

In-medium vector meson properties and low-mass dilepton production from hot hadronic matter

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The in-medium properties of the vector mesons are known to be modified significantly in hot and dense hadronic matter due to vacuum polarisation effects from the baryon sector in the Walecka model. The vector meson mass drops significantly in the medium due to the effects of the Dirac sea. In the variational approach adopted in the present paper, these effects are taken into account through a realignment of the ground state with baryon condensates. Such a realignment of the ground state becomes equivalent to summing of the baryonic tadpole diagrams in the relativistic Hartree approximation (RHA). The approximation scheme adopted here goes beyond RHA to include quantum effects from the scalar meson and is nonperturbative and self-consistent. It includes multiloop effects, thus corresponding to a different approximation as compared to the one loop approximation of including scalar field quantum corrections. In the present work, we study the properties of the vector mesons in the hot and dense matter as modified due to such quantum correction effects from the baryon as well as scalar meson sectors. These medium modifications of the properties of the vector mesons are reflected, through the shifting and broadening of the respective peaks, in the low mass dilepton spectra. There is broadening of the peaks due to corrections from scalar meson quantum effects as compared to the relativistic Hartree approximation. It is seen to be rather prominent for the ω meson in the invariant mass plot.

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I. INTRODUCTION

The in-medium properties of the vector mesons (ρ and ω) in hot and dense matter have an important role to play in the low mass dilepton production resulting from heavy ion collision experiments. This has, hence, been a topic of great interest in the recent past, both experimentally [1–5] and theoretically [6–11]. The experimental observation of enhanced dilepton production [1–3] in the low invariant mass regime could be due to a reduction in the vector meson masses in the medium. Brown and Rho suggested the hypothesis that the vector meson masses drop in the medium according to a simple (BR) scaling law [6], given as $m_V^*/m_V = f_\pi^*/f_\pi$, where f_π is the pion decay constant and asterisk refers to in-medium quantities. In the framework of Quantum Hadrodynamics (QHD) as a description of the hot hadronic matter, it is seen that the dropping of the vector meson masses has its dominant contribution arising from the vacuum polarisation effects from the baryon sector [12–15], which is not observed in the mean field approximation. This means that the quantum effects do play an important role in the medium modification of vector meson properties. There have also been approaches based on QCD sum rules [7,8] which confirm a scaling hypothesis [7] as suggested by Brown and Rho, with a saturation scheme that leads to a delta function at the vector meson pole and a continuum for higher energies for the hadronic spectral function. It is, however, seen that such a simple saturation scheme for finite densities does not work and a more realistic description for the hadronic spectral function is called for [9]. Using an effective Lagrangian model to calculate the hadronic spectral function, it is seen that such a universal scaling law is not suggested for in-medium vector meson masses [10].

The medium modifications of the vector mesons have been thus a subject of several investigations. It has been emphasized in the literature [15,16] that the properties of the hadrons are modified due to their interactions with the thermal bath, and such modifications are reflected in the dilepton and photon spectra emitted from a hot and dense matter [17]. Dileptons are interesting probes for the study of evolving matter arising from relativistic heavy ion collisions since they do not interact strongly and escape unthermalized from the hot and dense matter at all stages of the evolution. The temperature [18] and density [19] modifications of the dileptons from hot hadronic matter as well as from a Quark Gluon Plasma (QGP) resulting from heavy ion collisions have been a subject of extensive investigations in the recent past. Broadly, two types of modifications for the hadrons in the medium are expected:

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shift in the pole position for the hadron propagator, giving rise to a modification of the mass, and broadening of the spectral function. In the present work, we shall attempt to study the medium modification of vector meson masses and decay widths in the hot hadronic matter within the framework of Quantum Hadrodynamics (QHD) taking the vacuum polarisation effects into account, and their subsequent effect on the dilepton spectra.

In the conventional hadronic models [20,21], the masses of the vector mesons stay constant or increase slightly, in the mean field approximation, i.e., when the polarization from the Dirac sea is neglected [21]. With the inclusion of quantum corrections from the baryonic sector, however, in the Walecka model one observes a drop in the vector meson masses in the medium [12–14]. This medium modification of the vector meson masses plays an important role in the enhanced dilepton yield [16] for masses below the ρ resonance in the heavy ion collision experiments [1,2]. It has been emphasized recently that the Dirac sea contribution (taken into account through summing over baryonic tadpole diagrams in the relativistic Hartree approximation (RHA)) dominates over the Fermi sea contribution and $m_\omega^*/m_\omega < 1$ is caused by a large dressing of $\bar{N}N$ cloud in the medium [14]. Further quantum effects arising from the scalar mesons have been considered [20] along with the RHA for the baryon sector in the context of vector meson mass modifications in strange hadronic matter [22]. The present method is a step in that direction of studying medium modifications of the vector meson properties including such quantum effects from the scalar mesons in a self-consistent manner, along with those arising from the baryon sector.

It was earlier demonstrated in a nonperturbative formalism that a realignment of the ground state with baryon-antibaryon condensates is equivalent to the relativistic Hartree approximation [23]. The ground state for the nuclear matter was extended to include sigma condensates to take into account the quantum correction effects from the scalar meson sector [23]. Such a formalism includes multiloop effects and is self consistent [23–25]. The methodology was then generalized to consider hot nuclear matter [26] as well as to the situation of hyperon-rich dense matter [27] relevant for neutron stars. The effect of vacuum polarisations on the vector meson masses has also been recently studied [28]. In the present work, we study the effect of such quantum corrections on the in-medium vector meson properties. This apart, for studying ρ meson properties, we take into account the ρ - π - π interactions, and their effects on the dilepton production in the low invariant mass regime. The pionic interactions are seen to modify the ρ meson mass only marginally, as already stated in the literature [15,29]. Due to scalar meson contributions, there is broadening in the ω and ρ peaks, as compared to RHA. This is seen to be quite pronounced for the ω meson, which leads to smearing and ultimate disappearance of the ω peak at finite densities [30].

We organize the paper as follows. We first briefly recapitulate the nonperturbative framework used for studying the quantum correction effects in hot nuclear matter in section II. The medium modification of the ω and ρ meson masses and decay widths are considered in section III. We discuss the effect of these medium modifications on the low mass dilepton spectra in section IV. In section V, we summarize the results of the present work and give an outlook.

II. QUANTUM VACUUM IN HOT NUCLEAR MATTER

We briefly recapitulate here the vacuum polarisation effects arising from the nucleon and scalar meson fields in hot nuclear matter in a nonperturbative variational framework [26]. The method of thermofield dynamics (TFD) [31] is used here to study the “ground state” (the state with minimum thermodynamic potential) at finite temperature and density within the Walecka model with a quartic scalar self interaction. The temperature and density dependent baryon and sigma masses are also calculated in a self-consistent manner in the formalism. The ansatz functions involved in such an approach are determined through functional minimisation of the thermodynamic potential.

The Lagrangian density in the Walecka model is given as

$$\begin{aligned} \mathcal{L} = & \bar{\psi} (i\gamma^\mu \partial_\mu - M - g_\sigma \sigma - g_\omega \gamma^\mu \omega_\mu) \psi + \frac{1}{2} \partial^\mu \sigma \partial_\mu \sigma - \frac{1}{2} m_\sigma^2 \sigma^2 - \lambda \sigma^4 \\ & + \frac{1}{2} m_\omega^2 \omega^\mu \omega_\mu - \frac{1}{4} (\partial_\mu \omega_\nu - \partial_\nu \omega_\mu) (\partial^\mu \omega^\nu - \partial^\nu \omega^\mu). \end{aligned} \quad (1)$$

In the above, ψ , σ , and ω_μ are the fields for the nucleon, σ , and ω mesons with masses M , m_σ , and m_ω respectively. Since we are interested in symmetric nuclear matter, the isovector rho meson does not contribute. The quartic coupling term in σ is necessary for the sigma condensates to exist, through a vacuum realignment [23]. Our calculations thus include the quantum effects arising from the sigma meson in addition to the mean field contribution from the quartic self interaction of the scalar meson. We retain the quantum nature of both the nucleon and the scalar meson fields, whereas the vector ω meson is treated as a classical field, using the mean field approximation for the ω meson, given as $\langle \omega^\mu \rangle = \delta_{\mu 0} \omega_0$. The reason is that without higher-order term for the ω meson, the quantum effects generated due to the ω meson through the present variational ansatz turn out to be zero.

We then write down the expressions for the thermodynamic quantities including the quantum effects. The details regarding the formalism can be found in earlier references [23,26,28]. The energy density, after carrying out the renormalisation procedures for the baryonic and scalar meson sectors [26], is obtained as

$$\epsilon_{\text{ren}} = \epsilon_{\text{finite}}^{(N)} + \Delta\epsilon_{\text{ren}} + \epsilon_{\omega} + \Delta\epsilon_{\sigma}, \quad (2)$$

with

$$\epsilon_{\text{finite}}^{(N)} = \gamma(2\pi)^{-3} \int d\mathbf{k} (k^2 + M^{*2})^{1/2} (f_B + \bar{f}_B), \quad (3a)$$

$$\begin{aligned} \Delta\epsilon_{\text{ren}} = & -\frac{\gamma}{16\pi^2} \left(M^{*4} \ln \left(\frac{M^*}{M} \right) + M^3 (M - M^*) - \frac{7}{2} M^2 (M - M^*)^2 \right. \\ & \left. + \frac{13}{3} M (M - M^*)^3 - \frac{25}{12} (M - M^*)^4 \right), \end{aligned} \quad (3b)$$

$$\epsilon_{\omega} = g_{\omega} \omega_0 \rho_B^{\text{ren}} - \frac{1}{2} m_{\omega}^2 \omega_0^2, \quad (3c)$$

$$\begin{aligned} \Delta\epsilon_{\sigma} = & \frac{1}{2} m_R^2 \sigma_0^2 + 3\lambda_R \sigma_0^4 + \frac{M_{\sigma}^4}{64\pi^2} \left(\ln \left(\frac{M_{\sigma}^2}{m_R^2} \right) - \frac{1}{2} \right) - 3\lambda_R I_f^2 \\ & - \frac{M_{\sigma,0}^4}{64\pi^2} \left(\ln \left(\frac{M_{\sigma,0}^2}{m_R^2} \right) - \frac{1}{2} \right) + 3\lambda_R I_{f0}^2, \end{aligned} \quad (3d)$$

as the mean field result, contribution from the Dirac sea, and contributions from the ω and σ mesons, respectively. We might note here that the quantum effects arising here from the scalar meson sector through σ meson condensates amounts to a sum over a class of multiloop diagrams and, does not correspond to the one meson loop approximation for scalar meson quantum effects considered earlier [20,22]. In the above, f_B and \bar{f}_B are the thermal distribution functions for the baryons and antibaryons, given in terms of the effective energy, $\epsilon^*(k) = (k^2 + M^{*2})^{1/2}$, and the effective chemical potential, $\mu^* = \mu - g_{\omega}\omega_0$. $M^* = M + g_{\sigma}\sigma_0$ is the effective nucleon mass and ρ_B^{ren} is the baryon number density after subtracting out the vacuum contribution. The spin-isospin degeneracy factor is $\gamma = 4$ for symmetric nuclear matter. In the expression for the energy density arising from the scalar meson sector, the field dependent effective sigma mass, $M_{\sigma}(\beta)$, satisfies the gap equation in terms of the renormalised parameters as

$$M_{\sigma}(\beta)^2 = m_R^2 + 12\lambda_R \sigma_0^2 + 12\lambda_R I_f(M_{\sigma}(\beta)), \quad (4)$$

where

$$I_f(M_{\sigma}(\beta)) = \frac{M_{\sigma}(\beta)^2}{16\pi^2} \ln \left(\frac{M_{\sigma}(\beta)^2}{m_R^2} \right) + \frac{1}{(2\pi)^3} \int \frac{d\mathbf{k}}{(k^2 + M_{\sigma}(\beta)^2)^{1/2}} \frac{1}{e^{\beta\omega'(\mathbf{k},\beta)} - 1} \quad (5)$$

where $\omega'(\mathbf{k}, \beta) = (k^2 + M_{\sigma}(\beta)^2)^{1/2}$. The effective meson mass M_{σ} is determined through a self consistent solution of Eq. (4). In Eq. (3d), $M_{\sigma,0}$ and I_{f0} are the expressions as given by Eqs. (4) and (5) with $\sigma_0 = 0$. We might note here that the gap equation given by Eq. (4) is identical to that obtained through resumming the daisy and superdaisy graphs [24,25] and hence includes higher-order corrections from the scalar meson field in a self-consistent manner.

