J/ψ and η_c masses in isospin asymmetric hot nuclear matter: A QCD sum rule approach

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We study the in-medium masses of the charmonium states J/ψ and η_c in the nuclear medium using the QCD sum rule approach. These mass modifications arise owing to modifications of the scalar and the twist-2 gluon condensates in the hot hadronic matter. The scalar gluon condensate $\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{a\mu\nu} \rangle$ and the twist-2 tensorial gluon operator $\langle \frac{\alpha_s}{\pi} G^a_{\mu\sigma} G^a{}_\nu{}^\sigma \rangle$ in the nuclear medium are calculated from the medium modification of a scalar dilaton field introduced to incorporate trace anomalies of QCD within the chiral SU(3) model used in the present investigation. The effects of isospin asymmetry, density, and temperature of the nuclear medium on the in-medium masses of the lowest charmonium states J/ψ and η_c mesons are investigated in the present work. The results of the present investigation are compared with the existing results on the masses of these states. The medium modifications of the masses of these charmonium states (J/ψ and η_c) seem to be appreciable at high densities and should modify the experimental observables arising from the compressed baryonic matter produced in asymmetric heavy-ion collision experiments at the future facility of Facility for Antiproton and Ion Research, GSI.

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I. INTRODUCTION

The study of in-medium hadronic properties is of considerable interest, both experimentally and theoretically in present-day strong-interaction physics. The study of the in-medium properties of hadrons has direct relevance in experiments where hadronic matter is probed at high densities and/or temperatures. The compressed baryonic matter (CBM) experiment at the Facility for Antiproton and Ion Research (FAIR), GSI Helmholtz Centre for Heavy Ion Research (GSI), is planned to produce dense matter at high densities and moderate temperatures. The medium modifications of the strange and charm mesons and their effects on the experimental observables are among the topics that are intended to be studied extensively in these experiments. Therefore, the topic of the study of charm mesons in the medium has gotten considerable interest in the recent past. The medium modifications of the properties of the charm mesons D and $\bar{D}$, as well as the excited charmonium states, can have important consequences on the production of open charm and the suppression of the J/ψ in heavy-ion collision experiments. The suppression of J/ψ in heavy-ion collisions may lead to the signature of quark-gluon plasma (QGP) [1,2]. Also, it is observed that the effect of hadron absorption of J/ψ is not negligible [3–5]. In Ref. [6], it was reported that the charmonium suppression observed in Pb + Pb collisions in the NA50 experiment cannot be simply explained by nucleon absorption, but needs some additional density-dependent suppression mechanisms. It was suggested in these studies that the comover scattering [6–8] can explain the additional suppression of charmonium. An important difference between J/ψ suppression pattern in comovers interaction model and in a deconfining scenario is that, in the former case, the anomalous suppression sets in smoothly from peripheral to central collisions rather than in a sudden way when the deconfining threshold is reached [7]. The J/ψ suppression in nuclear collisions at Super Proton Synchrotron (SPS) energies has been studied in the covariant transport approach hadron-string dynamics (HSD) in Ref. [8]. The calculations show that the absorption of J/ψ's by both nucleons and produced mesons can explain reasonably not only the total J/ψ cross section but also the transverse energy dependence of J/ψ suppression measured in both proton-nucleus and nucleus collisions. In Ref. [9], the cross section of J/ψ dissociation by gluons is used to calculate the J/ψ suppression in an equilibrating parton gas produced in high-energy nuclear collisions. The large average momentum in the hot gluon gas enables gluons to break up the J/ψ, while hadron matter at reasonable temperature does not provide sufficiently hard gluons. The multigluon exchange can lead to an attractive potential between a $c\bar{c}$ meson and a nucleon, such that, for example, the η_c could form bound states even with light nuclei [10,11].

The D ($\bar{D}$) mesons are made up of a light (u or d) antiquark (quark) and one heavy charm quark (charm antiquark). In QCD sum rule calculations, the mass modifications of $D(\bar{D})$ mesons in the nuclear medium arise owing to interactions of light antiquark (quark) present in the $D(\bar{D})$ mesons with the light quark condensate [12,13]. There is appreciable change in the light-quark condensate in the nuclear medium and, hence, $D(\bar{D})$ meson mass, owing to its interaction with the light-quark condensate, changes appreciably in the hadronic matter. The medium modifications of the D mesons modify the decay widths of the charmonium states, which have been studied in Ref. [13]. The charmonium states are made up of a heavy charm quark and a charm antiquark. Within QCD sum rule calculations, it is suggested that these heavy quarkonium states interact with the nuclear medium through the gluon condensates [10], unlike the interaction of the light vector mesons with the nuclear medium, which is through the light-quark condensates [14]. This is because all the

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heavy-quark condensates can be related to the gluon condensates via heavy-quark expansion [15]. Also in the nuclear medium there are no valence charm quarks to leading order in density and any interaction with the medium is gluonic. The medium modifications of the gluon condensates are seen to be small and this leads to the mass modifications of J/ψ and η_c mesons, which are the lowest charmonium states, to be small in the nuclear medium [10]. The leading-order perturbative calculations [16] in the study of the charmonium states also shows that the mass of J/ψ is reduced slightly in the nuclear medium. In Ref. [17], the mass modifications of the charmonium states have been studied using QCD second-order Stark effect and the linear density approximation for the gluon condensate in the nuclear medium. This shows a small drop for the J/ψ mass at the nuclear-matter density, but there is seen to be significant shift in the masses of the excited states of charmonium [$\psi(3686)$ and $\psi(3770)$]. Using QCD second-order Stark effect, the masses of the charmonium states were also studied [18] in the asymmetric nuclear medium at finite temperatures. These medium modifications were investigated by computing the scalar gluon condensate in the hot nuclear medium from the medium modification of a scalar dilaton field within a chiral SU(3) model, which was introduced to incorporate broken scale invariance of QCD. This investigation showed a small drop in the J/ψ mass in the medium, whereas the masses of the excited charmonium states were observed to have appreciable drop at high densities.

In the present investigation, we study the in-medium modifications of the vector meson, J/ψ and the pseudoscalar meson, η_c, using QCD sum rules [10] and an effective chiral SU(3) model [19]. To apply the QCD sum rules for the study of in-medium modifications of J/ψ and η_c mesons, we consider the contributions of the scalar gluon condensates $\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{a\mu\nu}\rangle$ and twist-2 tensorial gluon operator $\langle \frac{\alpha_s}{\pi} G^a_{\mu\sigma} G^{a}{}_{\nu}{}^{\sigma}\rangle$ up to dimension four [10]. The scalar gluon condensate, as well as the twist-2 gluon operator in the nuclear medium, are calculated from the medium modification of a scalar dilaton field, χ, introduced within a chiral SU(3) model [19] through a scale symmetry-breaking term in the Lagrangian density leading to the QCD trace anomaly. The chiral SU(3) model [19] has been used successfully to study the medium modifications of kaons and antikaons in isospin asymmetric nuclear matter in Ref. [20] and in hyperonic matter in Ref. [21]. The chiral SU(3) model was generalized to SU(4) to study the mass modifications of D mesons arising from their interactions with the light hadrons in isospin symmetric hot hadronic matter in Ref. [22] and in isospin asymmetric nuclear matter at zero temperature [23] and finite temperatures [18], respectively. The in-medium properties of the vector mesons have also been studied within the model [24,25]. In the present investigation, we study the in-medium masses of the J/ψ and η_c mesons, calculated from the medium modifications of the dilaton field, χ, in the isospin asymmetric nuclear matter at finite temperatures within the chiral SU(3) model.

The outline of the article is as follows. In Sec. II, we give a brief introduction of the chiral SU(3) model used to study the in-medium masses of charmonium states J/ψ and η_c in the present investigation. The medium modifications of these charmonium states, J/ψ and η_c mesons, arise from the medium modification of the scalar gluon condensate in the nuclear medium, simulated by a scalar dilaton field introduced in the hadronic model to incorporate broken scale invariance of QCD leading to QCD trace anomaly and also owing to the medium modification of the expectation value of the twist-2 gluon operator. Section III discusses briefly the QCD sum rule approach used to calculate the masses of the charmonium states J/ψ and η_c. In Sec. IV, we discuss the results of the present investigation. Section V summarizes the conclusions of the present work.

II. THE HADRONIC CHIRAL SU(3) × SU(3) MODEL

We use an effective chiral SU(3) model for the present investigation [19]. The model is based on the nonlinear realization of chiral symmetry [26–28] and broken scale invariance [19,24,25]. This model has been used successfully to describe nuclear matter, finite nuclei, hypernuclei and neutron stars. The effective hadronic chiral Lagrangian density contains the following terms:

$$\mathcal{L} = \mathcal{L}_{\text{kin}} + \sum_{W=X,Y,V,A,u} \mathcal{L}_{\text{BW}} + \mathcal{L}_{\text{vec}} + \mathcal{L}_0 + \mathcal{L}_{\text{SB}}. \quad (1)$$

In Eq. (1), $\mathcal{L}_{\text{kin}}$ is the kinetic-energy term, $\mathcal{L}_{\text{BW}}$ is the baryon-meson-interaction term in which the baryon–spin-0–meson interaction term generates the vacuum baryon masses. $\mathcal{L}_{\text{vec}}$ describes the dynamical mass generation of the vector mesons via couplings to the scalar mesons and contains additionally quartic self-interactions of the vector fields. $\mathcal{L}_0$ contains the meson-meson interaction terms inducing the spontaneous breaking of chiral symmerty, as well as a scale-invariance-breaking logarthimic potential. $\mathcal{L}_{\text{SB}}$ describes the explicit chiral symmetry breaking.

