_\mu \bar{D}^{*\nu})D_\nu^* - \bar{D}^{*\nu}(\partial_\mu D_\nu^*)) \\ & + (\partial_\mu \psi^\nu) \bar{D}_\nu^* D^{*\mu} - \psi^\nu (\partial_\mu \bar{D}_\nu^*) D^{*\mu} \\ & + \bar{D}^{*\mu} (\psi^\nu (\partial_\mu D_\nu^*) - (\partial_\mu \psi^\nu) D_\nu^*)]. \end{aligned} \quad (26)$$

The following convention is employed for the isospin doublets of the pseudoscalar and vector D -mesons:

$$\begin{aligned} D &= \begin{pmatrix} D^+ \\ D^0 \end{pmatrix}, & \bar{D} &= (\bar{D}^- \quad \bar{D}^0), \\ D^* &= \begin{pmatrix} D^{*+} \\ D^{*0} \end{pmatrix}, & \bar{D}^* &= (\bar{D}^{*-} \quad \bar{D}^{*0}). \end{aligned} \quad (27)$$

If SU(4) symmetry treated as exact, it would establish the identical relationships among the coupling constants in the Lagrangian, which are taken to be $g_{\psi DD} = g_{\psi DD^*} = g_{\psi D^* D^*} = 7.64$. These coupling constants are determined by using the vector-meson dominance model [114,115]. The mass shift of J/ψ meson is given as

$$\Delta m_\psi = m_\psi^* - m_\psi, \quad (28)$$

with in-medium masses of J/ψ meson derived from

$$m_\psi^{*2} = (m_\psi^0)^2 + \Sigma_\ell(k^2 = m_\psi^{*2}). \quad (29)$$

The term $\Sigma_\ell(k^2)$ is the total J/ψ self-energy from different loop contributions with ℓ represents DD, DD^*, D^*D^* meson loops and m_ψ^0 stands for the bare mass. By taking the average over the three J/ψ polarization, J/ψ self-energy $\Sigma_\ell(k^2)$, can be expressed as follows [53,56]:

$$\Sigma_\ell(m_\psi^{*2}) = -\frac{g_{\psi\ell}^2}{3\pi^2} \int d\mathbf{q} \mathbf{q}^2 F_\ell^2(\mathbf{q}^2) K_\ell(\mathbf{q}^2), \quad (30)$$

where $K_\ell(\mathbf{q}^2)$ for each loop contribution are given by

$$K_{DD}(\mathbf{q}^2) = \frac{\mathbf{q}^2}{\omega_D^*} \left(\frac{1}{\omega_D^{*2} - m_\psi^{*2}/4} \right), \quad (31)$$

$$K_{DD^*}(\mathbf{q}^2) = \frac{\mathbf{q}^2}{\omega_D^* \omega_{D^*}^*} \frac{\bar{\omega}_{D^*}^*}{\omega_{D^*}^{*2} - m_\psi^{*2}/4}, \quad (32)$$

$$\begin{aligned} K_{D^* D^*}(\mathbf{q}^2) &= \frac{1}{4m_\psi \omega_{D^*}^*} \left[\frac{A(q^0 = \omega_{D^*}^*)}{\omega_{D^*}^{*2} - m_\psi^{*2}/2} + \frac{A(q^0 = \omega_{D^*}^* - m_\psi^*)}{\omega_{D^*}^{*2} + m_\psi^{*2}/2} \right], \end{aligned} \quad (33)$$

here $\omega_D^* = (\mathbf{q}^2 + m_D^{*2})^{1/2}$, $\omega_{D^*}^* = (\mathbf{q}^2 + m_{D^*}^{*2})^{1/2}$, $\bar{\omega}_{D^*}^* = (\omega_D^* + \omega_{D^*}^*)/2$, and $m_{D^*}^*$ represents the isospin-averaged mass of the D/D^* meson, calculated as the average of the charged and neutral states $m_{D^+}^*/m_{D^{*0}}^*$. In above expression, $A(q) = \sum_{i=1}^4 A_i(q)$, with

$$A_1(q) = -4q^2 \left[4 - \frac{q^2 + (q-k)^2}{m_{D^*}^{*2}} + \frac{[q \cdot (q-k)]^2}{m_{D^*}^{*4}} \right],$$

$$A_2(q) = 8 \left[q^2 - \frac{[q \cdot (q-k)]^2}{m_{D^*}^{*2}} \right] \left[2 + \frac{(q^0)^2}{m_{D^*}^{*2}} \right],$$

$$\begin{aligned} A_3(q) = & 8(2q^0 - m_\psi^*) \left[q^0 - (2q^0 - m_\psi^*) \frac{q^2 + q \cdot (q-k)}{m_{D^*}^{*2}} \right. \\ & \left. + q^0 \frac{[q \cdot (q-k)]^2}{m_{D^*}^{*4}} \right], \end{aligned}$$

$$\begin{aligned} A_4(q) = & -8 \left[q^0 - (q^0 - m_\psi^*) \frac{q \cdot (q-k)}{m_{D^*}^{*2}} \right] \\ & \times \left[(q^0 - m_\psi^*) - q^0 \frac{q \cdot (q-k)}{m_{D^*}^{*2}} \right]. \end{aligned}$$

In the above expressions, the four-vectors q and k are defined as $q = (q^0, \mathbf{q})$ and $k = (m_\psi^*, \mathbf{0})$. To regulate the self-energy loop integrals, we apply phenomenological form factors $F_\ell(\mathbf{q}^2)$. The $F_\ell(\mathbf{q}^2)$ signifies the product of the vertex form factors ($u_\ell(\mathbf{q}^2)$), which is applied to both vertices, given as [88,112,116,117].

$$u_{D,D^*}(\mathbf{q}^2) = \left(\frac{\Lambda_D^2 + m_\psi^{*2}}{\Lambda_D^2 + 4\omega_{D/D^*}^{*2}(\mathbf{q}^2)} \right)^2. \quad (34)$$

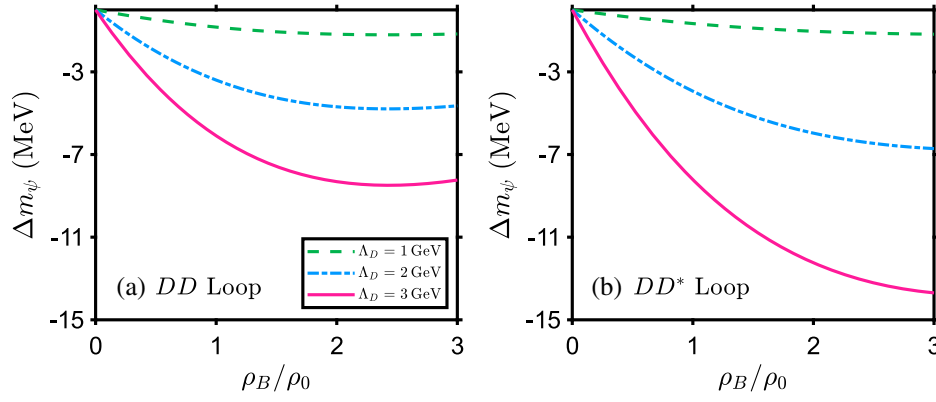


FIG. 4. Density dependence of the DD (a) and DD^* (b) loop contribution to the J/ψ mass shift (Δm_ψ) in asymmetric nuclear matter at $I = 0.3$ and $T = 100$ MeV for different values of the cutoff Λ_D .