Extremisation of the thermodynamic potential, $\Omega = -p = \epsilon_{\text{ren}} - \frac{\mathcal{S}}{\beta} - \mu\rho_B^{\text{ren}}$, with respect to the meson fields σ_0 and ω_0 yields the self-consistency conditions for σ_0 (and hence for the effective nucleon mass $M^* = M + g_{\sigma}\sigma_0$) and for the vector meson field ω_0 .

III. IN-MEDIUM PROPERTIES FOR ω AND ρ VECTOR MESONS

A. Vacuum polarisation

We now examine the medium modification to the masses and decay widths of the ω and ρ mesons including the quantum correction effects in hot nuclear matter in the relativistic random phase approximation. The interaction vertices for these mesons with nucleons are given as

$$\mathcal{L}_{\text{int}} = g_V \left(\bar{\psi} \gamma_\mu \tau^a \psi V_a^\mu - \frac{\kappa_V}{2M_N} \bar{\psi} \sigma_{\mu\nu} \tau^a \psi \partial^\nu V_a^\mu \right) \quad (6)$$

where $V_a^\mu = \omega^\mu$ or ρ_a^μ , M_N is the free nucleon mass, ψ is the nucleon field and $\tau_a = 1$ or $\vec{\tau}$, $\vec{\tau}$ being the Pauli matrices. g_V and κ_V correspond to the couplings due to the vector and tensor interactions for the corresponding vector mesons to the nucleon fields. The vector meson self energy is expressed in terms of the nucleon propagator, $G(k)$ modified by the quantum effects. This is given as

$$\Pi^{\mu\nu}(k) = -\gamma_I g_V^2 \frac{i}{(2\pi)^4} \int d^4p \text{Tr} \left[\Gamma_V^\mu(k) G(p) \Gamma_V^\nu(-k) G(p+k) \right], \quad (7)$$

where $\gamma_I = 2$ is the isospin degeneracy factor for nuclear matter, and $\Gamma_V^\mu(k) = \gamma^\mu \tau_a - \frac{\kappa_V}{2M_N} \sigma^{\mu\nu}$ represents the meson-nucleon vertex function.

The vector meson self energy can be written as the sum of two parts

$$\Pi^{\mu\nu} = \Pi_F^{\mu\nu} + \Pi_D^{\mu\nu}. \quad (8)$$

where $\Pi_F^{\mu\nu}$ is the contribution arising from the vacuum polarisation effects, described by the coupling to the $N\bar{N}$ excitations and $\Pi_D^{\mu\nu}$ is the density dependent contribution to the vector self energy. For the ω meson, the tensor coupling is generally small as compared to the vector coupling to the nucleons [13], and is neglected in the present work. The Feynman part of the self energy, $\Pi_F^{\mu\nu}$ is divergent and needs renormalization. We use dimensional regularization to separate the divergent parts and the renormalization condition $\Pi_F^{\text{ren}}(k^2)|_{M_N^* \rightarrow M_N} = 0$. Because of the tensor interaction, the vacuum self energy for the ρ meson is not renormalizable. We employ a phenomenological subtraction procedure [12,13] to extract the finite part using the renormalization condition $\left. \frac{\partial^n \Pi_F^\rho(k^2)}{\partial (k^2)^n} \right|_{M_N^* \rightarrow M_N} = 0$,

with $(n = 0, 1, 2, \dots \infty)$.

The effective mass of the vector meson is obtained by solving the equation, with $\Pi = \frac{1}{3} \Pi_\mu^\mu$,

$$k_0^2 - m_V^2 + \text{Re}\Pi(k_0, \mathbf{k} = 0) = 0. \quad (9)$$

B. Meson decay properties

The decay width for the process $\rho \rightarrow \pi\pi$ is calculated from the imaginary part of the self energy using the Cutkosky rule, and in the rest frame of the ρ meson is given by

$$\Gamma_\rho(k_0) = \frac{g_{\rho\pi\pi}^2}{48\pi} \frac{(k_0^2 - 4m_\pi^2)^{3/2}}{k_0^2} \left[\left(1 + f\left(\frac{k_0}{2}\right)\right) \left(1 + f\left(\frac{k_0}{2}\right)\right) - f\left(\frac{k_0}{2}\right) f\left(\frac{k_0}{2}\right) \right] \quad (10)$$

where, $f(x) = [e^{\beta x} - 1]^{-1}$ is the Bose-Einstein distribution function. The first and the second terms in the equation (10) represent the decay and the formation of the resonance, ρ . The medium effects have been shown to play a very important role for the ρ meson decay width. In the calculation for the ρ decay width, the pion has been treated as free, and any modification of the pion propagator due to effects such as delta-nucleon hole excitation [32] to yield a finite decay width for the pion, have not been taken into account. The coupling $g_{\rho\pi\pi}$ is fixed from the decay width of ρ meson in vacuum ($\Gamma_\rho=151$ MeV) decaying to two pions.

For the nucleon-rho couplings, we use the vector and tensor couplings as obtained from the N-N forward dispersion relation [13,15,33]. With the couplings as described above, we consider the temperature and density dependence of the ω and ρ meson properties in the hot nuclear matter due to quantum correction effects.

To calculate the decay width for the ω meson, we write down the effective Lagrangian for the ω meson as [34–36]

$$\mathcal{L}_\omega = -\frac{em_\omega^2}{g_\omega} \omega^\mu A_\mu + \frac{g_{\omega\pi\rho}}{m_\pi} \epsilon_{\mu\nu\alpha\beta} \partial^\mu \omega^\nu \partial^\alpha \rho_i^\beta \pi_i + \frac{g_{\omega 3\pi}}{m_\pi^3} \epsilon_{\mu\nu\alpha\beta} \epsilon_{ijk} \omega^\mu \partial^\nu \pi^i \partial^\alpha \pi^j \partial^\beta \pi^k. \quad (11)$$

In the above, the first term refers to the direct coupling of the vector meson ω to the photon, and hence to the dilepton pairs, as given by the vector dominance model. The decay width of the ω meson in vacuum is dominated by the channel $\omega \rightarrow 3\pi$. In the medium, the decay width for $\omega \rightarrow 3\pi$ is given as

$$\Gamma_{\omega \rightarrow 3\pi} = \frac{(2\pi)^4}{2k_0} \int d^3\tilde{p}_1 d^3\tilde{p}_2 d^3\tilde{p}_3 \delta^{(4)}(P - p_1 - p_2 - p_3) |M_{fi}|^2 \left[(1 + f(E_1))(1 + f(E_2))(1 + f(E_3)) - f(E_1)f(E_2)f(E_3) \right], \quad (12)$$

where $d^3\tilde{p}_i = \frac{d^3p_i}{(2\pi)^3 2E_i}$, p_i and E_i 's are 4-momenta and energies for the pions, and $f(E_i)$'s are their thermal distributions. The matrix element M_{fi} has contributions from the channels $\omega \rightarrow \rho\pi \rightarrow 3\pi$ (described by the second term in (11)) and the direct decay $\omega \rightarrow 3\pi$ resulting from the contact interaction (last term in (11)) [36–38]. For the $\omega\rho\pi$ coupling we take the value $g_{\omega\rho\pi} = 2$ according to Ref. [15], which is similar in value to the coupling taken in [37]. We fix the point interaction coupling $g_{\omega 3\pi}$ by fitting the partial decay width $\omega \rightarrow 3\pi$ in vacuum (7.49 MeV) to be 0.24. The contribution arising from the direct decay is upto around 15 % which is comparable to the results of [36,37].

With the modifications of the vector meson masses in the hot and dense medium, a new channel becomes accessible, i.e., the decay mode $\omega \rightarrow \rho\pi$ for $m_\omega^* > m_\rho^* + m_\pi$, which is not possible in the vacuum (since $m_\omega < m_\rho + m_\pi$). The in-medium decay width for this process is found from the cutting rules at finite temperature to be

$$\Gamma_{\omega \rightarrow \rho\pi} = \frac{g_{\omega\rho\pi}}{32\pi m_\omega^* m_\rho^* m_\pi^2} \lambda^{3/2}(m_\omega^*, m_\rho^*, m_\pi^2) \left[(1 + f(E_\pi))(1 + f(E_\rho^*)) - f(E_\pi)f(E_\rho^*) \right] \quad (13)$$

where $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2(xy + yz + xz)$ arises from phase space considerations, and f 's are the thermal distribution functions for the pions and ρ mesons. The first and second terms refer to the emission and absorption of pion and ρ meson.

IV. DILEPTON PRODUCTION

The pion annihilation ($\pi^+\pi^- \rightarrow e^+e^-$) dominates dilepton production in the low invariant mass region. The dropping of the vector meson masses in the hot and dense matter also gives rise to substantial contribution to the dilepton yield from the direct decays of the vector mesons into dileptons ($\rho \rightarrow e^+e^-$, $\omega \rightarrow e^+e^-$).

The thermal production rate per unit four volume for lepton pairs is related to the imaginary part of the one-particle irreducible photon self energy by [15,29,39]

$$\frac{dR}{d^4q} = \frac{\alpha}{12\pi^4} \cdot \frac{\text{Im}\Pi_\mu^\mu(q)}{q^2(e^{\beta q_0} - 1)}, \quad (14)$$

where $q^\mu = (q_0, \mathbf{q})$ is the 4-momentum of the virtual photon, and $\alpha = e^2/4\pi$ is the fine structure constant.