To study the hadron properties at finite temperature and densities in the present investigation, we use the mean-field approximation, where all the meson fields are treated as classical fields. In this approximation, only the scalar and vector fields contribute to the baryon-meson interaction $\mathcal{L}_{\text{BW}}$ because for all the other mesons the expectation values are zero. The interactions of the scalar mesons and vector mesons with the baryons are given as

$$\mathcal{L}_{\text{Bscal}} + \mathcal{L}_{\text{Bvec}} = -\sum_i \bar{\psi}_i [m_i^* + g_{\omega i}\gamma_0 \omega + g_{\rho i}\gamma_0 \rho + g_{\phi i}\gamma_0 \phi]\psi_i. \quad (2)$$

The interaction of the vector mesons and of the scalar fields and the interaction corresponding to the explicit symmetry breaking, in the mean field approximation, are given as

$$\mathcal{L}_{\text{vec}} = \frac{1}{2}\left(m_\omega^2 \omega^2 + m_\rho^2 \rho^2 + m_\phi^2 \phi^2\right)\frac{\chi^2}{\chi_0^2} + g_4(\omega^4 + 6\omega^2\rho^2 + \rho^4 + 2\phi^4), \quad (3)$$

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

and

$$\mathcal{L}_{SB} = -\left(\frac{\chi}{\chi_0}\right)^2\left[m_\pi^2 f_\pi\sigma+\left(\sqrt{2}m_k^2 f_k-\frac{1}{\sqrt{2}}m_\pi^2 f_\pi\right)\zeta\right]. \quad (5)$$

The effective mass of the baryon of species i is given as

$$m_i^* = -(g_{\sigma i}\sigma+g_{\zeta i}\zeta+g_{\delta i}\delta). \quad (6)$$

The baryon-scalar meson interactions, as can be seen from Eq. (6), generate the baryon masses through the coupling of baryons to the nonstrange σ, the strange ζ scalar mesons, and the scalar-isovector meson δ. In analogy to the baryon-scalar meson couplings, there exist two independent baryon-vector meson interaction terms corresponding to the F-type (antisymmetric) and D-type (symmetric) couplings. Here antisymmetric coupling is used because the universality principle [29] and vector meson dominance model suggest small symmetric couplings. Additionally, we choose the parameters [19,20] so as to decouple the strange vector field $\phi_\mu \sim \bar{s}\gamma_\mu s$ from the nucleon, corresponding to an ideal mixing between ω and ϕ mesons. A small deviation of the mixing angle from ideal mixing [30–32] has not been taken into account in the present investigation.

The concept of broken scale invariance leading to the trace anomaly in (massless) QCD, $\theta_\mu^\mu = \frac{\beta_{QCD}}{2g}G_{\mu\nu}^a G^{\mu\nu a}$, where $G_{\mu\nu}^a$ is the gluon field strength tensor of QCD, is simulated in the effective Lagrangian at tree level [33] through the introduction of the scale-breaking terms

$$\mathcal{L}_{\text{scale breaking}} = -\frac{1}{4}\chi^4\ln\left(\frac{\chi^4}{\chi_0^4}\right) + \frac{d}{3}\chi^4\ln\left[\left(\frac{I_3}{\det\langle X\rangle_0}\right)\left(\frac{\chi}{\chi_0}\right)^3\right], \quad (7)$$

where $I_3 = \det\langle X\rangle$, with X as the multiplet for the scalar mesons. These scale-breaking terms, in the mean-field approximation, are given by the last two terms of the Lagrangian density, $\mathcal{L}_0$ given by Eq. (4) [34]. Within the chiral SU(3) model used in the present investigation, the scalar gluon condensate $\langle\frac{\alpha_s}{\pi}G_{\mu\nu}^a G^{\mu\nu a}\rangle$, as well as the twist-2 gluon operator $\langle\frac{\alpha_s}{\pi}G_{\mu\sigma}^a G_\nu^{\sigma a}\rangle$, are simulated by the scalar dilaton field, χ. These are obtained from the energy momentum tensor

$$T_{\mu\nu} = (\partial_\mu\chi)\left(\frac{\partial\mathcal{L}_\chi}{\partial(\partial^\nu\chi)}\right)-g_{\mu\nu}\mathcal{L}_\chi, \quad (8)$$

derived from the Lagrangian density for the dilaton field, given as

$$\mathcal{L}_\chi = \frac{1}{2}(\partial_\mu\chi)(\partial^\mu\chi)-k_4\chi^4-\frac{1}{4}\chi^4\ln\left(\frac{\chi^4}{\chi_0^4}\right) + \frac{d}{3}\chi^4\ln\left[\left(\frac{(\sigma^2-\delta^2)\zeta}{\sigma_0^2\zeta_0}\right)\left(\frac{\chi}{\chi_0}\right)^3\right]. \quad (9)$$

In massless QCD, the energy momentum tensor can be written as [35,36]

$$T_{\mu\nu} = -ST\left(G_{\mu\sigma}^a G_\nu^{a\sigma}\right)+\frac{g_{\mu\nu}}{4}\frac{\beta_{QCD}}{2g}G_{\sigma\kappa}^a G^{a\sigma\kappa}, \quad (10)$$

where the first term is the symmetric traceless part and second term is the trace part of the energy momentum tensor. Writing

$$\left\langle\frac{\alpha_s}{\pi}G_{\mu\sigma}^a G_\nu^{a\sigma}\right\rangle = \left(u_\mu u_\nu-\frac{g_{\mu\nu}}{4}\right)G_2, \quad (11)$$

where u_μ is the four velocity of the nuclear medium, taken as $u_\mu = (1,0,0,0)$, we obtain the energy momentum tensor in QCD as

$$T_{\mu\nu} = -\left(\frac{\pi}{\alpha_s}\right)\left(u_\mu u_\nu-\frac{g_{\mu\nu}}{4}\right)G_2+\frac{g_{\mu\nu}}{4}\frac{\beta_{QCD}}{2g}G_{\sigma\kappa}^a G^{a\sigma\kappa}. \quad (12)$$

Equating the energy-momentum tensors given by Eqs. (8) and (12) and multiplying by $(u^\mu u^\nu - \frac{g^{\mu\nu}}{4})$, we obtain the expression for G_2 as

$$G_2 = -\frac{\alpha_s}{\pi}\left[(\partial_\alpha\chi)\left(\frac{\partial\mathcal{L}_\chi}{\partial(\partial_\alpha\chi)}\right)+\frac{4}{3}(\partial_i\chi)(\partial_i\chi)\right]. \quad (13)$$

We note here that by multiplying the energy momentum tensor of QCD given through Eq. (12), by $(u^\mu u^\nu - \frac{g^{\mu\nu}}{4})$, we project out the traceless part given by the first term of the energy momentum tensor, described by the function G_2. This is because $g_{\mu\nu}(u^\mu u^\nu - \frac{g^{\mu\nu}}{4}) = 0$, and, hence, there is no contribution from the trace part of the energy momentum tensor in QCD, when we multiply the same by $(u^\mu u^\nu - \frac{g^{\mu\nu}}{4})$. Similarly, by multiplying the energy momentum tensor given by Eq. (12) by $g^{\mu\nu}$, the first part gives zero and only the second term contributes to the trace of the energy momentum tensor. The effect of the logarithmic terms in the chiral SU(3) model, given by Eq. (9), is to break the scale invariance. Multiplying Eq. (8) by $g^{\mu\nu}$, we obtain the trace of the energy momentum tensor within the chiral SU(3) model as

$$T_\mu^\mu = (\partial_\mu\chi)\left(\frac{\partial\mathcal{L}_\chi}{\partial(\partial_\mu\chi)}\right)-4\mathcal{L}_\chi. \quad (14)$$

Using the Euler-Lagrange's equation for the χ field, the trace of the energy momentum tensor in the chiral SU(3) model can be expressed as [18,34]

$$T_\mu^\mu = \chi\frac{\partial\mathcal{L}_\chi}{\partial\chi}-4\mathcal{L}_\chi = -(1-d)\chi^4. \quad (15)$$

Multiplying Eq. (12) by $g^{\mu\nu}$, we obtain the trace of the energy momentum tensor in QCD as

$$T_\mu^\mu = \left\langle\frac{\beta_{QCD}}{2g}G_{\sigma\kappa}^a G^{a\sigma\kappa}\right\rangle. \quad (16)$$

Using the Euler-Lagrange equation for χ and dropping a total divergence term in Eq. (13), the expression for G_2 can be written as

$$G_2 = \frac{\alpha_s}{\pi}\left[\chi\frac{\partial\mathcal{L}_\chi}{\partial\chi} - \frac{4}{3}(\partial_i\chi)(\partial_i\chi)\right]$$
$$= \frac{\alpha_s}{\pi}\left\{-(1-d+4k_4)\chi^4 - \chi^4\ln\left(\frac{\chi^4}{{\chi_0}^4}\right) + \frac{4}{3}d\chi^4\ln\left[\left(\frac{(\sigma^2-\delta^2)\zeta}{\sigma_0^2\zeta_0}\right)\left(\frac{\chi}{\chi_0}\right)^3\right] - \frac{4}{3}(\partial_i\chi)(\partial_i\chi)\right\}. \quad (17)$$

The twist-2 gluon operator has a contribution only in the nuclear medium and is zero in vacuum [10]. Hence $(G_2)_{\rm vac}=0$, which implies that

$$-(1-d+4k_4){\chi_0}^4 - \tfrac{4}{3}\langle(\partial_i\chi)(\partial_i\chi)\rangle_{\rm vac} = 0. \quad (18)$$

Assuming the glueball field χ to be nonrelativistic, and hence assuming that $\langle(\partial_i\chi)(\partial_i\chi)\rangle_{\rm medium} \simeq \langle(\partial_i\chi)(\partial_i\chi)\rangle_{\rm vac}$ and using Eq. (18), the expression for G_2 is obtained from Eq. (17) as

$$G_2 = \frac{\alpha_s}{\pi}\left\{-(1-d+4k_4)(\chi^4-{\chi_0}^4) - \chi^4\ln\left(\frac{\chi^4}{{\chi_0}^4}\right) + \frac{4}{3}d\chi^4\ln\left[\left(\frac{(\sigma^2-\delta^2)\zeta}{\sigma_0^2\zeta_0}\right)\left(\frac{\chi}{\chi_0}\right)^3\right]\right\}. \quad (19)$$

The scalar gluon condensate and the twist-2 gluon operator, described in terms of the function G_2 given by Eqs. (15) and (19), are thus related to the χ field, which is solved from the coupled equations of motion of the scalar fields within the chiral SU(3) model.