The dipole form factors are given by $F_{DD}(\mathbf{q}^2) = u_D^2(\mathbf{q}^2)$, $F_{DD^*}(\mathbf{q}^2) = u_D(\mathbf{q}^2)u_{D^*}(\mathbf{q}^2)$, and $F_{D^*D^*}(\mathbf{q}^2) = u_{D^*}^2(\mathbf{q}^2)$, where, Λ_D is the cutoff mass parameter. The choice of Λ_D for the interaction of J/ψ meson with DD , DD^* , D^*D^* meson loops plays an important role, as they significantly influence the calculated results. These form factors are introduced to represent the finite spatial extent of the mesons and their overlapping region at the vertices. When the cutoff parameters are comparable to the meson masses, their associated Compton wavelengths are of the same order as the sizes of the mesons. This range of cutoff parameter values reduces the physical significance of the form factors [56]. Consequently, it is crucial to select cutoff values with care to accurately represent the finite-size effects of the interacting mesons. We choose $\Lambda_D = 1$ to 3 GeV to calculate the in-medium mass of J/ψ using heavy quark and meson symmetry in QCD, which is consistent with the parameters used in the study reported in Refs. [53,56].

Next, we proceed to present the results for the in-medium mass shift of the J/ψ meson, $\Delta m_{J/\psi}$, which are calculated from Eq. (28). To begin, we analyze the individual loop contributions of ψDD , ψDD^* , and ψD^*D^* separately, focusing on their impact on the mass shift of J/ψ in relation to baryonic density. This analysis is conducted for various cutoff mass values Λ_D and temperature conditions. In Fig. 4(a), the findings of the in-medium mass shift of the J/ψ meson for the DD loop are depicted as a function of the ρ_B/ρ_0 ratio in nuclear matter, with $I_a = 0.3$ and $T = 100$ MeV, using a cutoff mass parameter ranging from 1 to 3 GeV. We observe that as the density of the medium increases, the magnitude of $\Delta m_{J/\psi}$ becomes more negative for a given Λ_D . When moving to higher Λ_D values, a more substantial downward shift in $\Delta m_{J/\psi}$ is observed. From the DD^* loop contribution, a more significant negative mass shift of the J/ψ meson is noticed, as depicted in Fig. 4(b), which is consistent with the findings reported in Ref. [56]. This negative mass shift of J/ψ indicates the strong coupling of vector D^* mesons with the scalar and vector

fields of the medium compared to the pseudoscalar D meson. As a result, there is a greater drop for the DD^* meson loop contribution compared to the DD meson loop. To explicitly examine the temperature dependence, the mass shift of J/ψ meson has been evaluated at temperature $T = 0$ and 100 MeV while keeping Λ_D fixed at 2 GeV as illustrated in Fig. 5. The drop in magnitude of $\Delta m_{J/\psi}$ at $T = 100$ MeV is less pronounced than that at $T = 0$ MeV for both DD and DD^* loops, as one can see from Figs. 5(a) and 5(b). We listed the prediction of mass shift of J/ψ with the dependence of DD and DD^* meson-loop contribution at $\rho_0 = 0.15 \text{ fm}^{-3}$ within nuclear medium in Table III. Our analysis indicates a negative mass shift of the J/ψ within an isospin-asymmetric nuclear medium, varying from approximately -1.5 to -14 MeV at nuclear saturation density. This variation arises from the choice of cutoff mass parameter in the range $\Lambda_D = 1$ to 3 GeV and shows good consistency with the most recent theoretical analyses of the J/ψ meson [6], where the mass shift of the J/ψ is estimated to be between -4 and -12 MeV within same environment by using QMC model, attributed solely to the DD loop contribution [67]. In Ref. [81], the authors examined the mass shift of J/ψ within a symmetric nuclear medium and a strange medium, focusing only on the contributions from the DD meson loop. The study determined that the mass shift of J/ψ is -3.5 MeV in cold nuclear matter and -4.5 MeV in strange matter, respectively. In Fig. 6, we present the results of the D^*D^* meson loop contribution to J/ψ mass shift. An unexceptional large downward mass shift of J/ψ within an isospin symmetric nuclear medium is observed, ranging from -70 to -158 MeV at nuclear saturation density and depending on the value of the cutoff mass parameter. Our results are in close agreement with those reported in Ref. [56], where masses of J/ψ for D^*D^* meson loop were found to lie between -61 to -164 MeV using the quark meson coupling model. However, from a basic quantum mechanical perspective, it is anticipated that heavier loops would have a lesser impact. To ensure the accuracy of our

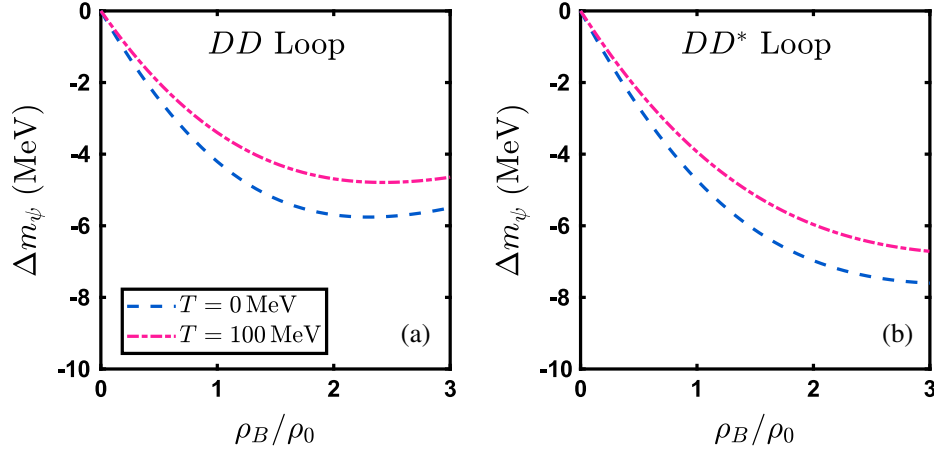


FIG. 5. Density dependence of the DD (a) and DD^* (b) loop contributions to the J/ψ mass shift (Δm_ψ) in asymmetric nuclear matter at $I = 0.3$ and $\Lambda_D = 2$ GeV for different values of temperatures.

minimum prediction and to avoid possible model-dependent uncertainties, we concentrate on the DD and DD^* meson-loop contribution [56]. These findings may offer significant insights into the mechanism responsible for J/ψ production and suppression, as observed in various experiments, such as the NA38, and NA60 collaborations at CERN SPS [118–120], the PHENIX experiments at RHIC [121,122], and the ALICE experiment at the LHC [123].