A. Pion Annihilation ($\pi^+\pi^- \rightarrow e^+e^-$)

The invariant mass distribution for the lepton pairs from pion annihilation is given by

$$\frac{dR}{dM} = \frac{M^3}{2(2\pi)^4} \left(1 - \frac{4m_\pi^2}{M^2}\right) \int dM_T M_T dy \sigma(q_0, |\mathbf{q}|) f(M_T \cosh y) \quad (15)$$

where, $q_0 = M_T \cosh y$, $|\mathbf{q}| = \sqrt{q_0^2 - M^2}$. In the above, $\sigma(q_0, |\mathbf{q}|)$ is the cross section for the pion annihilation, and is given by

$$\sigma(q_0, |\mathbf{q}|) = \frac{4\pi\alpha^2}{3M^2} \sqrt{1 - 4m_\pi^2/M^2} |F_\pi(q_0, |\mathbf{q}|)|^2, \quad (16)$$

where $F_\pi(q)$ is the electromagnetic form factor of the pion. The photon can also couple to the hadrons through vector mesons, which is the basis of the vector dominance model [34]. The photon can thus interact with the pions through a vector meson, which must be an isovector and is taken to be the ρ meson. The effective Lagrangian for the photon-pion-rho interaction can then be written as [15]

$$\mathcal{L}_{\text{int}} = -g_{\rho\pi\pi} \rho^\mu J_\mu - e A^\mu J_\mu - \frac{e}{2g_\rho} F^{\mu\nu} \rho_{\mu\nu}, \quad (17)$$

where $\rho_{\mu\nu}$ is the field tensor for the ρ field and $eJ_\mu = ie(\pi^-\partial_\mu\pi^+ - \pi^+\partial_\mu\pi^-)$. The coupling $g_{\rho\pi\pi}$ in the above can be determined from the decay $\rho \rightarrow \pi\pi$, and the direct photon-rho coupling g_ρ is determined from the process $\rho \rightarrow e^+e^-$ [40]. These, including a finite decay width for the ρ meson, then lead to the ρ -dominated pion form factor as

$$F_\pi(q^2) = 1 - q^2 \frac{g_{\rho\pi\pi}}{g_\rho} \frac{1}{q^2 - m_\rho^2 + im_\rho\Gamma_\rho}. \quad (18)$$

In the rest frame of the ρ meson, the invariant mass distribution of the lepton pair of invariant mass M from pion annihilation is given as

$$\frac{dR}{dM} = \frac{\sigma(M)}{(2\pi)^4} M^4 T \left(1 - \frac{4m_\pi^2}{M^2}\right) \sum_{n=1}^{\infty} \frac{K_1(nM/T)}{n} \quad (19)$$

with the cross section $\sigma(M)$ for the pion annihilation, given in terms of the pion form factor at $q^2 = M^2$. The summation over n arises from the Bose-Einstein distribution function for the ρ meson. Since we are mostly in the regime $M/T \gg 1$, the first term in the series, corresponding to the Boltzmannian distribution function is dominant, with quantum corrections being of the order of upto 10 %.

B. Vector meson decaying to dileptons

A vector meson V in a thermal bath can decay into leptons and photons, described by the final state $|f\rangle$ say, which escape the bath without thermalization. With q as the four momentum of the nonhadronic state $|f\rangle$, a resonance peak should appear in the invariant mass (M) plot for the number of lepton pair events, versus $q^2 = M^2$ at $M = m_V$, where m_V is the mass of the resonance in the thermal bath. Hence, the medium modifications of the vector mesons get reflected in the dilepton spectra through the shift and broadening of the peaks.

The invariant mass distribution of lepton pairs from the decays of vector mesons is given by a generalization of the Breit Wigner formula as [41]

$$\frac{dN_f}{d^4x d^4q} = (2J + 1) \frac{M^2}{4\pi^2} \frac{\Gamma_{\text{all}\rightarrow V} \Gamma_{V\rightarrow f}^{\text{vac}}(M)}{(M^2 - m_V^2)^2 + (m_V \Gamma_{\text{tot}})^2}. \quad (20)$$

In the above, dN_f is the number of lepton pair events, say in the space time and four momentum element $d^4x d^4q$, J being the spin of the resonance, $\Gamma_{\text{all}\rightarrow V}$ is the formation width, Γ_{tot} is the total width and $\Gamma_{V\rightarrow f}^{\text{vac}}$ is the partial width for the off-shell V (i.e. of mass M) to go into the nonhadronic state $|f\rangle$. The total width $\Gamma_{\text{tot}} = \Gamma_{V\rightarrow\text{all}} - \Gamma_{\text{all}\rightarrow V}$ is the rate at which the particles equilibrate and relax to chemical equilibrium. The invariant mass distribution of the lepton pairs from decays of vector mesons propagating with a momentum \mathbf{q} is obtained from equation (20) as

$$\frac{dR}{dM} = \frac{2J + 1}{2\pi^2} M \int M_T dM_T dy f(M_T \cosh y) \times \left[\frac{q_0 \Gamma_{\text{tot}}}{(M^2 - m_V^2 + \text{Re}\Pi)^2 + q_0^2 \Gamma_{\text{tot}}^2} \right] q_0 \Gamma_{V\rightarrow e^+e^-}^{\text{vac}} \quad (21)$$

In the rest frame of the vector meson, this reduces to

$$\frac{dR}{dM} = \frac{2J + 1}{\pi^2} M^2 T \sum_n \frac{K_1(nM/T)}{n} \times \frac{m_V^* \Gamma_{\text{tot}}/\pi}{(M^2 - m_V^{*2})^2 + m_V^{*2} \Gamma_{\text{tot}}^2} m_V^* \Gamma_{V\rightarrow e^+e^-}^{\text{vac}}(M). \quad (22)$$

In the above, $\Gamma_{V\rightarrow e^+e^-}^{\text{vac}}(M)$ is the partial width for the leptonic decay for the off-shell vector meson V with mass M , and is given as

$$\Gamma_{V\rightarrow e^+e^-}^{\text{vac}}(M) = \frac{4\pi\alpha^2 M}{3g_V^2} \quad (23)$$

The coupling g_V refers to the direct coupling of the vector meson V to the photon as given in the equations (11) and (17) for the ω and ρ mesons respectively. The coupling of ω to the photon taken to be one third of the coupling of ρ to the photon gives reasonable agreement with the ratio of the observed partial widths $\Gamma(\rho \rightarrow e^+e^-)/\Gamma(\omega \rightarrow e^+e^-)$ [40].

V. RESULTS AND DISCUSSIONS

In the following we present numerical results on the influence of vacuum effects on hadronic masses, decay widths of vector mesons, and dilepton emission from hot hadronic matter. As described in the previous sections, the study is based entirely on a hadronic scenario and makes use of the original Walecka model, extended by a nonlinear sigma self coupling.

In this work we concentrate on one particular aspect, namely on the role of scalar vacuum effects. For each quantity studied, we compare the results of the Relativistic Hartree Approximation (RHA), which was already treated in Ref. [15], with those additional effects arising from a scalar vacuum condensate. The quartic self coupling λ_R is taken as a free parameter; we have used $\lambda_R = 1.8$ and $\lambda_R = 5$ as representative values. The values of the incompressibility corresponding to these self couplings are 401 MeV (for $\lambda_R=1.8$) and 329 MeV (for $\lambda_R=5$). These values are smaller than the mean field result of 545 MeV as well as the relativistic Hartree result ($\lambda_R=0$) of 450 MeV. This is because an increase of λ_R leads to weaker potentials of the σ and ω mesons and to a softening of the equation of state.

The parameters of the model are adjusted to fit nuclear matter, corresponding to a saturation density of $\rho_0 = 0.193/\text{fm}^3$. The fitted values for the scaled coupling constants are $C_S^2 = g_\sigma^2 M_N^2 / m_\sigma^2 = 183.3, 167.5$ and 137.9 , and $C_V^2 = g_\omega^2 M_N^2 / m_\omega^2 = 114.7, 96.45$, and 63.7 for $\lambda_R = 0$ (i.e., RHA), 1.8, and 5, respectively.

When discussing dilepton production, we do not attempt to give a realistic description of specific experimental results. For this, the dynamic space-time evolution for the hot hadronic matter in the course of a nuclear collision would have to be addressed and additional effects contributing to the dilepton spectrum would have to be included. However, our study should help to elucidate, in which way in-medium vacuum effects can have an influence on experimentally observable quantities.

We first discuss the effect of quantum corrections on the in-medium nucleonic properties for hot nuclear matter [26]. In fig. 1a, we plot the temperature dependence of the nucleon mass for the baryon density $\rho_B = 0$ and the nuclear matter density $\rho = \rho_0$. The quantum corrections arising from the sigma field have the effect of softening the equation of state, and produce a higher value for the effective nucleon mass as compared to the RHA. The effective nucleon mass was seen to increase with the vacuum polarisation effects arising from the baryons as compared to the MFT calculations [21,23]. Though for $\rho_B = 0$, the effective nucleon mass decreases monotonically with temperature, for $\rho_B = \rho_0$, the nucleon mass first increases and then decreases as a function of temperature, which has also been observed earlier to be the trend for higher densities [42]. The variation in the nucleon mass is very slow up to a temperature $T \simeq 150$ MeV, beyond which there is a fast decrease. Though qualitative features are similar for the in-medium mass with inclusion of quantum effects from the scalar meson sector, it is observed that such effects lead to a higher value for the nucleon mass.

We next study the temperature and density dependence of the vector mesons in hot nuclear matter. The ω meson mass as a function of the temperature for $\rho_B=0$ and ρ_0 , is plotted in Fig. 1b. In the Walecka model, the Dirac sea has been shown to have a significant contribution over the Fermi sea, leading to a substantial drop in the omega meson mass in nuclear matter [12,13,15]. Specifically, at saturation density and at zero temperature, the decrease in the omega meson mass from the vacuum value, is around 150 MeV in RHA, whereas with sigma quantum effects, the drop in the mass reduces to around 117 MeV or 61 MeV for the two values of λ_R studied. A similar reduction in the mass drop, as compared to the RHA was also observed when scalar field quantum corrections were included within a one-loop approximation [20,22], unlike the approximation adopted here, which is self-consistent and includes multi-loop effects. This indicates that the quantum corrections do play an important role in the medium modification of the vector meson masses. In the present work, it is seen that the quantum corrections from the scalar mesons, over those arising from the baryonic sector, have the effect of giving rise to a higher value for the ω meson mass. This again is related to the fact that the effective nucleon mass is higher when such quantum effects are taken into account.

In figure 2, we illustrate the medium modification for the ρ meson mass with the vector and tensor couplings to the nucleons being fixed from the NN forward dispersion relation [13,15,33]. The values for these couplings are given as $g_{\rho N}^2/4\pi=0.55$ and $\kappa_\rho=6.1$. We notice that the decrease in the ρ meson mass with increase in temperature is much sharper than that of the ω meson. Such a behaviour of the ρ meson undergoing a much larger medium modification was also observed earlier [15] within the relativistic Hartree approximation for the nucleons. This indicates that the tensor coupling, which is negligible for the ω meson, plays a significant role for the ρ meson. With inclusion of quantum corrections from the scalar meson, the qualitative comparison between the ω and ρ vector mesons remains the same, though these effects are generally seen to lead to larger values for the vector meson masses. We have also calculated the contribution to the ρ meson mass arising from pion loops through the $\rho\pi\pi$ coupling, but this was found to lead to a marginal increase of at most 1%, in agreement with [10,29,43].