The coupled equations of motion for the nonstrange scalar field σ, the strange scalar field ζ, the scalar-isovector field δ, and the dilaton field χ are derived from the Lagrangian density and are given as

$$k_0\chi^2\sigma - 4k_1(\sigma^2+\zeta^2+\delta^2)\,\sigma - 2k_2(\sigma^3+3\sigma\delta^2) - 2k_3\chi\sigma\zeta - \frac{d}{3}\chi^4\left(\frac{2\sigma}{\sigma^2-\delta^2}\right) + \left(\frac{\chi}{\chi_0}\right)^2 m_\pi^2 f_\pi - \sum g_{\sigma i}\rho_i^s = 0, \quad (20)$$

$$k_0\chi^2\zeta - 4k_1(\sigma^2+\zeta^2+\delta^2)\zeta - 4k_2\zeta^3 - k_3\chi(\sigma^2-\delta^2) - \frac{d}{3}\frac{\chi^4}{\zeta} + \left(\frac{\chi}{\chi_0}\right)^2\left[\sqrt{2}m_k^2 f_k - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi\right] - \sum g_{\zeta i}\rho_i^s = 0, \quad (21)$$

$$k_0\chi^2\delta - 4k_1(\sigma^2+\zeta^2+\delta^2)\,\delta - 2k_2(\delta^3+3\sigma^2\delta) + k_3\chi\delta\zeta + \frac{2}{3}d\chi^4\left(\frac{\delta}{\sigma^2-\delta^2}\right) - \sum g_{\delta i}\rho_i^s = 0, \quad (22)$$

$$k_0\chi(\sigma^2+\zeta^2+\delta^2) - k_3(\sigma^2-\delta^2)\zeta + \chi^3\left[1+\ln\left(\frac{\chi^4}{\chi_0^4}\right)\right] + (4k_4-d)\chi^3 - \frac{4}{3}d\chi^3\ln\left[\left(\frac{(\sigma^2-\delta^2)\zeta}{\sigma_0^2\zeta_0}\right)\left(\frac{\chi}{\chi_0}\right)^3\right] + \frac{2\chi}{\chi_0^2}\left[m_\pi^2 f_\pi\sigma + \left(\sqrt{2}m_k^2 f_k - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi\right)\zeta\right] = 0. \quad (23)$$

In the preceding, ${\rho_i}^s$ are the scalar densities for the baryons, given as

$$\rho_i^s = \gamma_i\int\frac{d^3k}{(2\pi)^3}\frac{m_i^*}{E_i^*(k)}\times\left(\frac{1}{e^{[{E_i}^*(k)-{\mu_i}^*]/T}+1} + \frac{1}{e^{[{E_i}^*(k)+{\mu_i}^*]/T}+1}\right), \quad (24)$$

where ${E_i}^*(k) = (k^2+{m_i}^{*2})^{1/2}$ and ${\mu_i}^* = \mu_i - g_{\omega i}\omega - g_{\rho i}\rho - g_{\phi i}\phi$ are the single-particle energy and the effective chemical potential for the baryon of species i and, $\gamma_i = 2$ is the spin degeneracy factor [20].

The preceding coupled equations of motion are solved to obtain the density- and temperature-dependent values of the scalar fields (σ, ζ, and δ) and the dilaton field, χ, in the isospin asymmetric hot nuclear medium. As has been mentioned, the value of the χ is related to the scalar gluon condensate, as well as the twist-2 gluon operator in the hot hadronic medium, and is used to compute the in-medium masses of charmonium states in the present investigation. The isospin asymmetry in the medium is introduced through the scalar-isovector field δ and therefore the dilaton field obtained after solving the preceding equations is also dependent on the isospin asymmetry parameter η, defined as $\eta = (\rho_n-\rho_p)/(2\rho_B)$, where ρ_n and ρ_p are the number densities of the neutron and the proton and ρ_B is the baryon density. In the present investigation, we study the effect of isospin asymmetry of the medium on the masses of the charmonium states J/ψ and η_c.

The comparison of the trace of the energy momentum tensor arising from the trace anomaly of QCD with that of the present chiral model given by Eqs. (15) and (16) gives the relation of the dilaton field to the scalar gluon condensate. We have, in the limit of massless quarks [37],

$$T_\mu^\mu = \left\langle\frac{\beta_{\rm QCD}}{2g}G_{\mu\nu}^a G^{\mu\nu a}\right\rangle \equiv -(1-d)\chi^4. \quad (25)$$

In the case of finite quark masses, Eq. (25) gets modified to

$$T_\mu^\mu = \sum_i m_i\bar{q}_i q_i + \left\langle\frac{\beta_{\rm QCD}}{2g}G_{\mu\nu}^a G^{\mu\nu a}\right\rangle \equiv -(1-d)\chi^4, \quad (26)$$

where the first term of the energy-momentum tensor within the chiral SU(3) model is the negative of the explicit chiral symmetry-breaking term $\mathcal{L}_{\rm SB}$ given by Eq. (5).

The parameter d in Eq. (26) originates from the second logarithmic term of Eq. (7). To get an insight into the value of the parameter d, we recall that the QCD β function at the one-loop level for N_c colors and N_f flavors is given by

$$\beta_{\rm QCD}(g) = -\frac{11 N_c g^3}{48\pi^2}\left(1 - \frac{2N_f}{11 N_c}\right) + O(g^5). \tag{27}$$

In the preceding equation, the first term in the parentheses arises from the (antiscreening) self-interaction of the gluons and the second term, proportional to N_f, arises from the (screening) contribution of quark pairs. For massless quarks, Eqs. (25) and (27) suggest the value of d to be 6/33 for three flavors and three colors, and for the case of three colors and two flavors, the value of d turns out to be 4/33, consistent with the one-loop estimate of the QCD β function. These values give the order of magnitude about which the parameter d can be taken [34] because one cannot rely on the one-loop estimate for $\beta_{\rm QCD}(g)$. In the present investigation of the in-medium properties of the charmonium states owing to the medium modification of the dilaton field within chiral SU(3) model, we use the value of $d = 0.064$ [23]. This parameter, along with the other parameters corresponding to the scalar Lagrangian density $\mathcal{L}_0$ given by Eq. (4), are fitted so as to ensure extrema in the vacuum for the σ, ζ, and χ field equations to reproduce the vacuum masses of the η and η' mesons, the mass of the σ meson around 500 MeV, and pressure, $p(\rho_0) = 0$, with ρ_0 as the nuclear-matter saturation density [19,23].

The trace of the energy-momentum tensor in QCD, using the one-loop β function given by Eq. (27) for $N_c = 3$ and $N_f = 3$, and accounting for the finite quark masses [37] is given as

$$T_\mu^\mu = -\frac{9}{8}\frac{\alpha_s}{\pi} G^a_{\mu\nu}\, G^{a\mu\nu} + \left(\frac{\chi}{\chi_0}\right)^2 \left[m_\pi^2 f_\pi \sigma + \left(\sqrt{2} m_k^2 f_k - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi\right)\zeta\right]. \tag{28}$$

Using Eqs. (25) and (28), we can write

$$\left\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu}\, G^{a\mu\nu}\right\rangle = \frac{8}{9}\left\{(1-d)\chi^4 + \left(\frac{\chi}{\chi_0}\right)^2\left[m_\pi^2 f_\pi \sigma + \left(\sqrt{2} m_k^2 f_k - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi\right)\zeta\right]\right\}. \tag{29}$$

We thus see from Eq. (29) that the scalar gluon condensate $\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{\mu\nu a}\rangle$ is related to the dilaton field χ. For massless quarks, because the second term in Eq. (29) arising from explicit symmetry breaking is absent, the scalar gluon condensate becomes proportional to the fourth power of the dilaton field, χ, in the chiral SU(3) model. As mentioned earlier, the in-medium masses of charmonium states are modified owing to the scalar gluon condensate and the twist-2 gluon operators, which are calculated from the modification of the χ field.