V. J/ψ -NUCLEI BOUND STATE

In this section, we predict the existence of bound states of the J/ψ meson with nuclei O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} by calculating their binding energy and absorption decay width. In a finite nucleus, the scalar quark and gluon condensates are evaluated as a function of the radial coordinate r , within the range $0 \leq r \leq R$, where R denotes the cutoff radius beyond which the baryonic density is considered zero [124]. To determine their radial dependence, we solve the coupled equations of motion for mesonic fields, which are associated with baryonic densities within the chiral $SU(3)$ model. The total baryonic density $\rho_B(r)$ of nuclei is $\rho_B(r) = \sum_{i=p,n} \rho_i(r)$ [104,125]. For nuclei with atomic number $A < 18$, the harmonic potential type density distribution is used and is represented as follows [124,125]

$$\rho_i(r) = \rho_{i,0} \left(1 + a_i \left(\frac{r}{R_i} \right)^2 \right) \exp \left[- \left(\frac{r}{R_i} \right)^2 \right], \quad (35)$$

while the two-parameter Fermi distribution density is utilized for heavy nuclei, given as

$$\rho_i(r) = \frac{\rho_{i,0}}{1 + \exp \left[\frac{r - R_i}{a_i} \right]}. \quad (36)$$

The center density of nucleons in the above equations is represented by $\rho_{i,0}$, while the associated radius and diffuseness parameters are denoted by R_i and a_i , respectively (for nucleon of type i), and expressed as

$$R_i = R_{i,e} + \frac{5r_i^2 R_{i,e}}{15R_{i,e}^2 + 7\pi^2 a_{i,e}^2}, \quad (37)$$

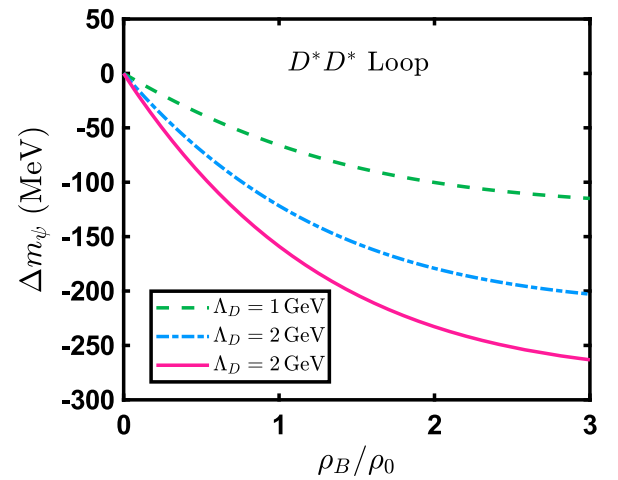


FIG. 6. The D^*D^* loop contribution to the J/ψ mass shift as function of the baryon density ratio ρ_B/ρ_0 is examined in symmetric nuclear matter at $T = 0$ MeV, with various cutoff values Λ_D .

TABLE III. Dependence of the J/ψ mass shift on cutoff parameter, Λ_D , through the contributions of DD and DD^* meson loops at $\rho_0 = 0.15 \text{ fm}^{-3}$ within an isospin asymmetric nuclear medium.

Λ_D (GeV)	$\Delta m_{J/\psi}^{DD}$ (MeV)	$\Delta m_{J/\psi}^{DD^*}$ (MeV)	$\Delta m_{J/\psi}^{DD+DD^*}$ (MeV)
1	-0.83	-0.66	-1.49
2	-3.40	-3.93	-7.32
3	-6.08	-8.20	-14.28

TABLE IV. Nucleons density parameters $R_{p,e}$, $R_{n,e}$, $a_{p,e}$ and $a_{n,e}$ for different nuclei O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} (in fm).

Nucleus	Proton		Neutron	
	$R_{p,e}$ (fm)	$a_{p,e}$ (fm)	$R_{n,e}$ (fm)	$a_{n,e}$ (fm)
O^{16}	1.833	1.544	1.815	1.529
Ca^{40}	3.51	0.563	3.43	0.563
Zr^{90}	4.484	0.55	4.94	0.55
Pb^{208}	6.624	0.549	6.89	0.549

and

$$a_i = \left(\frac{R_{i,e}^3 + \pi^2 a_{i,e}^2 R_{i,e} - R_{i,e}^3}{\pi^2 R_{i,e}} \right)^{1/2}, \quad (38)$$

here $r_i^2 = 0.69 \text{ fm}^2$, represents mean-squared radius of nucleons. Table IV presents the nucleon density parameter $R_{i,e}$ and $a_{i,e}$ for different nuclei [125]. Isospin asymmetry is crucial in understanding the properties of strongly interacting matter, particularly in neutron-rich heavy nuclei. The global asymmetry parameter is expressed as $I_a = (N - Z)/A$, where N represents the neutron number, A is the mass number, and Z denotes the proton number. While, the neutron-to-proton asymmetry is defined as $I_a = \frac{\rho_n(r) - \rho_p(r)}{2\rho_B(r)}$, which tends to reach unity at the nuclear surface [126,127].