In figure 3, the decay widths for the ρ meson with and without the scalar quantum corrections, over the relativistic Hartree approximation, are considered. Such quantum effects lead to an appreciable increase in the decay widths, which, as we shall see, is observed as a broadening of the ρ -peak in the dilepton spectra.

In figure 4, the decay width for the ω meson is seen to have a positive contribution arising from the σ -meson quantum effect as compared to RHA. The decay width is calculated as arising from the processes $\omega \rightarrow 3\pi$ as well as $\omega \rightarrow \rho\pi$, a channel which becomes possible with the mass modifications of the vector mesons in the medium. It is seen that the channel $\omega \rightarrow \rho\pi$ opens for temperature greater than around 180 MeV for $\rho_B = 0$, whereas it is observed even at zero temperature at nuclear matter saturation density. We might note that the collisional modification to the ω width, arising from the process $\omega\pi \rightarrow \pi\pi$ [15,30], has not been taken into account in the present calculations.

The dilepton spectra for the low invariant mass regime, at $T=200$ MeV are shown in figures 5 and 6 calculated for densities $\rho_B = 0$ and $\rho_B = \rho_0$. The total dilepton spectra along with contributions arising from pion annihilation and direct decays of the ρ and ω mesons, are displayed in figures 5a and 6a. The effect of sigma condensates shifts the peak position for the vector mesons to higher values as compared to RHA, and also results in a broadening of the peaks in the dilepton spectra as can be seen from figures 5b and 6b. The ρ and ω peaks are well separated due to the fact that the ρ meson mass gets much larger medium modification as compared to the ω meson, which would not be the case in the absence of medium effects. The omega peak is seen to be smeared due to scalar meson quantum corrections. The broadening of the ω peak due to σ -quantum effects is more clearly visible for $\rho_B = \rho_0$ than for zero density, in the dilepton spectra. The dominant contributions to the dilepton enhancement are from pion annihilation and ρ meson decay, which are plotted in figures 5c and 6c in comparison with the RHA results. The effect of σ condensates leads to a higher value for the peak position as well as broadening of the ρ -peak. The dilepton production rates due to the ω meson decay are displayed in figures 5d and 6d. It is observed that the broadening of the ω peak due to scalar quantum corrections is larger for ρ_0 than for zero density. This is a reflection of the fact that there is a larger modification of the omega width due to scalar quantum effects for higher densities. There have been various mechanisms under discussion which can lead to an increase in the ω width. It has been argued that the ω decay width can increase by an order of magnitude due to the reduction in the pion decay constant related with chiral symmetry restoration [44]. The medium modification due to scalar meson quantum effects giving rise to higher values for the ρ meson mass and decay width are reflected in the dilepton spectra.

The medium modification for the vector meson masses depend more sensitively on the density than on temperature. This can give rise to an enhancement of the dilepton production rate when matter is highly compressed at the initial stage of a heavy ion collision. In the DLS experiments for A-A collisions at 1 AGeV, there have been models that focus on density induced change for the ρ meson mass to describe the dilepton pair production for invariant mass of dilepton pair, $M \geq 0.4\text{GeV}$ [45,46]. Even though there is dilepton enhancement arising from the density-modified ρ mass, the results still underestimate the DLS data by a factor of 2-3 [46].

The mass modifications for the vector meson masses with density are plotted in figure 7, for temperature $T=100$ MeV. The trend is a sharp decrease of the masses with density up to $\rho_B \simeq \rho_0$, and then a rise for higher values of ρ_B . The observed trend is due to the competing effects of vacuum polarisation which leads to a reduction in the mass, and the density dependent dressing for the vector self energy causing an increase in the meson mass. Such a behaviour of the vector meson mass for the ω meson was also observed earlier [14]. A higher value of the scalar self interaction coupling, λ_R , leads to a higher value for the vector meson masses. The change in the ρ meson mass for higher densities is seen to be rather slow which means that there is a delicate cancellation between the Feynman and the density dependent parts of the vector self energy. As depicted in figure 7, the ω meson mass seems to be more dominated by the density dependent contributions for higher densities. The decay widths are plotted in figure 8. The slow change in ρ meson mass for higher densities is reflected in the in-medium ρ -decay width, while the ω decay width Γ_ω grows strongly with density and gets strongly enhanced by quantum effects from the σ meson.

The dilepton spectra for $T = 100$ MeV and $\rho_B = 2\rho_0$ are plotted in figures 9 and 10. These values should be typical for the compression zone of a nuclear collision at energies of a few GeV per nucleon. For example, for inclusive Ca+Ca collisions at 1 AGeV studied by the DLS collaboration [3], maximum temperatures and densities in this range were extracted from a transport model [47]. In figure 9a, the total dilepton spectrum is plotted, for $\lambda_R=1.8$, displaying the individual contributions. Figure 9b shows that the sigma quantum corrections lead to a broadening of the peaks in the dilepton spectra, which is seen to be more appreciable for the higher value of the scalar self coupling, λ_R . The individual contributions to the dilepton spectra, plotted in figure 10, illustrate the effect of σ quantum corrections.

Finally, the density dependence for the dilepton spectra is shown for temperature $T = 100$ MeV in figures 11 and 12 for couplings λ_R as 1.8 and 5, respectively, in comparison with the RHA results. The vector meson masses decreasing with density up to around twice the nuclear matter density and then slowly increasing or staying constant, are reflected in the dilepton spectra through the peak positions. One might notice that the broadening of the peaks due to σ - quantum effects is larger for higher densities and is more pronounced for the ω peak. The curves for densities twice and four times nuclear matter density in figure 12 are coincident up to a value of the invariant mass of about 0.6 GeV. This is due to the fact that the rho meson properties whose effects are relevant for that mass region, remain almost constant for $\lambda_R=5$. Beyond that however, they are different around the invariant mass region for the omega meson. Due to substantial increase in the decay width of the ω meson with density the ω peak is almost smeared out even at nuclear matter density and the omega loses its quasi particle nature. The corrections from the scalar meson

quantum effects are thus seen to be appreciable.

VI. SUMMARY

To summarize, we have considered in the present paper, the modifications of the vector meson (ρ and ω) properties due to quantum correction effects in hot and dense nuclear matter and their effect on dilepton spectra. The baryonic properties, as modified due to such effects, subsequently determine the vector meson masses and decay widths in the hot and dense hadronic matter. It has been recently emphasized that the Dirac sea contribution dominates over the Fermi sea, thus implying that the vacuum polarisation effects arising from the baryonic sector do play an important role in the vector meson properties in a medium. We consider the hadronic properties as modified arising due to the quantum effects from the baryon and scalar meson sectors and investigate their effect on the dilepton production rate. The effect of the scalar meson quantum corrections leads to an increase in the vector meson masses and decay widths. The effect on the omega meson decay width is particularly more pronounced at higher densities, than in the RHA. The value for $T = 100$ MeV is modified from 67 MeV (for RHA) to 125 MeV (for $\lambda_R=1.8$) and 251 MeV (for $\lambda_R=5$) for $\rho_B/\rho_0=2$, and for $\rho_B/\rho_0=4$ the RHA value of 101 MeV is modified to 175 MeV and 370 MeV, for $\lambda_R=1.8$ and 5, respectively .

In the present work, we have studied the effect of quantum corrections arising from the scalar meson over those resulting from the baryon sector (given through the summing over baryonic tadpole diagrams in the relativistic Hartree approximation). The quantum corrections from the σ -meson give rise to a higher value for the vector meson masses as compared to keeping only quantum effects from the baryons. There is also an appreciable increase in the decay widths in the presence of these quantum corrections, and this is reflected in the invariant mass plot for dilepton spectra through the broadening of the peaks. For the ω meson, an appreciable broadening arising due to sigma meson quantum effects leads to the disappearance of the peak from the spectra for high densities when ω ceases to exist as a quasi particle. The density dependence of in-medium properties of the vector mesons and their subsequent effects on the dilepton spectra have been studied, which will be relevant for highly compressed matter at the early stage of a heavy ion collision.

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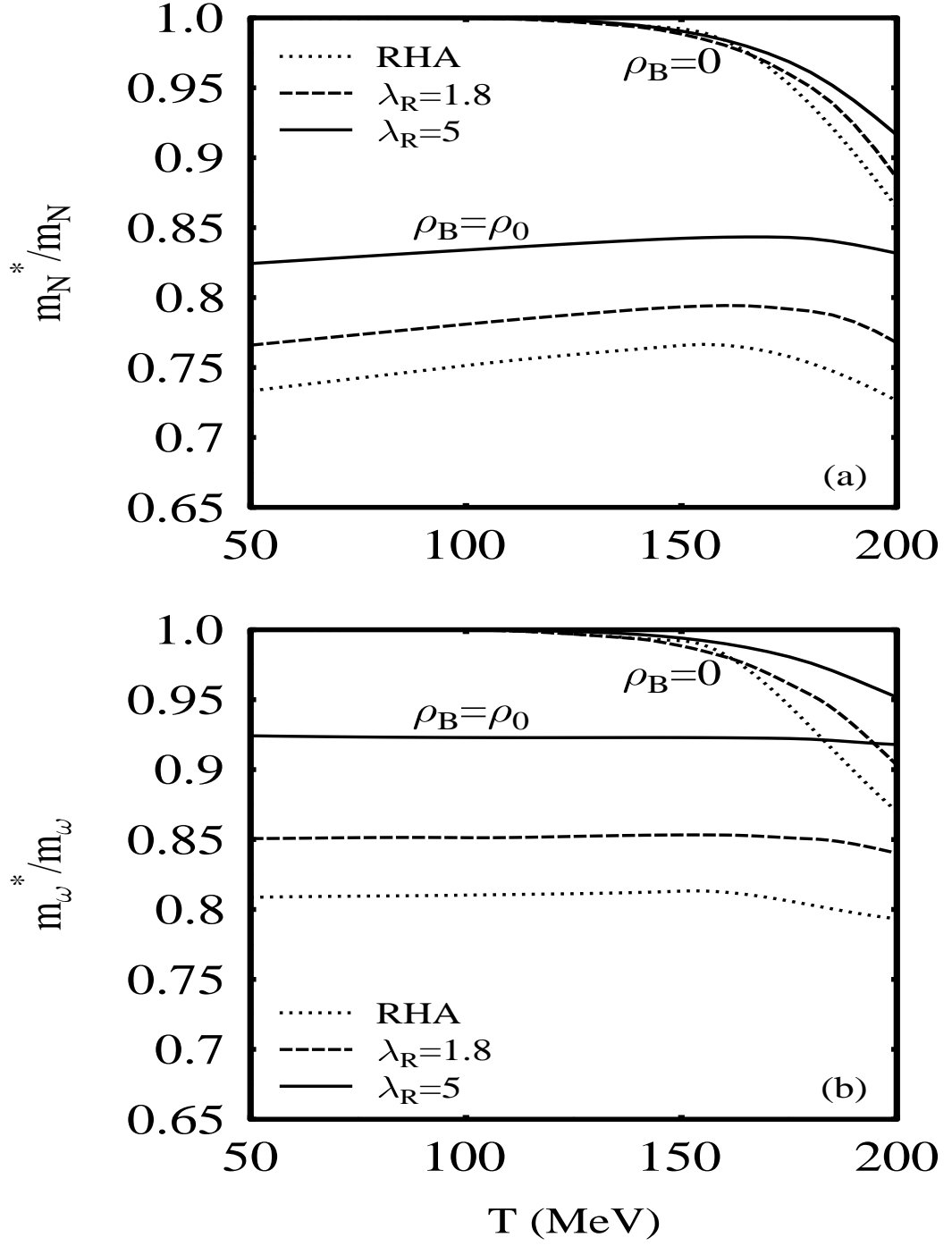


FIG. 1. Effective nucleon and omega meson masses as functions of the temperature for zero and nuclear matter density for RHA, $\lambda_R = 1.8$ and $\lambda_R = 5$. The quantum corrections from the scalar meson sector leads to an increase in the in-medium masses.