III. QCD SUM RULE APPROACH AND IN-MEDIUM MASSES OF J/ψ AND η_c

In the present section, we use the medium modifications of the gluon condensate, calculated from the dilaton field in the chiral effective model to compute the masses of the charmonium states J/ψ and η_c in isospin asymmetric hot nuclear matter. Using QCD sum rules [10] the in-medium masses of the lowest charmonium states can be written as

$$m^2 \simeq \frac{M^J_{n-1}(\xi)}{M^J_n(\xi)} - 4m_c^2\xi, \tag{30}$$

where M_n^J is the nth moment of the meson and ξ is the normalization scale. Using operator product expansion, the moment M_n^J can be written as [10]

$$M_n^J(\xi) = A_n^J(\xi)\left[1 + a_n^J(\xi)\alpha_s + b_n^J(\xi)\phi_b + c_n^J(\xi)\phi_c\right], \tag{31}$$

where $A_n^J(\xi)$, $a_n^J(\xi)$, $b_n^J(\xi)$, and $c_n^J(\xi)$ are the Wilson coefficients. The common factor A_n^J results from the bare loop diagram. The coefficients a_n^J take into account perturbative radiative corrections, while the coefficients b_n^J are associated with the scalar gluon condensate term

$$\phi_b = \frac{4\pi^2}{9}\frac{\left\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{a\mu\nu}\right\rangle}{(4m_c^2)^2}. \tag{32}$$

As already mentioned, the contribution of the scalar gluon condensate is taken through the dilaton field within the chiral SU(3) model used in the present investigation. Using Eq. (29), the preceding equation can be rewritten in terms of the dilaton field χ as

$$\phi_b = \frac{32\pi^2}{81(4m_c^2)^2}\left\{(1-d)\chi^4 + \left(\frac{\chi}{\chi_0}\right)^2 \times\left[m_\pi^2 f_\pi \sigma + \left(\sqrt{2} m_k^2 f_k - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi\right)\zeta\right]\right\}. \tag{33}$$

The coefficients A_n^J, a_n^J, and b_n^J are listed in Ref. [38]. The coefficients c_n^J are associated with the value of ϕ_c, which gives the contribution from twist-2 gluon operator and is given as

$$\phi_c = \frac{4\pi^2}{3(4m_c^2)^2} G_2, \tag{34}$$

where G_2 is given by Eq. (19). We calculate the in-medium masses of the charmonium states J/ψ and η_c in the hot asymmetric nuclear matter and compare the results with the contribution from the twist-2 gluon operator as calculated in the linear density approximation. In the low-density approximation, the term ϕ_c is given as [10]

$$\phi_c = -\frac{2\pi^2}{3}\frac{\left\langle \frac{\alpha_s}{\pi} A_G\right\rangle}{(4m_c^2)^2} m_N \rho_B. \tag{35}$$

In the preceding equation, A_G represents twice the momentum fraction carried by gluons in the nucleon and is set equal to 0.9 [10]. m_N and ρ_B are the nucleon mass and baryon density, respectively. The Wilson coefficients, c_n^J in the vector channel (for J/ψ) and the pseudoscalar channel (for η_c), can be found in Ref. [10]. The parameters m_c and α_s are the running charm quark mass and running coupling constant, respectively, and are ξ dependent [38]. These are given by

$$\frac{m_c(\xi)}{m_c} = 1 - \frac{\alpha_s}{\pi}\left[\frac{2+\xi}{1+\xi}\ln(2+\xi) - 2\ln 2\right], \qquad (36)$$

where $m_c \equiv m_c(p^2 = -m_c^2) = 1.26$ GeV [39], and

$$\alpha_s\left(Q_0^2 + 4m_c^2\right) = \alpha_s\left(4m_c^2\right) \Big/ \left(1 + \frac{25}{12\pi}\alpha_s\left(4m_c^2\right)\ln\frac{Q_0^2+4m_c^2}{4m_c^2}\right), \qquad (37)$$

with $\alpha_s(4m_c^2) \simeq 0.3$ and $Q_0^2 = 4m_c^2\xi$ [38].

In the next section, we present and discuss the results of our present investigation of the in-medium masses of J/Ψ and η_c in isospin asymmetric hot nuclear matter.

IV. RESULTS AND DISCUSSIONS

In this section, we first investigate the effects of density, isospin asymmetry, and temperature of the nuclear medium on the dilaton field χ in the chiral SU(3) model, from which we obtain the expectation value of the scalar gluon condensate in the medium. Using the QCD sum rule approach, the in-medium masses of charmonium states J/ψ and η_c are calculated from the medium dependence of the gluon condensates. The medium-dependent dilaton field χ is obtained by solving the equations of motion of the scalar fields, σ, ζ, δ, and χ, given by Eqs. (20)–(23). The values of the parameters used in the present investigation are $k_0 = 2.54$, $k_1 = 1.35$, $k_2 = -4.78$, $k_3 = -2.77$, $k_4 = -0.22$, and $d = 0.064$, which are the parameters occurring in the scalar meson interactions defined in Eq. (4). The vacuum values of the scalar isoscalar fields σ and ζ and the dilaton field χ are -93.3, -106.6, and 409.77 MeV, respectively. The values $g_{\sigma N} = 10.6$ and $g_{\zeta N} = -0.47$ are determined by fitting to the vacuum baryon masses. The other parameters fitted to the asymmetric nuclear-matter saturation properties in the mean-field approximation are $g_{\omega N} = 13.3$, $g_{\rho p} = 5.5$, $g_4 = 79.7$, $g_{\delta p} = 2.5$, $m_\zeta = 1024.5$ MeV, $m_\sigma = 466.5$ MeV, and $m_\delta = 899.5$ MeV. The nuclear-matter saturation density used in the present investigation is 0.15 fm^{-3}.

In Fig. 1, we show the variation of a dilaton field χ with temperature for both zero and finite baryon densities and for selected values of the isospin asymmetry parameter, $\eta = 0$, 0.1, 0.3, and 0.5 [18]. At zero baryon density, it is observed that the value of the dilaton field remains almost constant up to a temperature of about 130 MeV, above which it is seen to drop with increase in temperature. However, the drop in the dilaton field is seen to be very small up to a temperature of around 175 MeV, above which the drop is seen to be larger. The value of the dilaton field is seen to change from 409.8 MeV at $T = 0$ to about 409.7, 409.3, and 405.76 MeV at $T = 150$, 175, and 200 MeV, respectively. The thermal distribution functions have an effect of increasing the scalar densities at zero baryon density, that is, for $\mu_i^* = 0$, as can be seen from the expression of the scalar densities given by Eq. (24). This effect seems to be negligible up to a temperature of about 130 MeV. This leads to a decrease in the magnitudes of scalar fields σ and ζ. This behavior of the scalar fields is reflected in the value of χ, which is solved from the coupled equations of motion of the scalar fields, given by Eqs. (20)–(23), as a drop as we increase the temperature above a temperature of about 130 MeV. The scalar densities attaining nonzero values at high temperatures, even at zero baryon density, indicates the presence of baryon-antibaryon pairs in the thermal bath and has already been observed in the literature [25,40]. This leads to the baryon masses being different from their vacuum masses above this temperature, arising from modifications of the scalar fields σ and ζ.

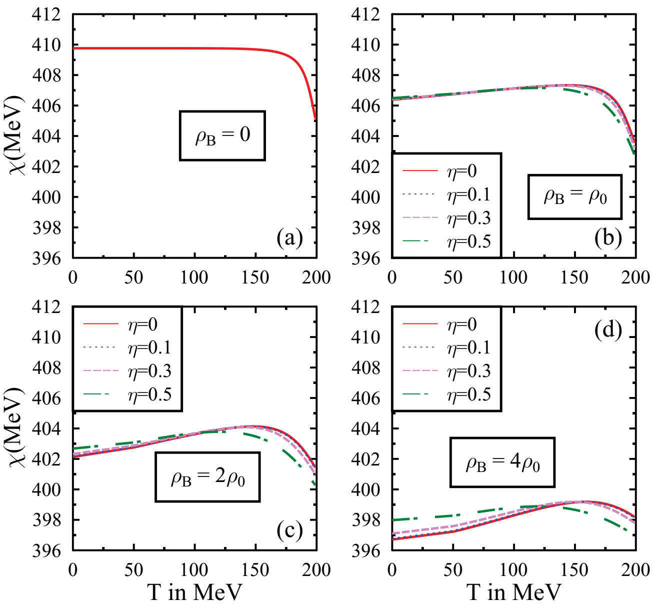

FIG. 1. (Color online) The dilaton field χ plotted as a function of temperature, at given baryon densities, for different values of the isospin asymmetry parameter η.

For finite-density situations, the behavior of the χ field with temperature is seen to be very different from the zero-density case, as can be seen in panels (b), (c), and (d) of Fig. 1, where the χ field is plotted as a function of temperature for densities ρ_0, $2\rho_0$, and $4\rho_0$, respectively. At finite densities, one observes first a rise and then a decrease of the dilaton field with temperature. This is related to the fact that at finite densities, the magnitude of the σ field (as well as of the ζ field) first show an increase and then a drop with further increase of the temperature [18], which is reflected in the behavior of χ field, because it is solved from the coupled equations of the scalar fields. The reason for the different behavior of the scalar fields (σ and ζ) at zero and finite densities can be understood in the following manner [25]. As has already been mentioned, the thermal distribution functions in Eq. (24) have an effect of increasing the scalar densities at zero baryon density, that is, for $\mu_i^* = 0$. However, at finite densities, that is, for nonzero values of the effective chemical potential, $\mu_i{}^*$, for increasing temperature, there are contributions also from higher momenta, thereby increasing the denominator of the integrand on the right-hand side of Eq. (24). This leads to a decrease in the scalar density. The competing effects of the thermal distribution functions and the contributions of the higher-momenta states give rise to the observed effect of the scalar density and hence of the σ and ζ fields with temperature at finite baryon densities [25]. This kind of behavior of the scalar σ field on temperature at finite densities has also been observed in the Walecka model by Li *et al.* [41], which was reflected as an increase in the mass of the nucleon with temperature at finite densities in the mean-field calculations. The effects of the behavior of the scalar fields on the value of the χ field, obtained from solving the coupled equations (20)–(23) for the scalar fields, are shown in Fig. 1.

In Fig. 1, it is observed that for a given value of isospin asymmetry parameter η, the dilaton field χ decreases with increase in the density of the nuclear medium. The drop in the value of χ with density is seen to be much larger than that seen with its modification with temperature at a given density. For isospin symmetric nuclear medium ($\eta = 0$) at temperature $T = 0$, the reduction in the dilaton field χ from its vacuum value ($\chi_0 = 409.8$ MeV) is seen to be about 3 MeV at $\rho_B = \rho_0$ and about 13 MeV at $\rho_B = 4\rho_0$. As we move from isospin symmetric medium with $\eta = 0$ to isospin asymmetric medium at temperature $T = 0$ for a given value of density, there is seen to be an increase in the value of the dilaton field χ. However, the effect of isospin asymmetry of the medium on the value of the dilaton field is observed to be negligible up to about a density of nuclear-matter saturation density, and is appreciable only at higher values of densities, as can be seen in Fig. 1. At nuclear-matter saturation density ρ_0, the value of dilaton field χ changes from 406.4 MeV in the symmetric nuclear medium ($\eta = 0$) to 406.5 MeV in the isospin asymmetric nuclear medium ($\eta = 0.5$). At a density of about $4\rho_0$, the values of the dilaton field are modified to 396.7 and 398 MeV at $\eta = 0$ and 0.5, respectively. Thus, the increase in the dilaton field χ with isospin asymmetry of the medium is seen to be greater at zero temperature as we move to higher densities.