At relatively low momenta, the meson experiences an effective potential within the nucleus, which serves as input to the Klein-Gordon equation to investigate the bound states of mesic-nuclei. The Klein-Gordon equation, incorporating a local complex potential $V(r)$, can be expressed as follows [128]

$$(-\nabla^2 + (\mu + V(r))^2)\psi(r) = \epsilon^2\psi(r), \quad (39)$$

here, the parameter μ denotes the reduced mass of the mesic-nuclei system and ϵ stands for the energy eigenvalue of Eq. (39). For a local complex potential, $V(r)$, we write [83,129]

$$V(r) = U(r) - \frac{i}{2}W(r). \quad (40)$$

The real part of the potential $U(r)$ reflects the mass modification of the meson, whereas the imaginary part $W(r)$ accounts for their absorption within the nucleus. The J/ψ meson self-energy $\Pi(r)$ can be used to describe a meson-nucleus potential in the local-density approximation as $\Pi(r) = 2V(r)\sqrt{\mathbf{p}^2 + m_\psi^2}$, which enters in dispersion relation, $E^2 - \mathbf{p}^2 - m_\psi^2 - \Pi(r) = 0$ [83,130]. For a slow meson $|\mathbf{p}| \leq m_\psi$, therefore, the optical potential of the bound system can be written as $V(r) = \frac{\Pi(r)}{2m_\psi}$. By

incorporating in-medium meson masses, $m_\psi^* = m_\psi - \Delta m_\psi(\rho_B)$, the dispersion relation can be modified as follows

$$E^2 - \mathbf{p}^2 - m_\psi^{*2} = E^2 - \mathbf{p}^2 - m_\psi^2 + 2m_\psi\Delta m_\psi(\rho_B). \quad (41)$$

In this case, terms quadratic in Δm_ψ can be omitted as it is assumed that the mass modification Δm_ψ is minimal in relation to the meson mass. Therefore, for the real part of the potential $U(r)$, we have

$$U(r) = \Delta m_\psi(\rho_B(r)) = \Delta m_\psi(\rho_{i,0}) \frac{\rho_B(r)}{\rho_{i,0}} \quad (42)$$

A microscopic calculation of the absorptive part would require the in-medium $J/\psi N$ scattering amplitude and the corresponding open inelastic channels [131,132]. Since this is not the main focus of the present work, we follow a standard phenomenological optical-potential formalism and parametrize the local absorption width as proportional to the local baryon density

$$W(r) = -\frac{1}{2}\Gamma(r) = -\frac{\text{Im}\Pi(r)}{2m} = -\frac{1}{2}\Gamma(\rho_B = \rho_{i,0}) \frac{\rho_B(r)}{\rho_{i,0}}. \quad (43)$$

In this study, we determine the value $\Gamma(\rho_B = \rho_{i,0})$ by employing the approach used in the QMC model. This model connects to the mass shift through a parameter κ , using following relation [128]

$$W(r) = (-\kappa\Delta m_\psi(\rho_{i,0}) + \Gamma_{\text{vac}}) \frac{\rho_B(r)}{\rho_{i,0}}. \quad (44)$$

In the above equation, $\Delta m_\psi(\rho_{i,0})$ is the mass shift of the J/ψ meson at the central baryonic density of nuclei, κ mimics the absorption effect, and Γ_{vac} is the vacuum decay width, which is not taken in the calculation. The cases $\kappa = 0, 0.5$, and 1 are used to estimate the sensitivity of the bound-state widths to absorption. Following this, we calculate the binding energy and absorption decay width of the J/ψ meson by solving Eq. (39) in momentum space. This calculation is performed using a Fourier transform, which incorporates partial wave decomposition of the potential [133,134]. This method enables us to derive complex eigenvalues, where the real part signifies the binding energy E , and the imaginary part indicates the decay width Γ , as defined by the equations $E = \text{Re } \epsilon - \mu$ and $\Gamma = -2 \text{Im } \epsilon$, respectively.

Now, we present the results of the J/ψ mass shift as a function of radial distance r from the center of the nuclei O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} , evaluated at various values of the cutoff parameter Λ_D . These results include contributions from the DD and DD^* loop interactions as depicted in Figs. 7 and 8. As seen in Fig. 7(a), at $\Lambda_D = 2 \text{ GeV}$,

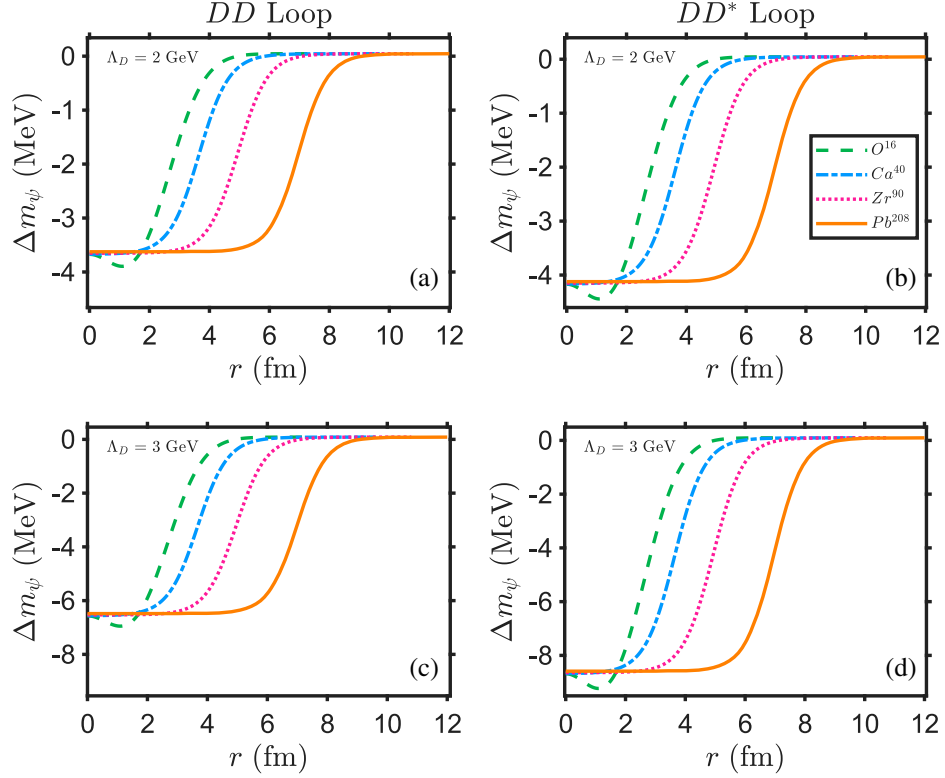


FIG. 7. The result of J/ψ mass shift as a function of r in nuclei O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} at different value of cut off parameter Λ_D , from the DD loop contribution, illustrated in subfigures (a) and (c), and the DD^* loop contribution, presented in subfigures (b) and (d).

a reduction in the magnitude of Δm_ψ is observed as one moves away from the center of the nucleus, from the contribution of DD loop interactions across all nuclei. It is noted that heavy nuclei exhibit a larger negative J/ψ mass shift compared to lighter nuclei. This effect becomes more pronounced for DD^* loop interactions, as depicted in Fig. 7(b). Furthermore, an increase in the cutoff parameter Λ_D leads to a more substantial downward shift in J/ψ mass within finite nuclei, as shown in Figs. 7(c) and 7(d). This pattern becomes more evident in Fig. 8, which illustrates the mass shift of the J/ψ mesic-nuclei,

resulting from the combined effects of DD and DD^* loop interactions. The results of J/ψ mass shift in nuclei O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} arises from the total contribution of DD and DD^* loops for various values of r and Λ_D are listed in Table V. For the J/ψ meson, we observed a negative mass shift, indicating the potential formation of a meson-nucleus bound state. The calculated binding energy and absorption decay width for both the ground and excited states of nuclei are detailed in Table VI. In the case of the light nuclei O^{16} nucleus, the bound states with negative binding energy are restricted to the $1s$ and $1p$ orbitals for