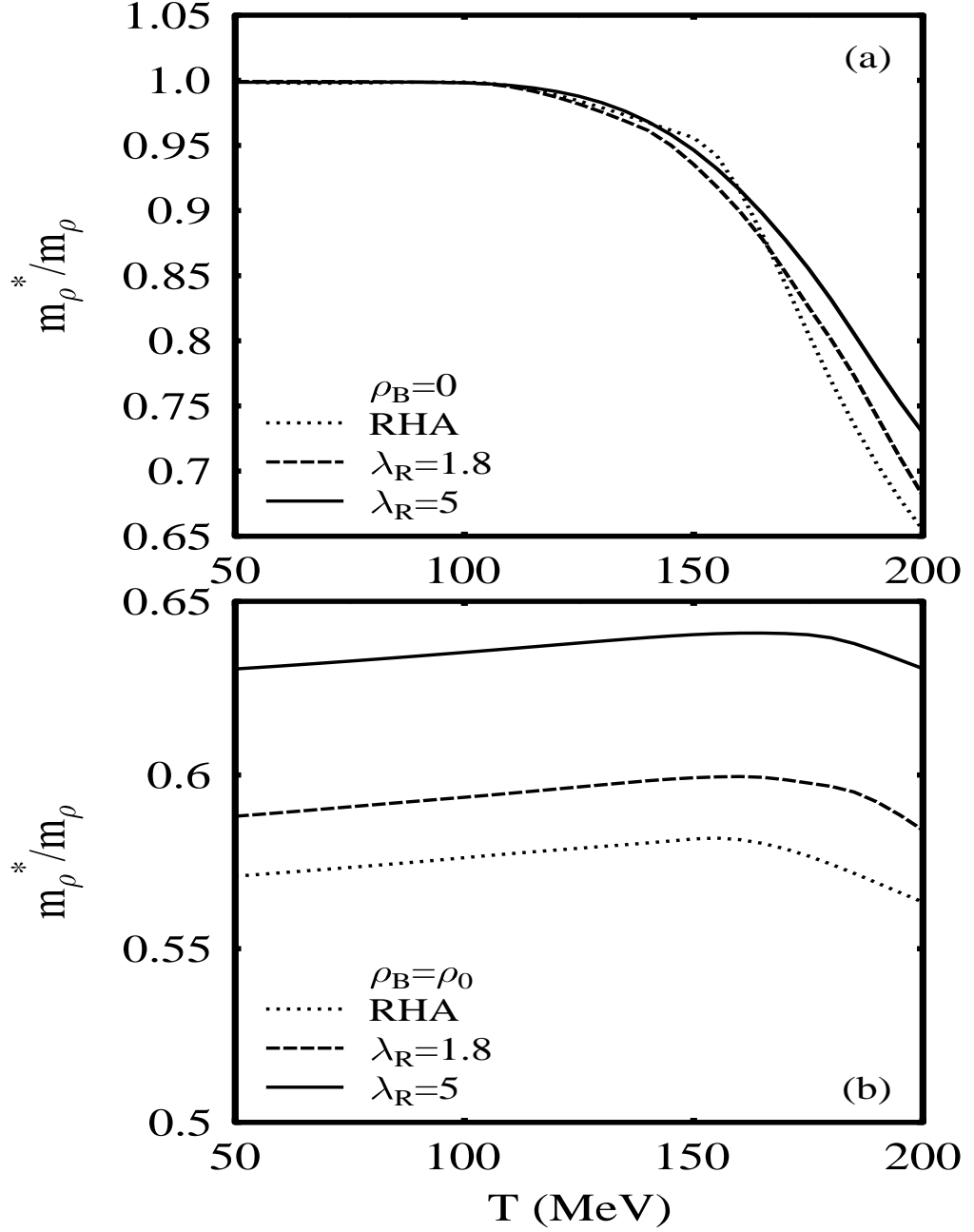


FIG. 2. Effective ρ meson mass as a function of the temperature for zero and nuclear matter densities for the three cases RHA, $\lambda_R = 1.8$ and $\lambda_R = 5$. The quantum corrections from the scalar meson sector increases the in-medium mass.

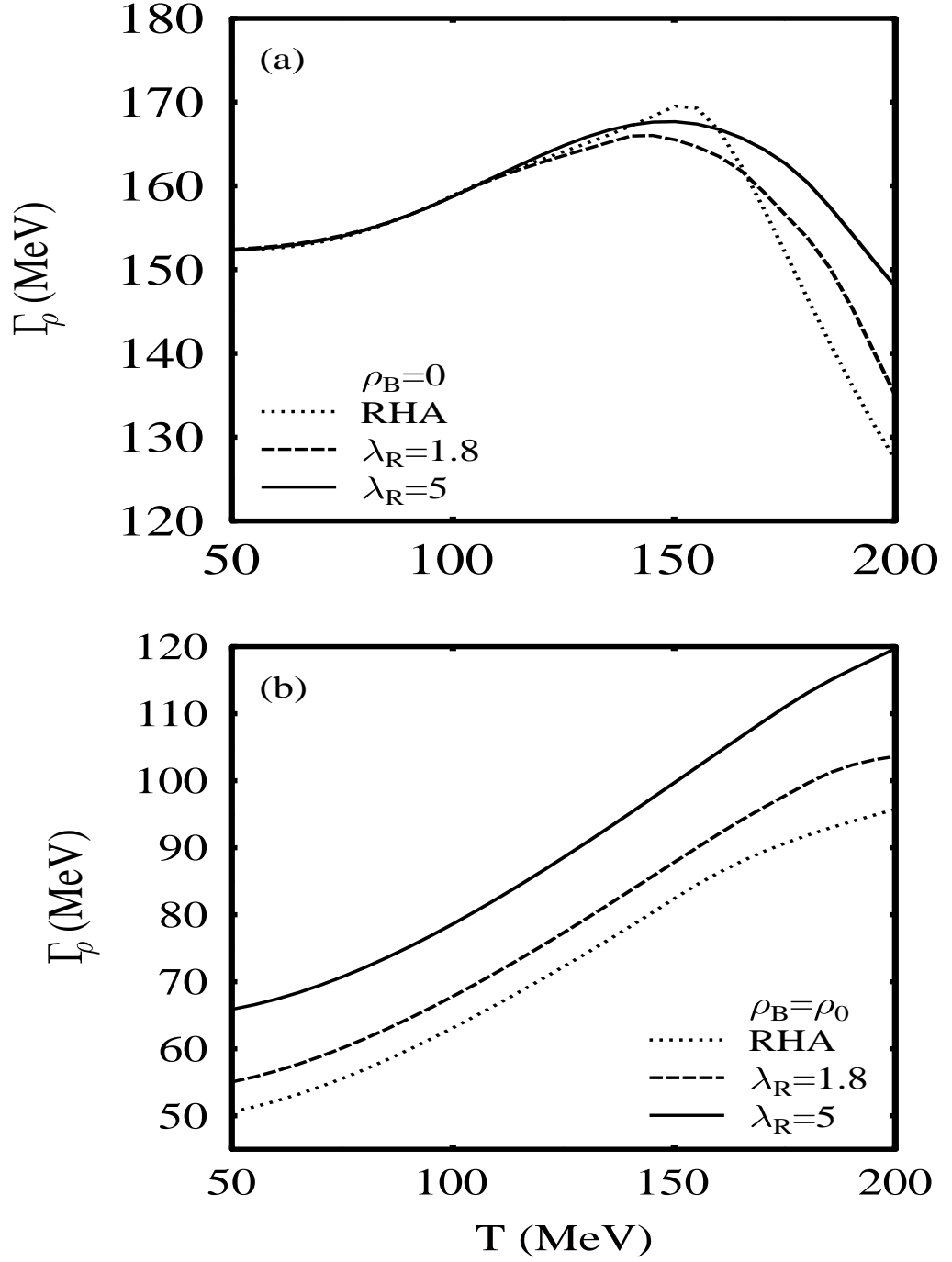


FIG. 3. In-medium decay width for the ρ meson as a function of temperature and density.

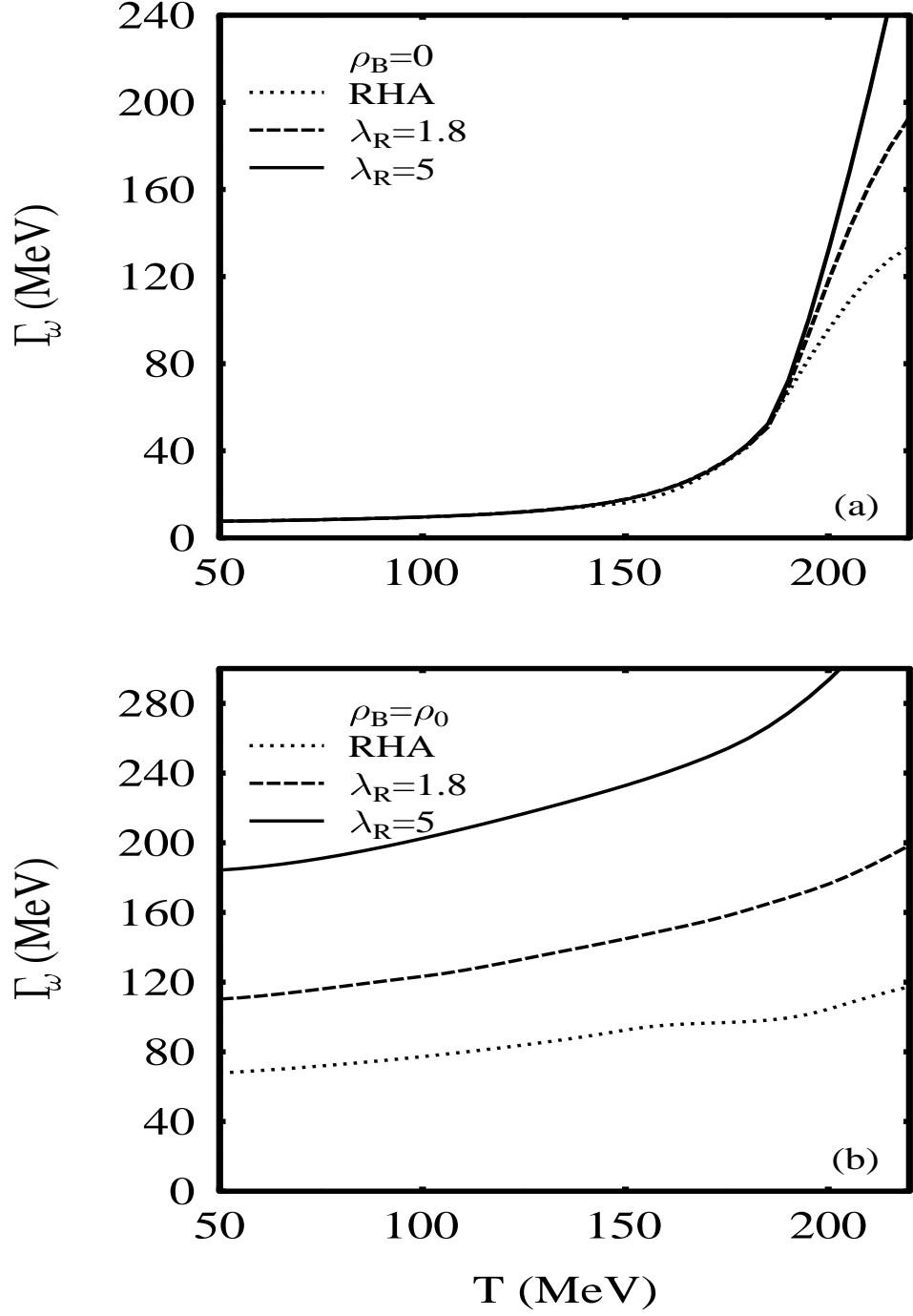


FIG. 4. In-medium decay width for the ω meson arising from the channels $\omega \rightarrow 3\pi$ and $\omega \rightarrow \rho\pi$, the latter becoming accessible because of the medium modification of the meson masses.