At a finite density, ρ_B, and for given isospin asymmetry parameter η, the dilaton field χ is seen to first increase with temperature, and above a particular value of temperature, it is seen to decrease with further increase in temperature. At the nuclear-matter saturation density $\rho_B = \rho_0$ and in isospin symmetric nuclear medium ($\eta = 0$) the value of the dilaton field χ increases up to a temperature of about $T = 145$ MeV, above which there is a drop in the dilaton field. For $\rho_B = \rho_0$ in the asymmetric nuclear matter with $\eta = 0.5$, there is seen to be a rise in the value of χ up to a temperature of about 120 MeV, above which it starts decreasing. As has already been mentioned, at zero temperature and for a given value of density, the dilaton field χ is found to increase with increase in the isospin asymmetry of the nuclear medium. However, from Fig. 1, it is observed that at high temperatures and for a given density, the value of the dilaton field χ becomes higher in symmetric nuclear medium as compared to isospin asymmetric nuclear medium; for example, at nuclear saturation density $\rho_B = \rho_0$ and temperature $T = 150$ MeV the values of dilaton field χ are 407.3 and 407 MeV at $\eta = 0$ and 0.5, respectively. At density $\rho_B = 4\rho_0$, $T = 150$ MeV, the values of dilaton field χ are seen to be 399.1 and 398.7 MeV for $\eta = 0$ and 0.5, respectively. This observed behavior of the χ is related

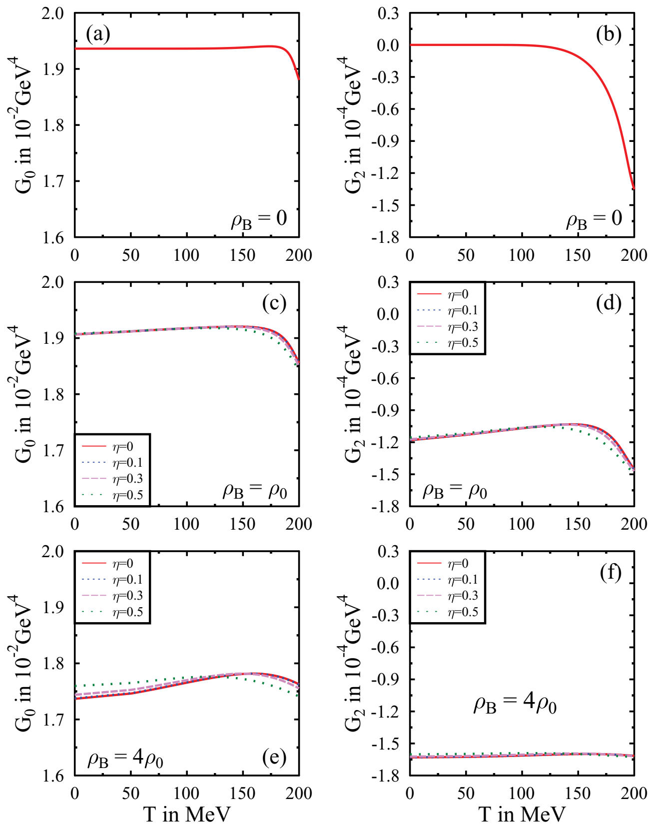

FIG. 2. (Color online) The functions G_0 and G_2 describing the trace and nontrace parts of the energy momentum tensor are plotted as functions of the density at different temperatures and for different values of the isospin asymmetry parameter η.

to the fact that at finite densities and for isospin asymmetric matter, there are contributions from the scalar isovector δ field whose magnitude is seen to decrease for higher temperatures for given densities, whereas δ field has zero contribution for isospin symmetric matter.

In Fig. 2, we show the variation of the trace and the nontrace parts of the energy momentum tensor given by Eq. (12) with temperature for different values of the baryon density and isospin asymmetry parameter, η. The trace part, $G_0 = \langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{a\mu\nu} \rangle$, is given by Eq. (29), and G_2, which is related to the nontrace part of the energy momentum tensor, is given by Eq. (19), both obtained from the SU(3) model used in the present investigation. The value of the trace part, G_0, is plotted as a function with temperature for densities $\rho_B = 0$, ρ_0, and $4\rho_0$ in panels (a), (c), and (e) in Fig. 2. For zero density, there is seen to be an increase of G_0 with temperature up to a temperature of about 175 MeV, and then a drop with further increase in the temperature. The values of G_0 are obtained as 1.9361×10^{-2}, 1.9362×10^{-2}, 1.9381×10^{-2}, and 1.88 $\times 10^{-2}$ GeV^4 at values of the temperature, $T = 0$, 100, 150, and 200 MeV, respectively. We might note here that the calculations of the scalar gluon condensate, G_0, given by Eq. (29), have been performed by accounting for the finite quark masses in the trace anomaly. In the absence of finite quark masses, the scalar gluon condensate becomes proportional to the fourth power of the dilaton field, as can be seen from Eq. (29). The dilaton field is observed to decrease with temperature at zero baryon density, as can be seen from Fig. 1. The values of G_0, for the limit of zero quark masses, also decrease accordingly with temperature for $\rho_B = 0$, with the values of G_0 given as 2.3455×10^{-2}, 2.34547×10^{-2}, 2.3437×10^{-2}, and 2.323×10^{-2} GeV^4 at values of temperature T of 0, 100, 150, and 200 MeV, respectively. A similar behavior of G_0 with temperature at zero baryon density has also been observed in Ref. [36]. In the present investigation, the finite quark mass term leads to a decrease in the value of G_0, as can be seen from Eq. (29). At finite densities, the dilaton field χ is seen to increase up to a temperature above which it starts decreasing, as can be seen from Fig. 1. Accounting for the finite quark masses, we get a positive contribution to G_0 from the temperature effects from the second term in Eq. (29), leading to an increase in the scalar condensate up to a temperature above which there is seen to be a decrease with further rise in temperature. For baryon densities of $\rho_B = \rho_0$ and $4\rho_0$, the values up to which G_0 increases with temperature are about 145 and 175 MeV, respectively. In isospin symmetric nuclear matter, for $\rho_B = \rho_0$, the values of G_0 are observed to be 1.90646×10^{-2}, 1.91755×10^{-2}, 1.92×10^{-2}, and 1.8554×10^{-2} GeV^4 for temperatures of 0, 100, 150, and 200 MeV, respectively. For the same values of the temperature, in the absence of finite quark masses, the values of G_0 are observed to be 2.269×10^{-2}, 2.2857×10^{-2}, 2.29×10^{-2}, and 2.2×10^{-2} GeV^4 for $\rho_B = \rho_0$ and $\eta = 0$. In isospin symmetric nuclear matter for $\rho_B = 4\rho_0$ the values of G_0 are given as 1.7367×10^{-2} GeV^4 (2.06×10^{-2} GeV^4), 1.7656×10^{-2} GeV^4 (2.094×10^{-2} GeV^4), 1.78×10^{-2} GeV^4 (2.112×10^{-2} GeV^4), and 1.7626×10^{-2} GeV^4 (2.09×10^{-2} GeV^4) for values of temperature, $T = 0$, 100, 150, and 200 MeV, respectively, for the cases of the finite (zero) quark masses in the trace anomaly.

The nontrace part of the energy momentum tensor, G_2, is plotted as a function of temperature in subplots (b), (d), and (f) of Fig. 2 for densities, $\rho_B = 0$, ρ_0, and $4\rho_0$. It may be noted that value of G_2 is zero in vacuum, and this has a nonzero contribution only for finite density and/or temperature. The magnitude of the quantity G_2 is observed to increase with increase in the temperature of the nuclear medium for zero density, with the values of G_2 at $\rho_B = 0$ given as -7.106×10^{-12}, -2.386×10^{-7}, -1.106×10^{-5}, and -3.528×10^{-5} GeV^4 at values of temperature, T as 50, 100, 150, and 200 MeV, respectively. The observed behavior of the magnitude of G_2 increasing as a function of temperature at zero baryon density has also been observed in Ref. [36]. At nuclear saturation density $\rho_B = \rho_0$ there is seen to be a decrease in the magnitude of G_2 with temperature and then an increase with further rise in temperature. In isospin symmetric medium, for $\rho_B = \rho_0$, the values of G_2 are given as -1.181×10^{-4}, -1.130×10^{-4}, -1.069×10^{-4}, -1.034×10^{-4}, and -1.4527×10^{-4} GeV^4 at values of temperature $T = 0$, 50, 100, 150, and 200 MeV, respectively. For density $4\rho_0$ and $\eta = 0$, the values of G_2 are given as -1.63×10^{-4}, -1.626×10^{-4}, -1.613×10^{-4}, -1.5992×10^{-4}, and -1.6156×10^{-4} GeV^4 for $T = 0$, 50, 100, 150, and 200 MeV, respectively. In isospin asymmetric medium, $\eta = 0.5$, at $\rho_B = 4\rho_0$, the values of G_2 are -1.602×10^{-4}, -1.598×10^{-4}, -1.591×10^{-4}, -1.5991×10^{-4}, and -1.6265×10^{-4} GeV^4 at temperature $T = 0$, 50, 100, 150, and 200 MeV, respectively. In the present investigation, the effects of isospin asymmetry and temperature of the nuclear medium on the values of G_0 and G_2 are observed to be small and the effect of density seems to be the dominant effect. This is related to the fact that the dilaton field and the scalar fields, σ, ζ, and δ in the hot isospin asymmetric nuclear medium are strongly dependent on the density of the medium and the effects of temperature and isospin asymmetry on these scalar fields are much smaller than the density effects.