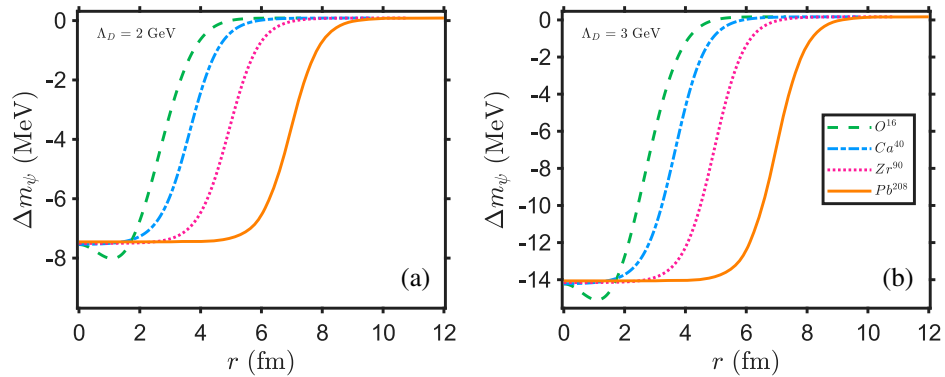


FIG. 8. The J/ψ mass shift as a function of r in nuclei O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} arises from the total contribution of DD and DD^* loops, depicted in subfigures (a) for $\Lambda_D = 2$ GeV, and (b) for $\Lambda_D = 3$ GeV.

TABLE V. The J/ψ mass shift in nuclei O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} arises from the total contribution of DD and DD^* loops for various values of r and Λ_D .

A	r (fm)	Δm_ψ (MeV) for $DD + DD^*$ loops	
		$\Lambda_D = 2$ GeV	$\Lambda_D = 3$ GeV
$^{16}_8O$	0	-7.53	-14.2
	2	-6.74	-12.7
	4	-0.63	-1.22
$^{40}_{20}Ca$	0	-7.53	-14.2
	2	-7.26	-13.7
	4	-2.42	-4.7
$^{90}_{40}Zr$	0	-7.50	14.1
	2	-7.48	-14.1
	4	-6.46	-12.2
	6	-0.70	-1.37
$^{208}_{82}Pb$	0	-7.45	-14.0
	2	-7.45	-14.0
	4	-7.43	-14
	6	-6.54	-12.4

phenomenological parameters $\kappa = 0, 0.5$, and 1 . As κ approaches unity, the magnitude of binding energy of meson is observed to decrease while its decay width values increase. This implies that the absorption of mesons in the nucleus is greater at higher κ values, which may lead to peak broadening and hinder the clear identification of bound states. The Ca^{40} nucleus exhibits stability for the $1s$, $1p$, $1d$, and $2s$ bound states, whereas

Zr^{90} , and Pb^{208} support a broader range of bound states, extending from $1s$ to $2d$ and $2p$, respectively, as illustrated in Table VI. These findings indicate that the greater the magnitude of the attractive potential in heavy nuclei, the more significantly stable states are expected. The J/ψ -nucleus system is notably stable, as its energy lies below the nucleon knockout threshold. The exothermic reaction $J/\psi N \rightarrow \eta_c N$ results in a minimal decay width, allowing this bound state to remain sufficiently narrow for experimental observation [132]. In contrast to the η and ω mesons within different nuclei [135], which experience significant broadening due to strong interactions, the J/ψ can be captured into discrete bound states through recoilless kinematics. For example, as discussed in Ref. [135], the η meson in O^{16} (Pb^{208}) nuclei exhibits a binding energy of $-32.6(-56.3)$ MeV and a decay width of $26.7(33.2)$ MeV. Similarly, the ω meson binds at $-46(-118)$ MeV with a corresponding decay width of $31.7(33.1)$ MeV. Earlier works, including the Refs. [88,67] have calculated the binding energies of J/ψ -nuclei using the QMC model, their corresponding values are shown in Table VII and are compared with our findings. Moreover, medium and heavy nuclei are expected to show a rich range of bound state spectra. However, experimental studies depend on producing J/ψ mesons with minimal relative momentum to the nucleus. By carefully choosing the beam and target nuclei, it may be possible to achieve such kinematics, as current findings suggest that numerous nuclei are likely to form J/ψ -nuclear bound states [88].

TABLE VI. The binding energy E and absorption decay width Γ of J/ψ meson for both the ground and excited states of O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} nuclei (in units of MeV).

A	nl	E (MeV)			Γ (in MeV)		
		$\kappa = 0$	$\kappa = 0.5$	$\kappa = 1$	$\kappa = 0$	$\kappa = 0.5$	$\kappa = 1$
$^{16}_8O$	$1s$	-7.827	-7.745	-7.521	0.0	6.121	12.371
	$1p$	-2.133	-1.925	-1.391	0.0	4.264	8.91
$^{40}_{20}Ca$	$1s$	-8.563	-8.513	-8.374	0.0	5.626	11.318
	$1p$	-5.054	-4.967	-4.726	0.0	4.56	9.244
	$1d$	-1.903	-1.763	-1.39	0.0	3.395	7.012
	$2s$	-2.335	-2.19	-1.801	0.0	3.322	6.873
$^{90}_{40}Zr$	$1s$	-9.982	-9.946	-9.846	0.0	6.015	12.076
	$1p$	-7.306	-7.246	-7.078	0.0	5.265	10.611
	$1d$	-4.773	-4.685	-4.442	0.0	4.488	9.100
	$2s$	-4.963	-4.872	-4.624	0.0	4.459	9.046
	$2p$	-2.807	-2.68	-2.335	0.0	3.603	7.402
	$2d$	-0.958	-0.77	-0.291	0.0	2.635	5.598
$^{208}_{82}Pb$	$1s$	-11.146	-11.121	-11.05	0.0	6.309	12.649
	$1p$	-9.231	-9.19	-9.075	0.0	5.799	11.653
	$1d$	-7.367	-7.309	-7.148	0.0	5.282	10.641
	$2s$	-7.439	-7.38	-7.216	0.0	5.271	10.622
	$2p$	-4.08	-3.98	-3.705	0.0	4.172	8.49

TABLE VII. J/ψ -nucleus bound state energies for different references (in the units of MeV).