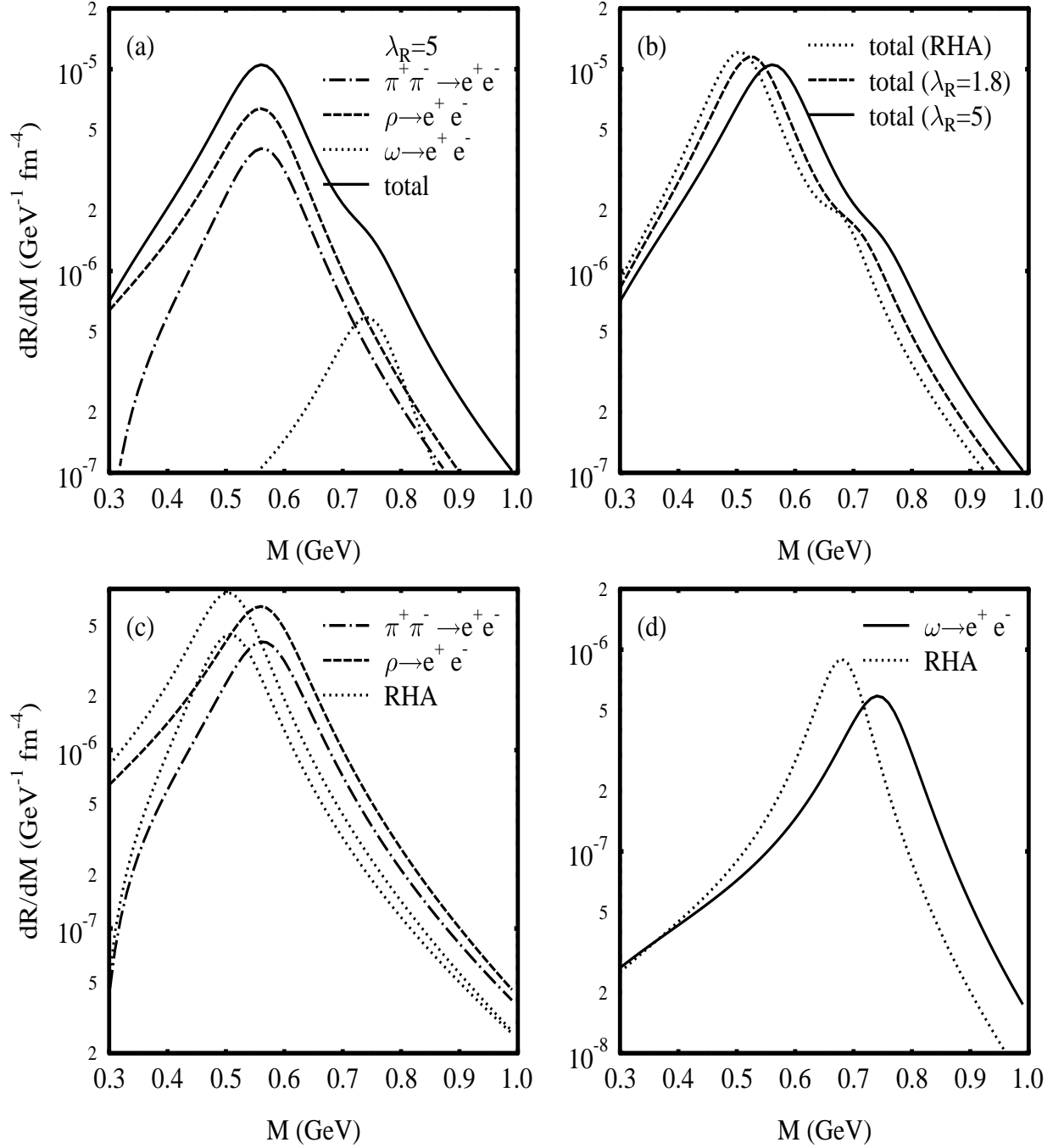


FIG. 5. Total dilepton production rate, as well as the individual contributions from the pion annihilation and direct decays of the ρ and ω mesons are shown, with and without scalar meson quantum effects for $\rho_B = 0$ and $T = 200$ MeV. One observes a positive shift as well as a broadening of the peaks due to σ quantum effects.