TABLE I. The mass shifts of J/ψ and η_c are shown at densities of ρ_0, $2\rho_0$, and $4\rho_0$ at values of the isospin asymmetric parameter $\eta = 0$ and 0.5 for $\xi = 0.874$. This value of ξ reproduces the vacuum mass of J/ψ as 3097 MeV.

ρ_B	J/ψ		η_c	
	$\eta = 0$	$\eta = 0.5$	$\eta = 0$	$\eta = 0.5$
ρ_0	−4.48	−4.34	−5.21	−5.06
$2\rho_0$	−10	−9.29	−9.14	−8.64
$4\rho_0$	−16.77	−15.19	−13.12	−12.19

After obtaining the medium modification of the scalar gluon condensate from the value of the dilaton field using Eq. (29) and of the twist-2 gluon operator by using Eqs. (11) and (19), we next determine the in-medium mass shifts of J/ψ and η_c mesons using the QCD sum rule approach. We use the moments in the range $5 \leqslant n \leqslant 12$ and fix the value of parameter $\xi = 0.874$ so that we can reproduce the vacuum value of mass of J/ψ, $m_{J/\psi} = 3097$ MeV. For this value of ξ, the parameter $\alpha_s = 0.2667$ and the running charm quark mass, $m_c = 1.232 \times 10^3$ MeV. We consider the contributions from scalar gluon condensate $\langle \frac{\alpha_s}{\pi} G^a_{\sigma\kappa} G^{a\sigma\kappa} \rangle$ and the twist-2 tensorial gluon operator $\langle \frac{\alpha_s}{\pi} G^a_{\mu\sigma} G^{a\sigma}_{\nu} \rangle$ through the dilaton field χ within the chiral SU(3) model used in the present investigation. We obtain the value of ϕ_b, which is related to the scalar gluon condensate by Eq. (32), from the medium-dependent χ field using Eq. (33). The value for ϕ_c arising from the twist-2 tensorial gluon operator is calculated within the chiral SU(3) model by using Eq. (34). We also compare our results with the twist-2 gluon operator as calculated from the formula obtained in the low-density approximation as given by Eq. (35) [10]. The value of ϕ_c at nuclear saturation density $\rho_0 = 0.15$ fm^{-3} is calculated to be -4.2158×10^{-5} within the SU(3) chiral model used in the present investigation, which may be compared to the value of -1.4685×10^{-5} in the linear density approximation [10]. In Table I, we summarize the results for the mass shifts of J/ψ and η_c, as obtained in the present investigation, at zero temperature, for values of the baryon densities of ρ_0, $2\rho_0$, and $4\rho_0$ and for isospin-asymmetry parameters of $\eta = 0$ and 0.5 for the value of $\xi = 0.874$. As has already been mentioned, this value of ξ is fixed so as to obtain the observed vacuum mass of J/ψ as 3097 MeV.

In Fig. 3, we show the variation of masses of J/ψ and η_c mesons with n for a fixed value of baryon density, $\rho_B = \rho_0$, and for different values of isospin asymmetry parameter η. We show the results for values of temperature $T = 0$, 100,

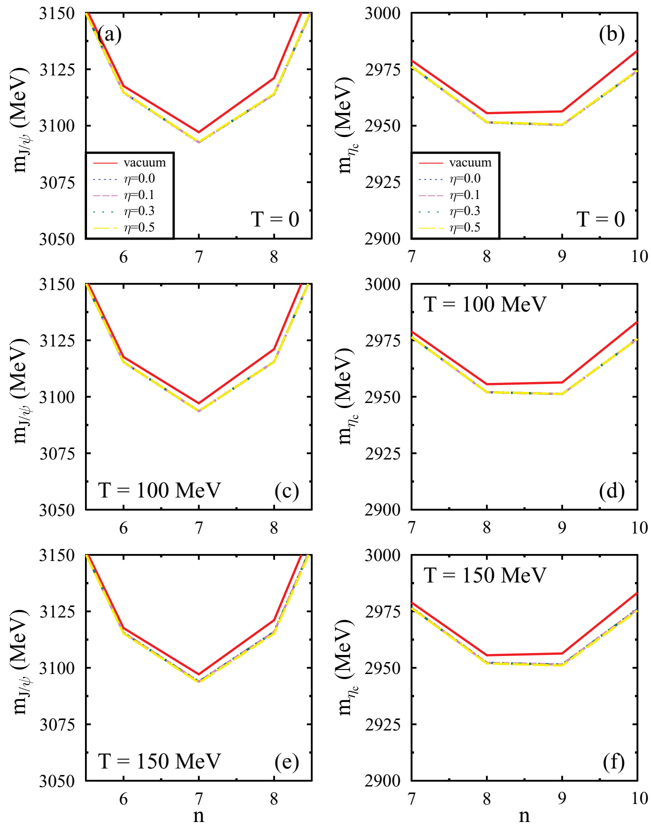

FIG. 3. (Color online) The in-medium masses of the J/ψ and η_c mesons plotted as functions of n for nuclear-matter saturation density ρ_0 at different temperatures and for different values of the isospin asymmetry parameter η. The value of parameter ξ is taken as 0.874, which reproduces the vacuum mass of J/ψ as 3097 MeV.

and 150 MeV. In symmetric nuclear matter, at nuclear-matter saturation density, $\rho_B = \rho_0$, and at temperature $T = 0$, we obtain the mass shifts for J/ψ and η_c mesons to be equal to -4.48 and -5.21 MeV, respectively, as can be seen from Table I. These values of mass shifts for for J/ψ and η_c mesons may be compared with the mass shifts of -7 and -5 MeV, respectively obtained in the linear density approximation in Ref. [10]. In the present investigation, we calculate the values of ϕ_b and ϕ_c from the medium modification of the dilaton field, χ, within the chiral SU(3) model by using Eqs. (33) and (34). In isospin symmetric nuclear medium, at baryon densities, $\rho_B = 0$ and ρ_0, the values of the dilaton field, χ are 409.76 and 406.38 MeV, respectively, and, hence, using Eq. (29), the values of the scalar gluon condensate $\langle \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{a\mu\nu} \rangle$ turn out to be (373 MeV)^4 and (371.6 MeV)^4 for densities $\rho_B = 0$ and ρ_0, respectively. We note here that when we neglect the quark masses in the trace anomaly, the values of the scalar gluon condensate at these densities are modified to (391 MeV)^4 and (388 MeV)^4, respectively. We thus observe an increase of the values of the scalar gluon condensate by about 20% when we do not account for the finite masses of the quarks. The values of ϕ_b, accounting for the finite quark masses, turn out to be 2.3×10^{-3} and 2.27×10^{-3} in the vacuum and at nuclear-matter saturation density, ρ_0, respectively. These may be compared with the values of ϕ_b to be equal to 1.7×10^{-3} and 1.6×10^{-3}, respectively, for $\rho_B = 0$ and for $\rho_B = \rho_0$, obtained from the values of scalar gluon condensate of (350 MeV)^4 and (344.81 MeV)^4, respectively, in vacuum and at nuclear saturation density ρ_0 in Ref. [10] in the linear density approximation. We might note here that the value of nuclear-matter saturation density used in the present calculations is 0.15 fm^{-3} and in Ref. [10] it was taken to be 0.17 fm^{-3}. When the quark masses are neglected and ϕ_c is as calculated in the chiral SU(3) model used in the present investigation, the values of the mass shifts for J/ψ and η_c turn out to be -8.01 and -5.13 MeV, respectively. When we calculate ϕ_b within the chiral SU(3) model, but calculate the contribution of the twist-2 operator through ϕ_c calculated in the linear density approximation given by Eq. (35) [10], we obtain the mass shifts for J/ψ and η_c at $\rho_B = \rho_0$ for symmetric nuclear matter at zero temperature to be given as -2.88 and -2.02 MeV, respectively. The value of the mass shift of J/ψ of about -4.48 MeV at the nuclear matter density ρ_0 in symmetric nuclear matter at zero temperature obtained in the present investigation may be compared to the value of the mass shift of -8 MeV obtained using QCD second-order Stark effect using the value of the scalar gluon condensate obtained using a linear density approximation [17], as well as a value of -8.6 MeV, when the scalar gluon condensate was obtained from the expectation value of the scalar dilaton field in a chiral SU(3) model [18]. We observe in Fig. 3 that the isospin dependence of the mass shifts of J/ψ and η_c are very small. This is attributable to the fact that the dependence of χ on the isospin asymmetry is very small, as can be seen from Fig. 1.

Figures 4 and 5 show the mass shifts of J/ψ and η_c for baryon densities $\rho_B = 2\rho_0$ and $4\rho_0$, respectively, at different temperatures and different values of the isospin asymmetry parameter η. In isospin symmetric nuclear medium, at density $\rho_B = 2\rho_0$ and temperature $T = 0$, the mass shifts for J/ψ and η_c mesons are observed to be -10 and -9.14 MeV, respectively. The effects of isospin asymmetry of the medium on the mass shift of the J/ψ and η_c mesons are seen to be almost negligible, as can be seen from Table I. This is attributable to the very small changes in the dilaton field with the isospin asymmetry of the medium, as can be seen from Fig. 1. In isospin asymmetric nuclear medium ($\eta = 0.5$), at nuclear saturation density ρ_0, the mass shifts in J/ψ and η_c mesons at zero temperature are observed to be -4.34 and -5.06 MeV from their vacuum values, which may be compared with the values of -4.48 and -5.21 MeV, respectively, for the isospin symmetric nuclear matter. At values of the baryon density ρ_B of $2\rho_0$ and $4\rho_0$, as can be seen from Table I, the isospin dependence of the mass shifts for J/ψ and η_c mesons are seen to be negligibly small.