A	nl	Bound state energies		
		Ref. [88] ($\Lambda_D = 2$ GeV)	Ref. [67] ($\Lambda_D = 3$ GeV)	Our results ($\Lambda_D = 3$ GeV)
$^{40}_{J/\psi}\text{Ca}$	1s	-17.2	-5.44	-8.56
	1p	-12.9	-2.32	-5.054
	1d	-8.21	...	-1.90
	2s	-7.48	...	-2.33
$^{90}_{J/\psi}\text{Zr}$	1s	-18.69	-6.40	-9.98
	1p	-16.07	-4.42	-7.30
	1d	-13.06	-2.13	-4.77
	2s	-12.22	-1.56	-4.96
$^{208}_{J/\psi}\text{Pb}$	1s	-19.1	-7.08	-11.14
	1p	-17.59	-5.86	-9.23
	1d	-15.81	-4.38	-7.36
	2s	-15.26	-3.81	-7.43

VI. SUMMARY

In this work, we have examined how the contribution of D and D^* meson loops in an effective Lagrangian framework affects properties of J/ψ meson in asymmetric nuclear matter at zero and finite temperatures. The hadronic chiral SU(3) model is employed to calculate scalar condensates, which are then incorporated in the QCD sum rules method to estimate the medium dependence of D (D^+ , D^0 , D^{*+} , and D^{*0}) meson masses. The in-medium reduction in masses of vector D^* mesons is more significant compared to pseudoscalar D mesons at finite temperature and density of medium, which is attributed to the stronger attractive interaction of the D^* meson with the nuclear medium. Our analysis reveals that the J/ψ meson exhibits a negative mass shift correlated with an increase in baryonic density. At a temperature of 100 MeV and isospin asymmetry $I_a = 0.3$, the results derived from sum rules indicate that the mass shifts of the D and D^* mesons are observed to be -2.66 and -4.68 MeV, respectively. This leads to an approximate reduction of -3.39 and -7.32 MeV in the J/ψ mass for the DD and DD^* meson loop contributions within asymmetric nuclear matter at

nuclear saturation density, with cutoff parameter $\Lambda_D = 2$ GeV, respectively. In examining the impact of the D^*D^* meson loop on the J/ψ self-energy, we noted a much larger negative mass shift of the J/ψ meson. However, the heavier loop is expected to contribute less. The observed negative mass shift suggests that nuclear mean fields exert an attractive influence on the J/ψ meson, such a mass shift may be sufficient to facilitate the formation of meson-nucleus bound states. Moreover, we have computed the binding energy and absorption decay width of J/ψ meson for both the ground and excited states of O^{16} , Ca^{40} , Zr^{90} , and Pb^{208} nuclei by solving the Klein-Gordon equation. Our finding indicates that the J/ψ meson forms bound states with all the nuclei we examined, and heavier mesons exhibit a greater probability of forming a mesic-nuclei bound state. Mainly, the theoretical predictions of J/ψ - Pb^{208} bound state strongly encourage experimental studies of nuclei after the upgraded 12 GeV CEBAF at JLab facility [28]. This study may serve as a useful reference for experimental studies of quarkonium suppression in heavy-ion collisions and medium effects through the formation of bound states of mesons, associated with the FAIR project, such as CBM and PANDA [22–24], and J-PARC collaboration [29], and CEBAF experiment at JLab [28].

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

- [1] A. Hosaka, T. Hyodo, K. Sudoh, Y. Yamaguchi, and S. Yasui, *Prog. Part. Nucl. Phys.* **96**, 88 (2017).
- [2] G. Krein, *AIP Conf. Proc.* **1701**, 020012 (2016).
- [3] S. Leupold, V. Metag, and U. Mosel, *Int. J. Mod. Phys. E* **19**, 147 (2010).
- [4] R. S. Hayano and T. Hatsuda, *Rev. Mod. Phys.* **82**, 2949 (2010).

- [5] S. A. Mir, S. Uddin, S. K. Tiwari, and M. Faizal, *Fortschr. Phys.* **74**, 3 (2026).
- [6] J. J. Cobos-Martínez, G. N. Zeminiani, and K. Tsushima, *Symmetry* **17**, 787 (2025).
- [7] T. Matsui and H. Satz, *Phys. Lett. B* **178**, 4162 (1986).
- [8] A. Bazavov and J. H. Weber, *Prog. Part. Nucl. Phys.* **116**, 103823 (2021).