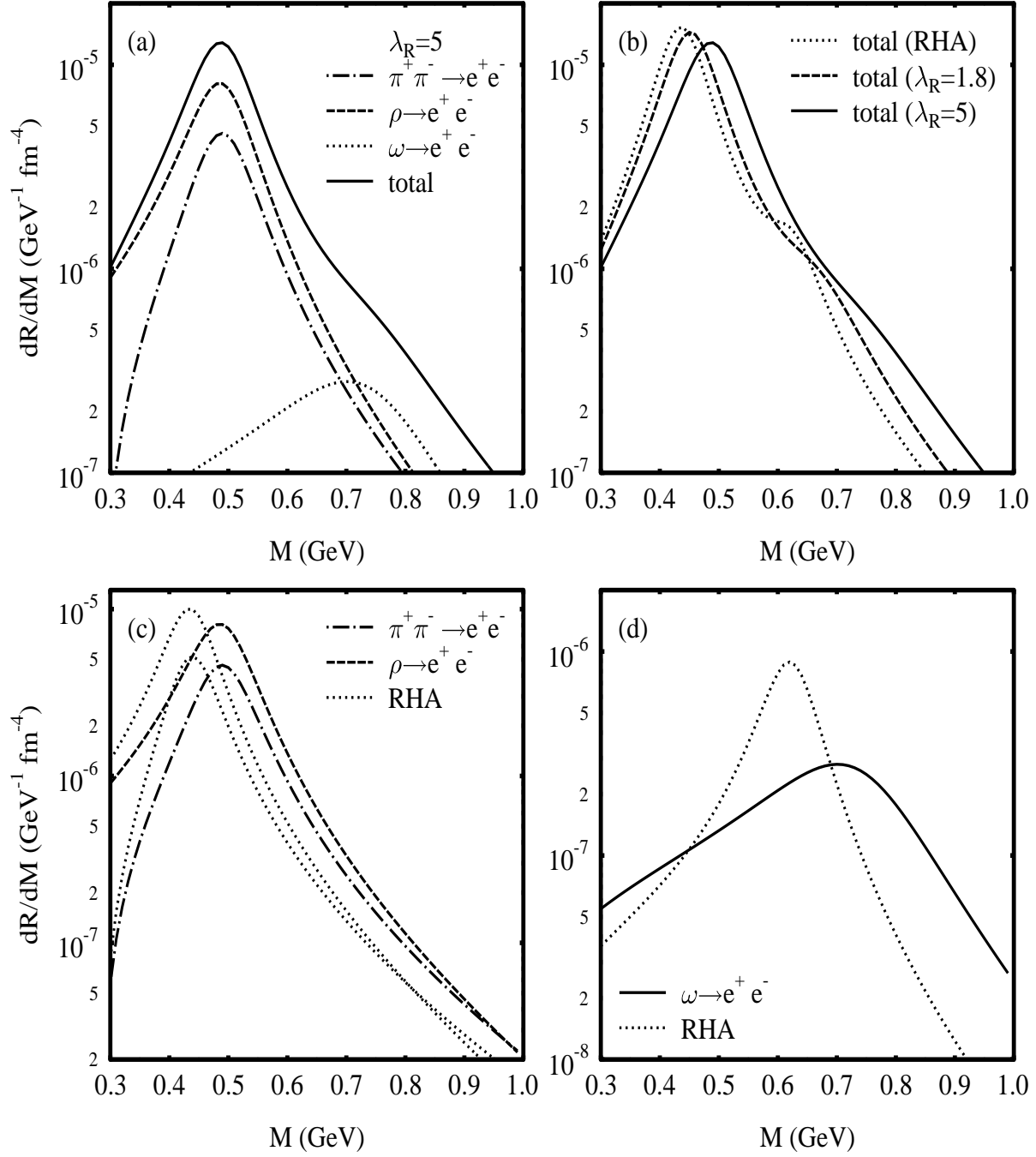


FIG. 6. Same as Fig. 5 for $\rho_B = \rho_0$

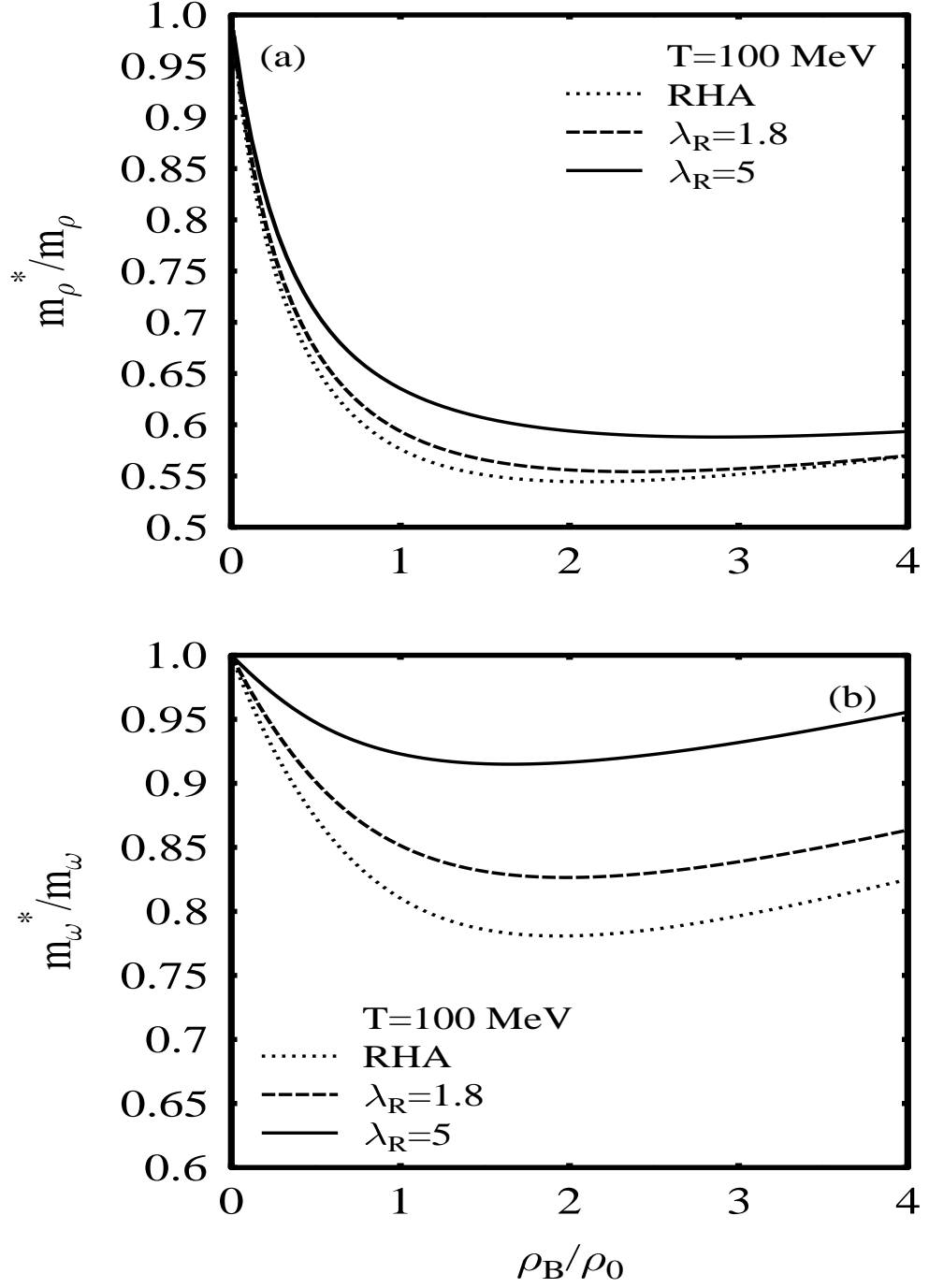


FIG. 7. Density dependence of the ρ and ω meson masses for $T = 100$ MeV. The masses are seen to depend more sensitively on density than on temperature. The σ quantum effects lead to an increase in the in-medium masses.

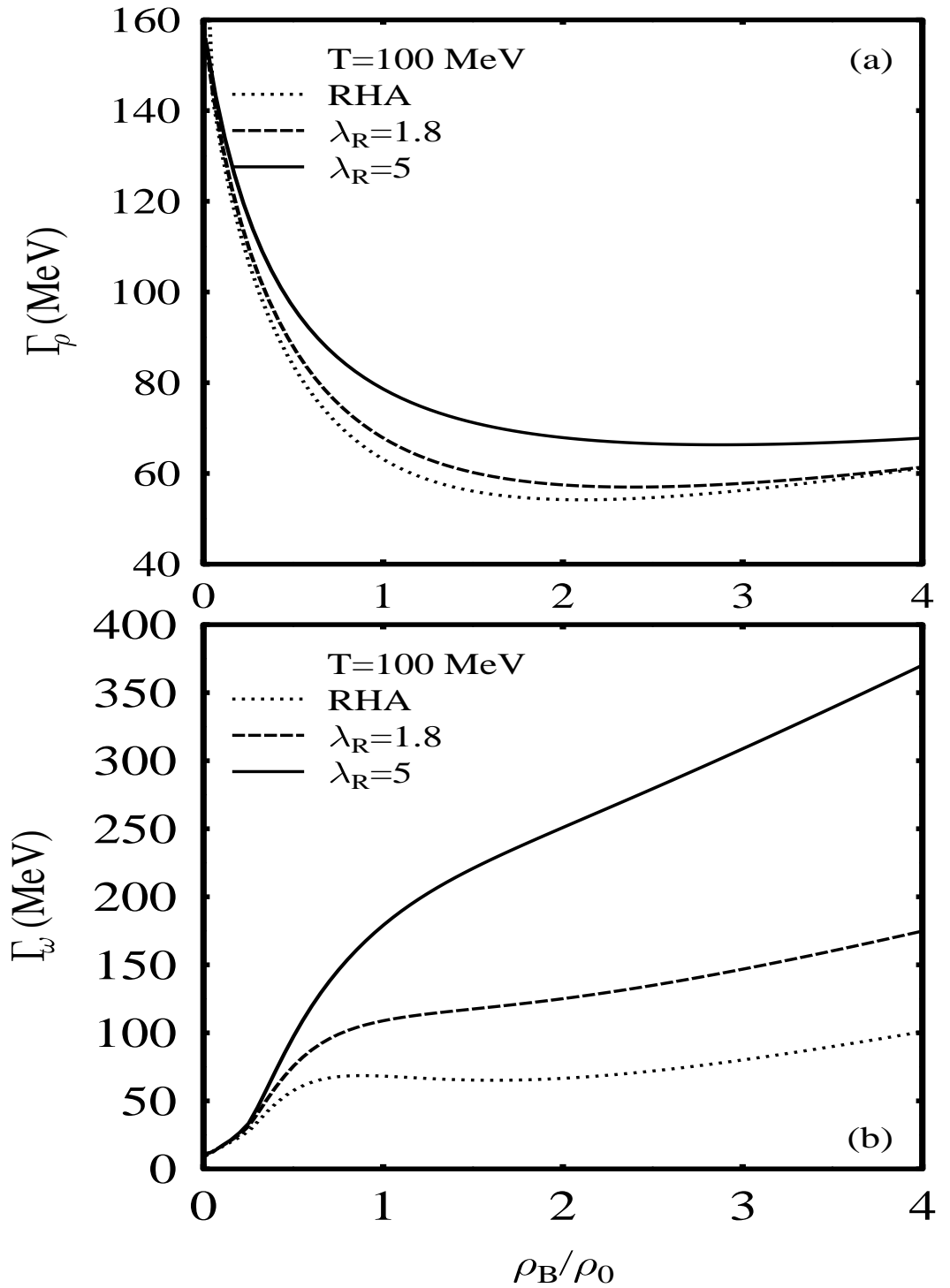


FIG. 8. Density dependence of the ρ and ω meson widths. With growing strength of the scalar self coupling, λ_R , the widths increase appreciably.

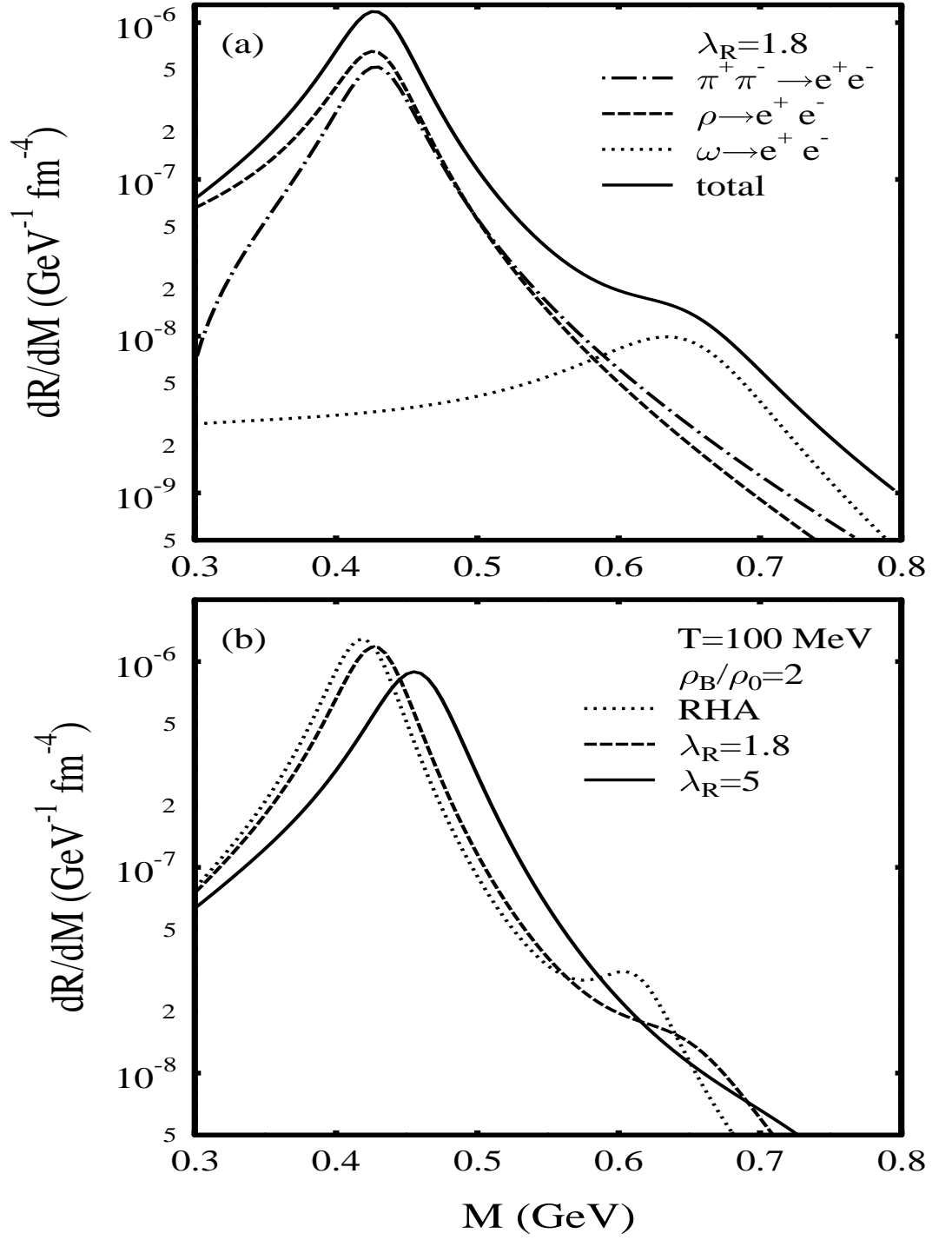


FIG. 9. Dilepton spectra for $T = 100$ MeV and $\rho_B = 2\rho_0$. The quantum effects from σ lead to a positive shift and broadening of the ρ and ω peaks. For $\lambda_R = 5$ the ω peak is no longer visible.