The effects of temperature on the dilaton field χ is very small and this is reflected in the small change in the mass shifts of J/ψ and η_c mesons with temperature T. In isospin symmetric nuclear medium ($\eta = 0$), at nuclear saturation density $\rho_B = \rho_0$, the mass shifts for J/ψ meson from their vacuum values are observed to be -4.01, -3.5, and -3.23 MeV at temperatures $T = 50$, 100, and 150 MeV, respectively. At baryon density $\rho_B = 4\rho_0$, the values of mass shift for J/ψ meson change to -16.13, -14.82, and

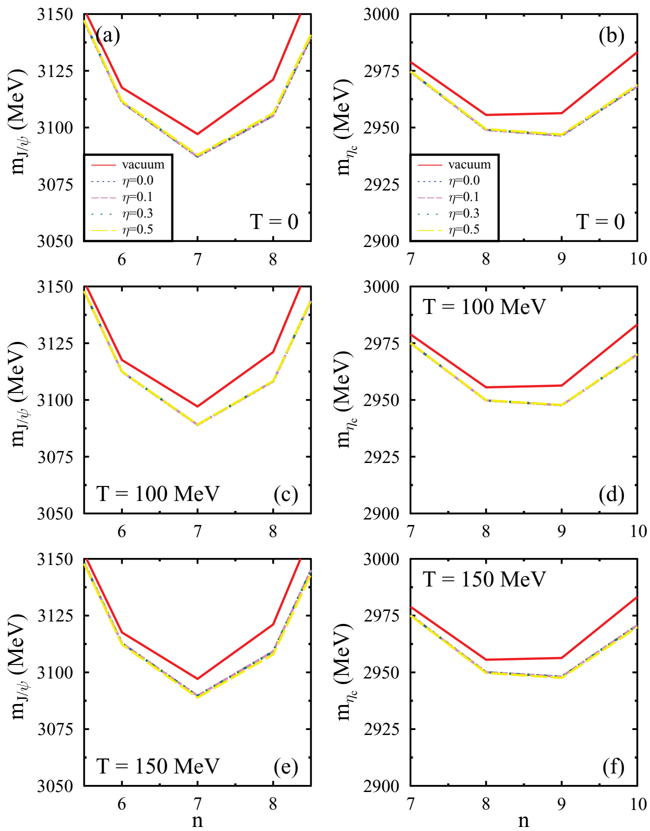

FIG. 4. (Color online) The in-medium masses of the J/ψ and η_c mesons plotted as functions of n for baryon density $\rho_B = 2\rho_0$ at different temperatures and for different values of the isospin asymmetry parameter η. The value of parameter ξ is taken as 0.874, which reproduces the vacuum mass of J/ψ as 3097 MeV.

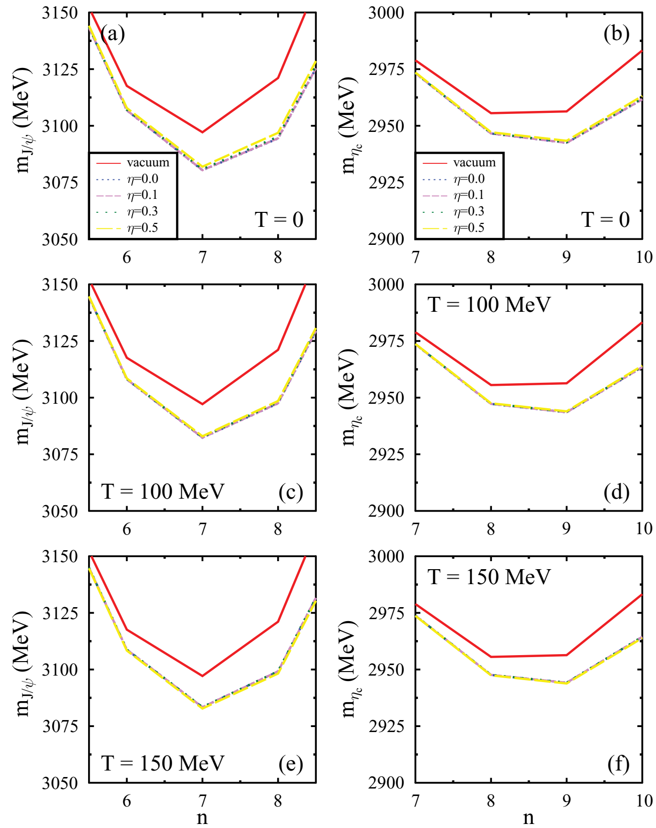

FIG. 5. (Color online) The in-medium masses of the J/ψ and η_c mesons plotted as functions of n for baryon density $\rho_B = 4\rho_0$ at different temperatures and for different values of the isospin asymmetry parameter η. The value of parameter ξ is taken as 0.874, which reproduces the vacuum mass of J/ψ as 3097 MeV.

−13.76 MeV at temperatures $T = 50$, 100, and 150 MeV, respectively. The value of the mass shift obtained at finite value of temperature is observed to be smaller than that seen in the zero-temperature case. This is because, at finite value of baryon density ρ_B, the dilaton field χ increases with increase in the temperature of the nuclear medium, but the increase is very small. In isospin asymmetric nuclear medium $\eta = 0.5$, at density $\rho_B = 4\rho_0$, the mass shifts for J/ψ mesons from their vacuum values are −14.82, −14.16, and −14.36 MeV at temperatures $T = 50$, 100, and 150 MeV, respectively.

For the pseudoscalar meson η_c, the mass shifts at nuclear saturation density ρ_0 in nuclear medium with $\eta = 0(0.5)$ are −4.81(−4.73), −4.352(−4.345), and −4.1(−4.54) MeV at temperatures $T = 50$, 100, and 150 MeV, respectively. At density $\rho_B = 4\rho_0$, with $\eta = 0$ (0.5), these values are modified to −12.77(−11.98), −12.02(−11.6), and −11.41 (−11.73) MeV, respectively for $T = 50$, 100, and 150 MeV. It may be noted that at high values of temperatures, for example, at $T = 150$ MeV, the mass shift is greater in the isospin asymmetric nuclear medium ($\eta = 0.5$) than in the isospin symmetric nuclear medium ($\eta = 0$). This is opposite to the zero-temperature case. The reason is that at high temperatures the dilaton field χ has larger drop in the isospin asymmetric nuclear medium ($\eta = 0.5$) than in the isospin symmetric nuclear medium ($\eta = 0$), as can be seen in Fig. 1.

As mentioned earlier, for the preceding calculations we had fixed the value of parameter ξ so as to reproduce the vacuum value of J/ψ mass. However, with this value of ξ, the vacuum value of η_c meson comes out to be 2955.6 MeV. We can reproduce the vacuum value of pseudoscalar meson η_c = 2980.5 MeV if we fix the value of $\xi = 0.8995$. For this value of ξ, the parameter $\alpha_s = 0.266$ and the running charm quark mass $m_c = 1.2313 \times 10^3$ MeV. For these values of parameters, the mass shifts for J/ψ and η_c mesons, in nuclear medium at zero temperature, at densities ρ_0, $2\rho_0$, and $4\rho_0$ for $\eta = 0(0.5)$ are summarized in Table II.

TABLE II. The mass shifts of J/ψ and η_c are shown at densities of ρ_0, $2\rho_0$, and $4\rho_0$ at values of the isospin asymmetric parameter $\eta = 0$ and 0.5 for $\xi = 0.8995$. This value of ξ reproduces the vacuum mass of η_c as 2980.5 MeV.

ρ_B	J/ψ		η_c	
	$\eta = 0$	$\eta = 0.5$	$\eta = 0$	$\eta = 0.5$
ρ_0	−4.27	−4.14	−5.69	−5.54
$2\rho_0$	−9.55	−8.87	−9.39	−8.92
$4\rho_0$	−16.02	−14.51	−13.12	−12.25

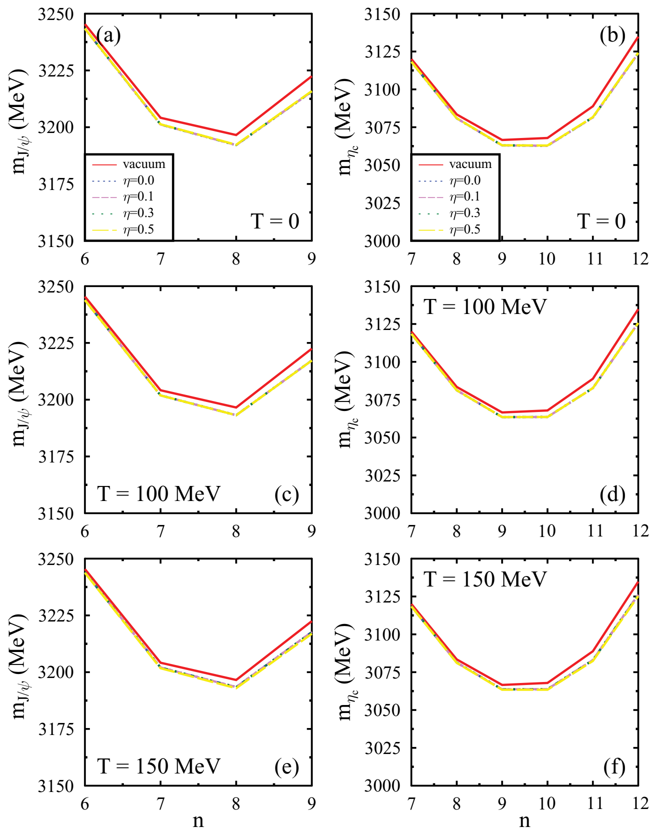

FIG. 6. (Color online) The in-medium masses of the J/ψ and η_c mesons plotted as functions of n for nuclear matter saturation density ρ_0 at different temperatures and for different values of the isospin asymmetry parameter η, with $\xi = 1$.