- [9] G. Calucci and E. Cattaruzza, *Eur. Phys. J. C* **46**, 439 (2006).
- [10] H.-S. Shao, *Phys. Rev. D* **102**, 034023 (2020).
- [11] N. Brambilla *et al.*, *Eur. Phys. J. C* **71**, 1534 (2011).
- [12] A. Mocsy and P. Petreczky, *Phys. Rev. D* **73**, 074007 (2006).
- [13] A. Andronic *et al.*, *Eur. Phys. J. A* **60**, 88 (2024).
- [14] X. Du and R. Rapp, *Nucl. Phys. A* **943**, 147 (2015).
- [15] P. Braun-Munzinger and J. Stachel, *Landolt- Bornstein* **23**, 424 (2010).
- [16] P. Braun-Munzinger and J. Stachel, *Phys. Lett. B* **490**, 196 (2000).
- [17] K. J. Eskola, P. Paakkinen, H. Paukkunen, and C. A. Salgado, *Eur. Phys. J. C* **77**, 163 (2017).
- [18] E. G. Ferreira, *Phys. Lett. B* **749**, 98 (2015).
- [19] A. Capella and E. G. Ferreira, *Eur. Phys. J. C* **42**, 419 (2005).
- [20] N. Brambilla, M. A. Escobedo, J. Ghiglieri, and A. Vairo, *J. High Energy Phys.* **05** (2013) 130.
- [21] S. Chen and M. He, *Phys. Lett. B* **786**, 260 (2018).
- [22] P. Senger, *Central Eur. J. Phys.* **10**, 1289 (2012).
- [23] K. Agarwal, *Phys. Scr.* **98**, 034006 (2023).
- [24] M. Durante *et al.*, *Phys. Scr.* **94**, 033001 (2019).
- [25] E. Tomas-Gustafsson, *Hyperfine Interact.* **239**, 30 (2018).
- [26] A. Belias, *J. Instrum.* **15**, C10001 (2020).
- [27] L. Frankfurt, M. Strikman, A. Larionov, A. Lehrach, R. Maier, H. van Hees, C. Spieles, V. Vovchenko, and H. Stoecker, *Eur. Phys. J. A* **56**, 171 (2020).
- [28] P. A. Adderley *et al.*, *Phys. Rev. Accel. Beams* **27**, 084802 (2024).
- [29] K. Tanida, *J. Phys. Soc. Jpn. Conf. Proc.* **10**, 010019 (2016).
- [30] E. Eichten, K. Gottfried, T. Kinoshita, K. D. Lane, and T.-M. Yan, *Phys. Rev. D* **17**, 3090 (1978).
- [31] F. Karsch, M. T. Mehr, and H. Satz, *Z. Phys. C* **37**, 617 (1988).
- [32] C. Bonati, M. D’Elia, and A. Rucci, *Phys. Rev. D* **92**, 054014 (2015).
- [33] A. Hayashigaki, *Phys. Lett. B* **487**, 96 (2000).
- [34] S. Kim and S. H. Lee, *Nucl. Phys. A* **679**, 517 (2001).
- [35] F. Klingl, S. Kim, S. H. Lee, P. Morath, and W. Weise, *Phys. Rev. Lett.* **82**, 3396 (1999).
- [36] A. Kumar and A. Mishra, *Phys. Rev. C* **82**, 045207 (2010).
- [37] K. Morita and S. H. Lee, *Phys. Rev. C* **85**, 044917 (2012).
- [38] T. Hilger, B. Kampfer, and S. Leupold, *Phys. Rev. C* **84**, 045202 (2011).
- [39] S. Zschocke, T. Hilger, and B. Kampfer, *Eur. Phys. J. A* **47**, 151 (2011).
- [40] R. Chhabra and A. Kumar, *Eur. Phys. J. A* **53**, 105 (2017).
- [41] A. Kumar and R. Chhabra, *Phys. Rev. C* **92**, 035208 (2015).
- [42] R. Kumar and A. Kumar, *J. Phys. Conf. Ser.* **1690**, 012112 (2020).
- [43] L. Tolos, J. Schaffner-Bielich, and A. Mishra, *Phys. Rev. C* **70**, 025203 (2004).
- [44] L. Tolos, A. Ramos, and T. Mizutani, *Phys. Rev. C* **70**, 015207 (2008).
- [45] J. Hofmann and M. F. M. Lutz, *Nucl. Phys. A* **763**, 90 (2005).
- [46] R. Molina, D. Gamermann, E. Oset, and L. Tolos, *Eur. Phys. J. A* **42**, 31 (2009).
- [47] M. Neubert, *Phys. Rep.* **245**, 2596 (1994).
- [48] S. Yasui and K. Sudoh, *Phys. Rev. C* **87**, 015202 (2013).
- [49] S. Yasui and K. Sudoh, *Phys. Rev. C* **89**, 015201 (2014).
- [50] K. Tsushima *et al.*, *Phys. Rev. C* **59**, 2824 (1999).
- [51] K. Tsushima and F. C. Khanna, *Phys. Lett. B* **552**, 138 (2003).
- [52] K. Saito and A. W. Thomas, *Phys. Lett. B* **327**, 9 (1994).
- [53] G. Krein, A. W. Thomas, and K. Tsushima, *Phys. Lett. B* **697**, 136 (2011).
- [54] G. Krein, A. W. Thomas, and K. Tsushima, *Prog. Part. Nucl. Phys.* **100**, 161 (2018).
- [55] K. Tsushima, D. H. Lu, G. Krein, and A. W. Thomas, *AIP Conf. Proc.* **1354**, 39 (2011).
- [56] G. N. Zeminiani, J. J. Cobos-Martinez, and K. Tsushima, *Eur. Phys. J. A* **57**, 259 (2021).
- [57] S. H. Lee and C. M. Ko, *Phys. Rev. C* **67**, 038202 (2003).
- [58] A. B. Kaidalov and P. E. Volkovitsky, *Phys. Rev. Lett.* **69**, 3155 (1992).
- [59] M. E. Luke, A. V. Manohar, and M. J. Savage, *Phys. Lett. B* **288** (1992) 355.
- [60] J. Tarrús Castellà and G. Krein, *Phys. Rev. D* **98**, 014029 (2018).
- [61] A. Sibirtsev and M. B. Voloshin, *Phys. Rev. D* **71**, 076005 (2005).
- [62] M. E. Peskin, *Nucl. Phys. B* **156**, 365 (1979).
- [63] D. Kharzeev, *Proc. International School of Physics “Enrico Fermi”* **130**, 105 (1996).
- [64] S. R. Beane, E. Chang, S. D. Cohen, W. Detmold, H.-W. Lin, K. Orginos, A. Parreño, and M. J. Savage, *Phys. Rev. D* **91**, 114503 (2015).
- [65] M. Alberti, G. S. Bali, S. Collins, F. Knechtli, G. Moir, and W. Söldner, *Phys. Rev. D* **95**, 074501 (2017).
- [66] A. Hayashigaki, *Prog. Theor. Phys.* **101**, 923 (1999).
- [67] J. Cobos Martinez, K. Tsushima, G. Krein, and A. W. Thomas, *Proc. Sci. PANIC2021* (2022) 199.
- [68] P. Papazoglou *et al.*, *Phys. Rev. C* **59**, 411 (1999).
- [69] W. Bardeen and B. Lee, *Phys. Rev. C* **177**, 2389 (1969).
- [70] S. Weinberg, *Phys. Rev.* **166**, 1568 (1968).
- [71] D. Zschesche *et al.*, *Phys. Rev. C* **70**, 045202 (2004).
- [72] A. Kumar and A. Mishra, *Phys. Rev. C* **81**, 065204 (2010).
- [73] A. Mishra and D. Pathak, *Phys. Rev. C* **90**, 025201 (2014).
- [74] A. Mishra, E. L. Bratkovskaya, J. Schaffer-Bielich, S. Schramm, and H. Stöcker, *Phys. Rev. C* **69**, 015202 (2004).
- [75] D. Pathak and A. Mishra, *Phys. Rev. C* **91**, 045206 (2015).
- [76] T. Barnes, F. E. Close, P. R. Page, and E. S. Swanson, *Phys. Rev. D* **55**, 4157 (1997).
- [77] A. Le Yaouanc, L. Oliver, O. Pene, and J. C. Raynal, *Phys. Rev. D* **8**, 2223 (1973).
- [78] T. S. H. Lee, S. Sakinah, and Y. Oh, *Eur. Phys. J. A* **58**, 252 (2022).
- [79] S.-H. Kim, T. S. H. Lee, and R. B. Wiringa, *Phys. Rev. C* **112**, 015206 (2025).
- [80] F. He and I. Zahed, *Phys. Rev. C* **112**, 014303 (2025).
- [81] R. Chhabra and K. Tsushima, *J. Phys. Conf. Ser.* **1690**, 012108 (2020).
- [82] T. Nishi *et al.*, *Nat. Phys.* **19**, 788 (2023).