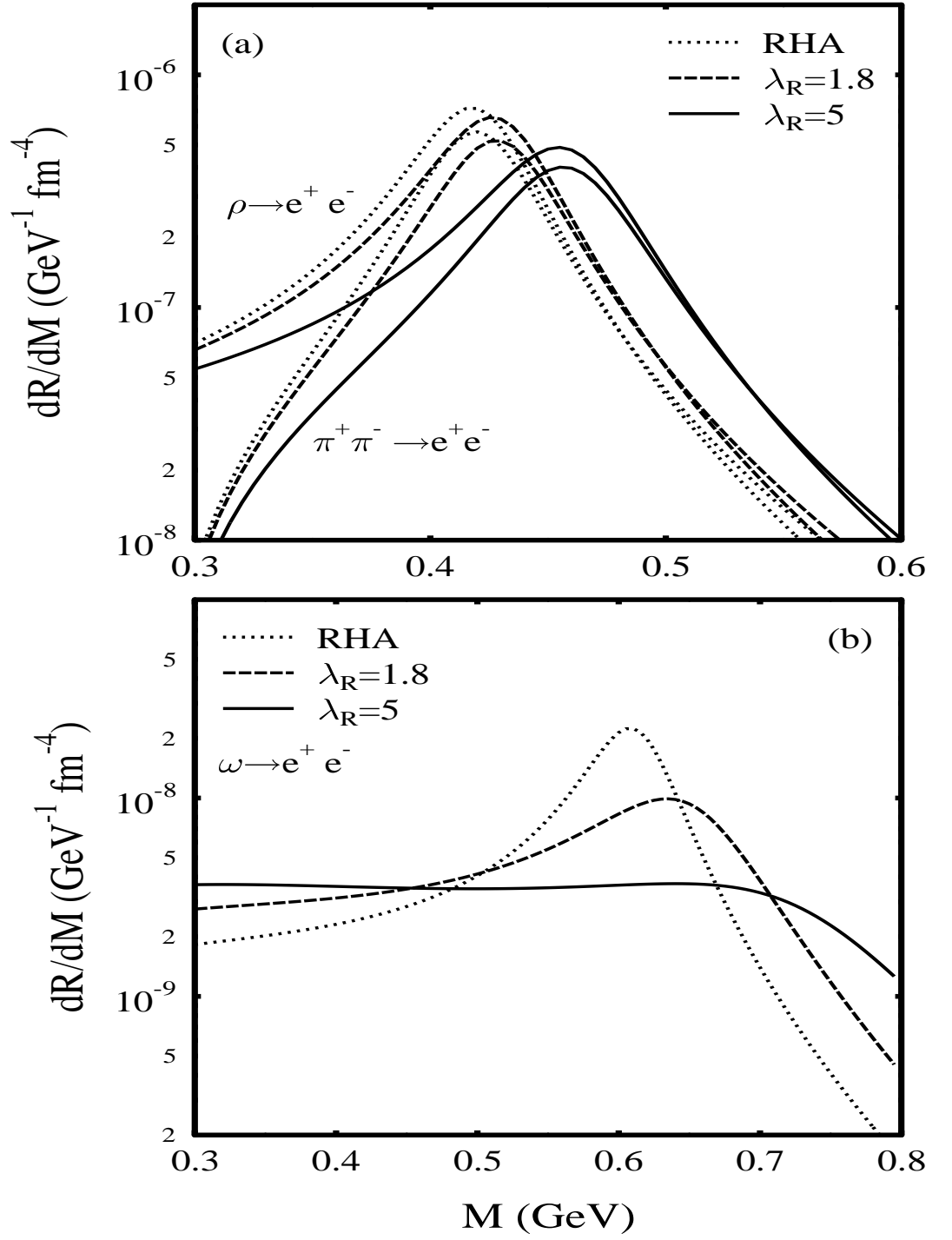


FIG. 10. Contributions from the processes of pion annihilation and decay of vector mesons to the dilepton spectra. For the higher value of λ_R the ω meson peak is completely smeared out.

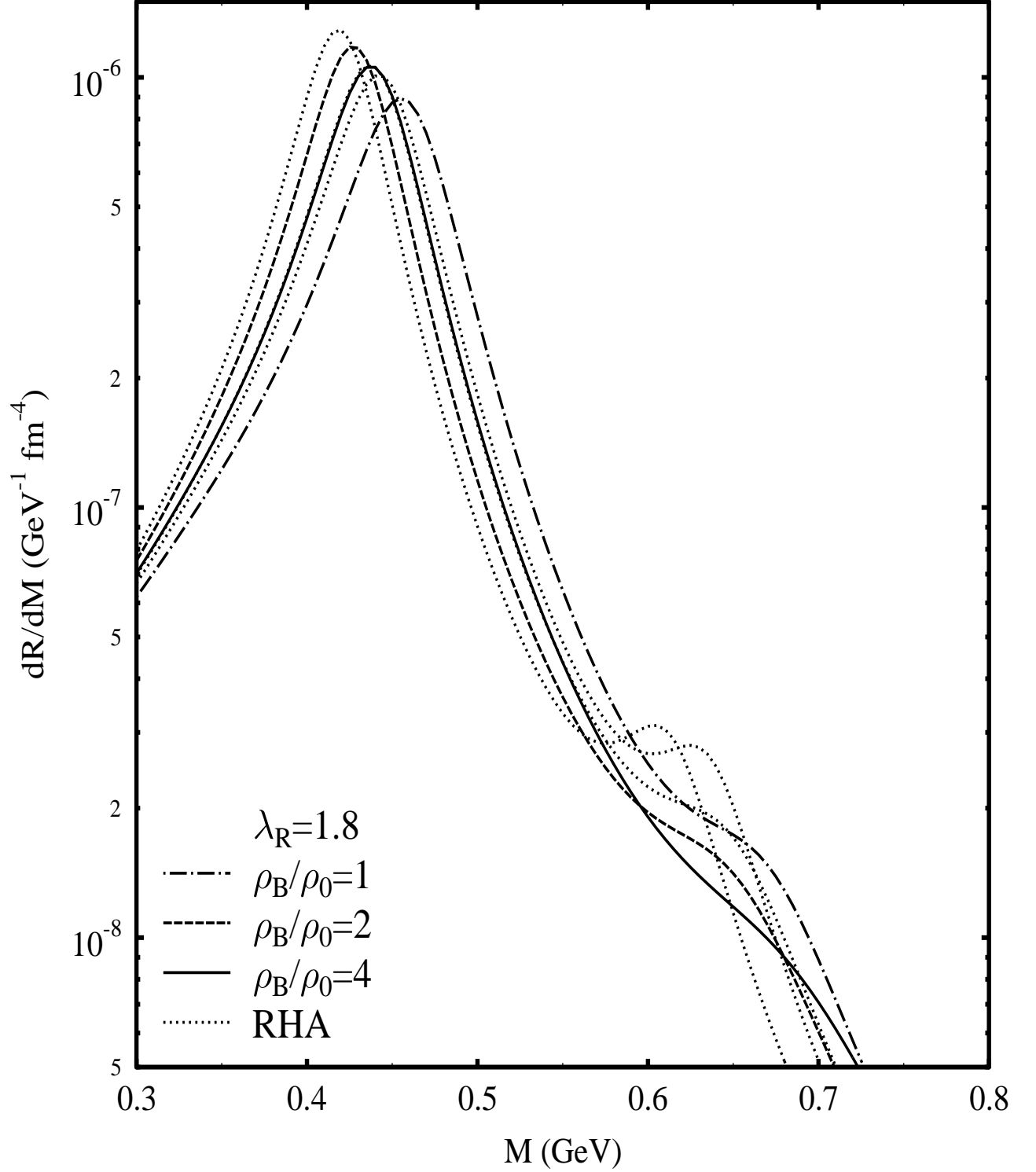


FIG. 11. Density dependence of dilepton spectra peak positions and widths for $T = 100$ MeV. Scalar quantum effects lead to a broadening of the peaks. For the ω meson, the broadening is more pronounced for higher densities.

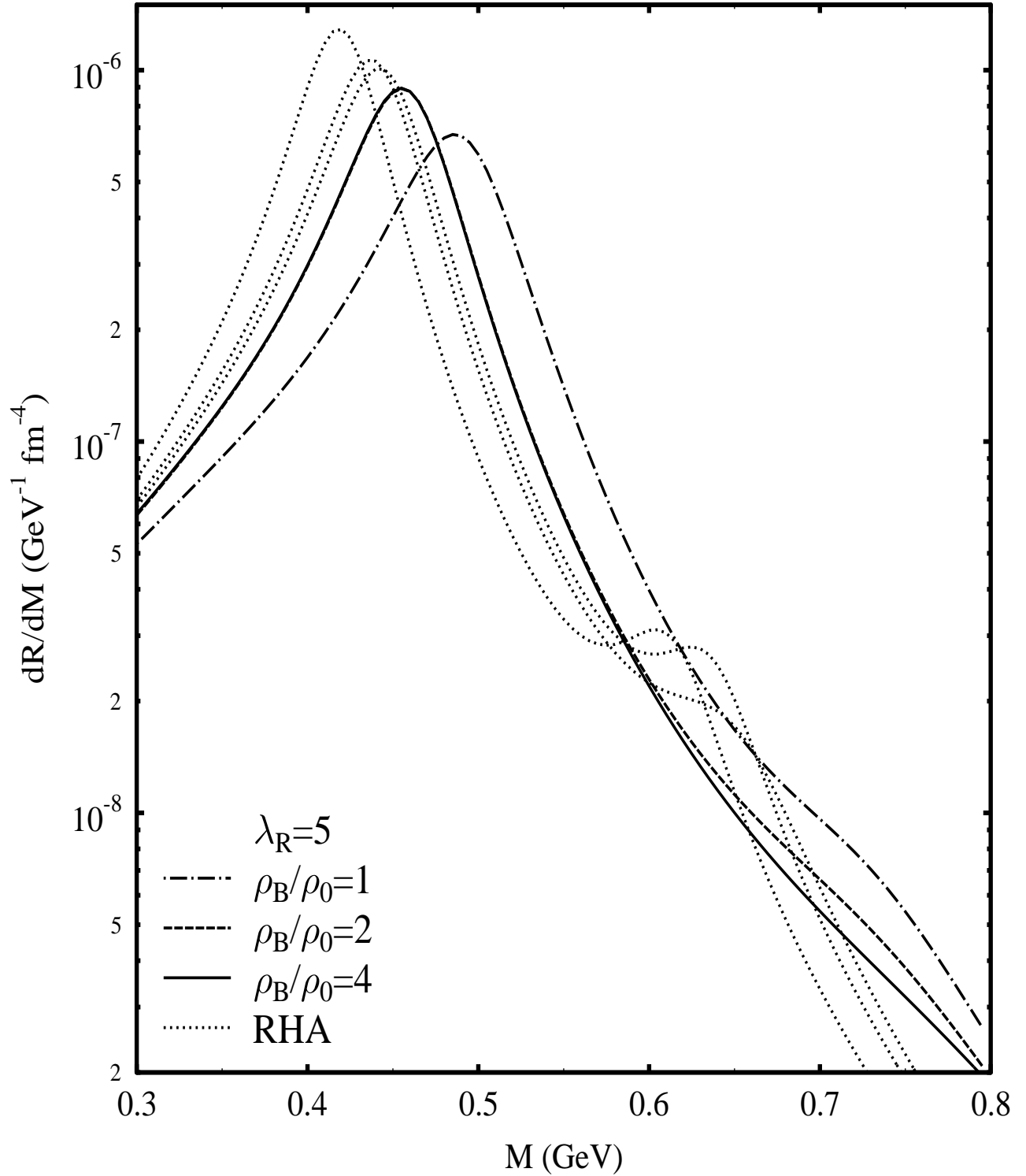


FIG. 12. Same as Fig. 11 for $\lambda_R = 5$. This leads to the disappearance of the ω peak even at nuclear saturation density.