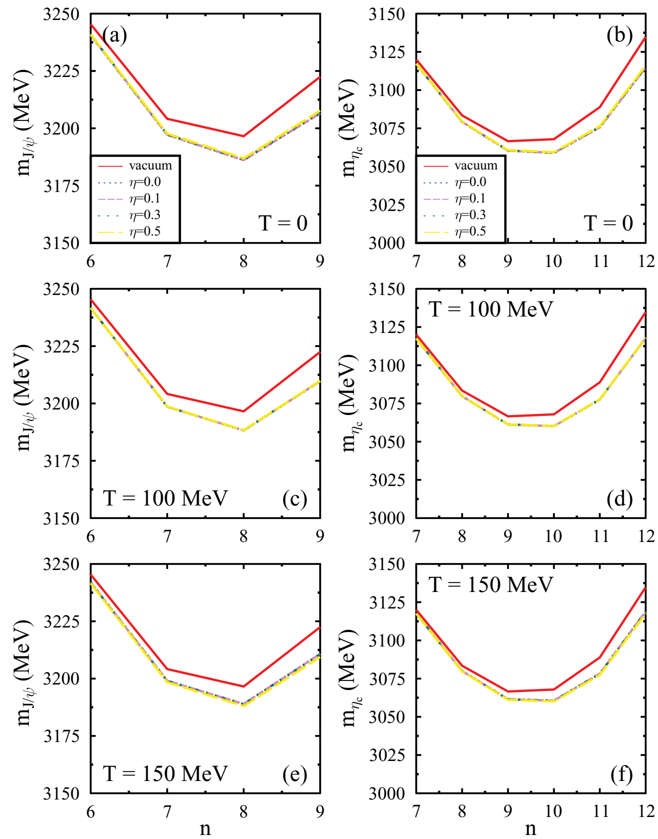

FIG. 7. (Color online) The in-medium masses of the J/ψ and η_c mesons plotted as functions of n for baryon density of $2\rho_0$ at different temperatures and for different values of the isospin asymmetry parameter η, with $\xi = 1$.

We also show the results for the mass modifications of J/ψ and η_c mesons, if we consider the value of parameter, $\xi = 1$ [10], leading to the value of α_s as 0.21 and of m_c as 1.24×10^3 MeV. Figures 6, 7, and 8 show the temparature and isospin-asymmetry dependence of the mass modifications of J/ψ and η_c mesons for baryon densities of ρ_0, $2\rho_0$, and $4\rho_0$, respectively, with parameter $\xi = 1$. We observe that with $\xi = 1$, the vacuum values of the masses of J/ψ and η_c mesons are given as 3196.56 and 3066.57 MeV, respectively. With $\xi = 1$, the results for the mass shifts for J/ψ and η_c at different densities with $\eta = 0$ and 0.5, and zero temperature, obtained in the present investigation are summarized in Table III. The values of the mass shifts for J/ψ meson in isospin symmetric medium, with $\xi = 1$, at nuclear saturation density $\rho_B = \rho_0$ are observed to be -3.92, -3.38, and -3.1 MeV for $T = 50$, 100, and 150 MeV, respectively. At baryon density $\rho_B = 4\rho_0$, these values of the mass shift change to -17.23, -15.77, and -14.59 MeV at temperature $T = 50$, 100, and 150 MeV, respectively. For pseudoscalar meson η_c, the mass shifts at $\rho_B = \rho_0$ are obtained to be -3.42, -3.06, and -2.91 MeV for $T = 50$, 100, and 150 MeV, respectively, whereas at $\rho_B = 4\rho_0$ these values of mass shift are seen to be modified to -11.47, -10.68, and -10.04 MeV, respectively.

In Ref. [15] the operator product expansion was carried out up to dimension six and the mass shift for J/ψ was found to be -4 MeV at nuclear saturation density ρ_0 and temperature $T = 0$. The effect of temperature on the J/ψ in deconfinement phase was studied in Refs. [42,43]. In these investigations, it was reported that J/ψ mass is essentially constant in a wide range of temperatures and above a particular value of the temperature, T, there is a sharp change in the mass of J/ψ in the deconfined phase; for example, in Ref. [44] the mass shift for J/ψ was reported to be about 200 MeV at $T = 1.05T_c$. The pseudoscalar charmonium spectral function for different temperatures was studied using a screened potential in Ref. [45]. The effect of rising temperature was observed to melt the higher excited states by $1.1T_C$ and to shift the continuum threshold to lower energies. In these studies, it

TABLE III. The mass shifts of J/ψ and η_c are shown at densities of ρ_0, $2\rho_0$, and $4\rho_0$ at values of the isospin asymmetric parameter $\eta = 0$ and 0.5 for $\xi = 1$.

ρ_B	J/ψ		η_c	
	$\eta = 0$	$\eta = 0.5$	$\eta = 0$	$\eta = 0.5$
ρ_0	-4.43	-4.28	-3.8	-3.66
$2\rho_0$	-10.43	-9.66	-7.67	-7.18
$4\rho_0$	-17.93	-16.19	-11.85	-10.87

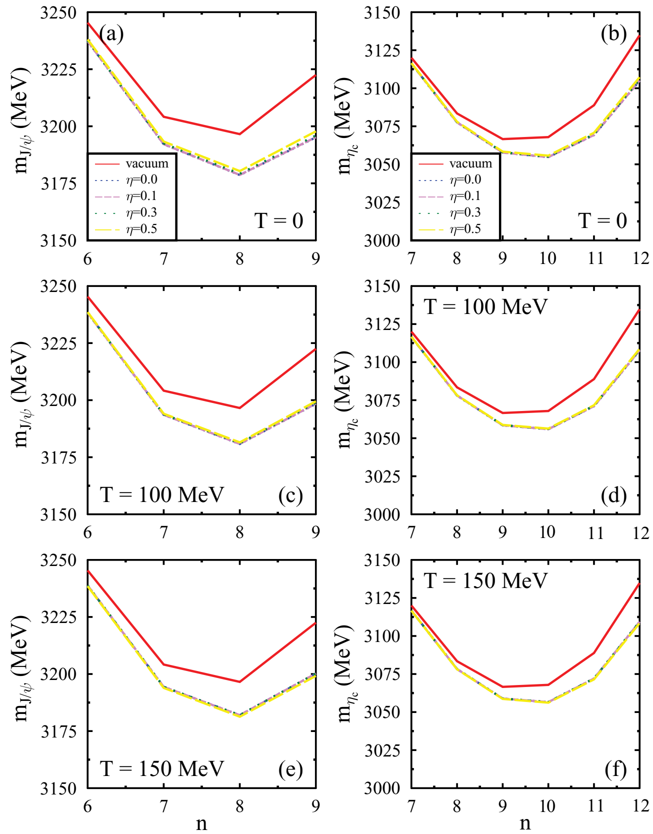

FIG. 8. (Color online) The in-medium masses of the J/ψ and η_c mesons plotted as functions of n for baryon density of $4\rho_0$ at different temperatures and for different values of the isospin asymmetry parameter η, with $\xi = 1$.

was observed that the charmonium η_c survives even in the deconfined phase. In Refs. [45,46], the effect of temperature on η_c in the deconfinement phase was studied. In Ref. [46], it was reported that the J/ψ and η_c survive as distinct resonances in the plasma even up to $T \simeq 1.6T_c$ and that they eventually dissociate between $1.6T_c$ and $1.9T_c$. This suggests that the deconfined plasma is nonperturbative enough to hold heavy-quark bound states. In the present work, we have studied the effects of temperature on the mass modifications of J/ψ and η_c mesons in the confined phase dueowing to modifications of the scalar gluon condensate and twist-2 tnesorial gluon operator, simulated by a medium-dependent scalar dilaton field in chiral SU(3) model and the temperature effects are found to be very small as compared to the density effects.

V. SUMMARY

In summary, in the present investigation, we have studied the mass modifications of the charmonium states J/ψ and η_c in the nuclear medium using the QCD sum rule approach and using modification of a dilaton field (which simulates the gluon condensates) within a chiral SU(3) model. The in-medium modifications of the J/ψ and η_c are studied as arising owing to changes in the scalar and twist-2 gluon condensates in the nuclear medium, obtained from the medium modificaion of the χ field. The value of the dilaton field in the hot nuclear matter is obtained by solving the coupled Eqs. (20)–(23), which are the equations of motion of the σ, ζ, δ, and χ fields. The dilaton field χ thus depends on the scalar isovector field δ, which is related to the isospin asymmetry of the nuclear medium. The isospin asymmetry dependence of the χ, in turn, leads to the isospin asymmetry dependence of the charmonium states J/ψ and η_c. The modification of the χ field is observed to be small with the isospin asymmetry of the medium, as can be seen from Fig. 1. This is related to the fact that the magnitude of the obtained value of δ after solving the coupled equations for the scalar fields turns out to be much smaller (about few percent) as compared to the magnitudes of σ and ζ and hence the isospin asymmetry (through δ) only gives rise to a very small modification of the dilaton field χ [18]. It is observed that the temperature effect on the χ field is also very small and the modification of the dilaton field with density is seen to be the dominant medium effect in the present investigation. The negligible dependence of the dilaton field on isospin asymmetry as well as on temperature is reflected in the small isospin asymmetry and temperature dependence of the masses of the J/ψ and η_c states in the nuclear medium. Experimentally, measurements of dileptons (diphotons) in heavy-ion collisions may provide a clue to the properties of vector (pseudoscalar) mesons in hot and dense matter [46]. The present study of the in-medium properties of J/ψ and η_c mesons will be helpful for the experiments in the future facility of the FAIR, GSI, where the compressed baryonic matter at high densities and moderate temperature will be produced.

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