- [83] V. Metag, M. Nanova, and E. Y. Paryev, *Prog. Part. Nucl. Phys.* **97**, 199 (2017).
- [84] R. Szmajkowski, *Phys. Rev. A* **98**, 042507 (2018).
- [85] C. Garcia-Recio, J. Nieves, and L. Tolos, *Phys. Lett. B* **690**, 369 (2010).
- [86] J. Yamagata-Sekihara, C. Garcia-Recio, J. Nieves, L. L. Salcedo, and L. Tolos, *Phys. Lett. B* **754**, 26 (2016).
- [87] S. Schadmand, *Prog. Theor. Phys. Suppl.* **186**, 373 (2010).
- [88] K. Tsushima, D. H. Lu, G. Krein, and A. W. Thomas, *Phys. Rev. C* **83**, 065208 (2011).
- [89] K. Saito, K. Tsushima, and A. W. Thomas, *Phys. Rev. C* **56**, 566 (1997).
- [90] J. J. Cobos-Martínez, K. Tsushima, G. Krein, and A. W. Thomas, *Phys. Lett. B* **811**, 135882 (2020).
- [91] J. J. Cobos-Martínez, G. N. Zeminiani, and K. Tsushima, *Phys. Rev. C* **105**, 025204 (2022).
- [92] A. Mondal and A. Mishra, *Phys. Rev. C* **110**, 055201 (2024).
- [93] M. Kaur and A. Kumar, *Phys. Rev. D* **110**, 114054 (2024).
- [94] A. Mishra, *Phys. Rev. C* **91**, 035201 (2015).
- [95] Z.-G. Wang, *Phys. Rev. C* **92**, 065205 (2015).
- [96] R. Kumar and A. Kumar, *Phys. Rev. C* **101**, 015202 (2020).
- [97] S. Weinberg, *Phys. Rev.* **166**, 1568 (1968).
- [98] R. Thomas, T. Hilger, and B. Kampfer, *Nucl. Phys.* **A795**, 19 (2007).
- [99] K. Saito, K. Tsushima, and A. W. Thomas, *Mod. Phys. Lett. A* **13**, 769 (1998).
- [100] G.-Q. Li and C. M. Ko, *Phys. Lett. B* **338**, 118 (1994).
- [101] T. D. Cohen, R. J. Furnstahl, and D. K. Griegel, *Phys. Rev. C* **45**, 1881 (1992).
- [102] R. Chhabra and A. Kumar, *Eur. Phys. J. C* **77**, 726 (2017).
- [103] Z.-G. Wang, *Int. J. Mod. Phys. A* **28**, 1350049 (2013).
- [104] A. Kumar and A. Mishra, *Chin. Phys. C* **49**, 074106 (2025).
- [105] M. Kaur and A. Kumar, *Phys. Rev. D* **112**, 014030 (2025).
- [106] T. Hilger, R. Thomas, and B. Kampfer, *Phys. Rev. C* **79**, 025202 (2009).
- [107] R. Chhabra and A. Kumar, *EPJ Web Conf.* **112**, 04001 (2016).
- [108] C. E. Jimenez-Tejero, A. Ramos, L. Tolos, and I. Vidana, *Phys. Rev. C* **84**, 015208 (2011).
- [109] J. S. Schwinger, *J. Math. Phys. (N.Y.)* **2**, 407 (1961).
- [110] L. V. Keldysh, *Sov. Phys. JETP* **20**, 1018 (1965).
- [111] K.-c. Chou, Z.-b. Su, B.-l. Hao, and L. Yu, *Phys. Rep.* **118**, 1 (1985).
- [112] Z.-w. Lin and C. M. Ko, *Phys. Rev. C* **62**, 034903 (2000).
- [113] J. Haidenbauer, G. Krein, U.-G. Meissner, and A. Sibirtsev, *Eur. Phys. J. A* **33**, 107 (2007).
- [114] J. J. Sakurai, *Ann. Phys. (N.Y.)* **11**, 1 (1960).
- [115] Z.-w. Lin and C. M. Ko, *Phys. Lett. B* **503**, 104 (2001).
- [116] D. B. Leinweber, A. W. Thomas, K. Tsushima, and S. V. Wright, *Phys. Rev. D* **64**, 094502 (2001).
- [117] K. Tsushima, *Phys. Rev. D* **99**, 014026 (2019).
- [118] C. Baglin *et al.*, *Can. J. Phys.* **67**, 1222 (1989).
- [119] C. Baglin *et al.*, *Phys. Lett. B* **345**, 617 (1995).
- [120] R. Arnaldi *et al.*, *Nucl. Phys.* **A783**, 261 (2007).
- [121] A. Adare *et al.*, *Phys. Rev. Lett.* **98**, 232301 (2007).
- [122] A. Adare *et al.*, *Phys. Rev. C* **84**, 054912 (2011).
- [123] B. Abelev *et al.*, *Phys. Rev. Lett.* **109**, 072301 (2012).
- [124] C. W. De Jager, H. De Vries, and C. De Vries, *At. Data Nucl. Data Tables* **14**, 479 (1974).
- [125] J. Nieves, E. Oset, and C. Garcia-Recio, *Nucl. Phys.* **A554**, 509 (1993).
- [126] A. W. Steiner, M. Prakash, J. M. Lattimer, and P. J. Ellis, *Phys. Rep.* **411**, 325 (2005).
- [127] A. W. Steiner, M. Prakash, J. M. Lattimer, and P. J. Ellis, *Acta Phys. Hung.* **A 25**, 203 (2006).
- [128] J. J. Cobos-Martínez and K. Tsushima, *Phys. Rev. C* **109**, 025202 (2024).
- [129] K. Tsushima, *Nucl. Phys.* **A670**, 198 (2000).
- [130] E. Friedman and A. Gal, *Phys. Rep.* **452**, 89 (2007).
- [131] Y. Oh, W. Liu, and C. M. Ko, *Phys. Rev. C* **75**, 064903 (2007).
- [132] A. Sibirtsev, K. Tsushima, and A. W. Thomas, *Phys. Rev. C* **63**, 044906 (2001).
- [133] R. H. Landau, *Phys. Rev. C* **27**, 2191 (1983).
- [134] Y. R. Kwan and F. Tabakin, *Phys. Rev. C* **18**, 932 (1978).
- [135] K. Tsushima, D.-H. Lu, A. W. Thomas, and K. Saito, *Phys. Lett. B* **443**, 26 (1998).