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In-medium effects at low temperatures in the photon bremsstrahlung

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The influence of the dense nuclear medium on the process of photon bremsstrahlung in nuclear collisions at low excitation energies is investigated. The in-medium nucleon-nucleon scattering $\mathcal{T}$ matrix is incorporated in the calculations of photon bremsstrahlung cross section for nucleon-nucleus and nucleus-nucleus scattering cases. Calculations are performed in the soft photon approximation. It is shown that the Pauli blocking of intermediate scattering states and quasiparticle energy shifts due to the mean field interaction essentially suppress the photon bremsstrahlung cross section for temperatures $T \leqslant 10$ MeV and densities $n \leqslant n_s$ (n_s being the nuclear saturation density). In the case of nucleus-nucleus collisions, when the photons are emitted from the thermalized zone, a temperature and density dependent reduction factor for the bremsstrahlung rate is calculated. It implies a typical bremsstrahlung rate which is $\sim 15\%$ of that in the free scattering case.

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I. INTRODUCTION

The process of photon production by the nucleon-nucleon bremsstrahlung in the course of intermediate-energy heavy-ion collisions is traditionally regarded as a direct probe of the reaction dynamics. Because of the weak electromagnetic coupling, energetic photons ($\omega_\gamma \geqslant 30$ MeV) created mainly in incoherent neutron-proton (np) collisions are not strongly affected by the nuclear medium and may carry undisturbed information at the early stage of reaction. Thus the measurement of energetic photons provides a possibility for reconstructing the dynamics of heavy-ion collisions at intermediate energies and investigating the properties of hot and dense nuclear matter.

Early experimental data on photon spectra from $np\gamma$ bremsstrahlung and photon production in low- and intermediate-energy heavy-ion reactions was based on inclusive measurements, which were done with a rather limited statistics and involved only comparatively light colliding systems. This experimental situation was summarized in the review article [1]. The observed general trends of hard-photon emission from nuclear collisions in this energy range were quite well understood theoretically (see [2] and references therein). More recently, energetic photon production revived interest as a result of the latest experimental results [3–5] obtained with high statistical accuracy which allow one to measure hard photons in coincidence with fragments or/and light particles, to investigate the impact parameter dependence of these characteristics, to extend measurements to very heavy colliding systems, and even to make first photon-photon coincidence experiments [6], which throws light on the space-time evolution of nuclear collisions.

The bremsstrahlung emission in elementary nucleon-nucleon collisions, which can serve as a probe of the nucleon-nucleon potentials, is a basic subprocess in the consideration of incoherent nuclear bremsstrahlung and was theoretically investigated in a number of works; see the review [2]. However, when photons are emitted from a hot and dense source formed in the course of a nucleus-nucleus collision, the influence of the dense medium on the nucleon-nucleon scattering process has to be taken into account. Several aspects of the problem of photon bremsstrahlung in a nuclear medium, such as the validity of the so-called soft photon approximation, off-energy-shell effects in the nucleon-nucleon scattering, softening of the one-pion-exchange mode of nucleon-nucleon interaction, have been discussed; see Refs. [2,7–12]. However, no attempts were made to treat the problem of photon bremsstrahlung beyond the mean field description of the strong nuclear interaction which includes the full two-body scattering $\mathcal{T}$ matrix obtained from the many-body theory. The two-nucleon in-medium scattering states and correlations in dense nuclear matter in other contexts have been investigated in Refs. [9,13–21]. Most of these calculations are restricted to the nonrelativistic many-body theory (see Refs. [9,17,19] for relativistic approaches to the in-medium $\mathcal{T}$ matrix). Restricting ourselves to the low-temperature regime, we will employ the finite-temperature Matsubara Green functions technique with nonrelativistic propagators [22,23].

The goal of the present paper is to investigate the influence of the dense medium on the photon bremsstrahlung reaction cross section and photon emission rates for the $np\gamma$ process. We will consider two different situations using the soft photon approximation to find the respective reaction amplitudes: the case of photon bremsstrahlung through nucleon-nucleus scattering and the case of bremsstrahlung from the thermalized low-temperature medium. It will be shown that the $np\gamma$ reaction is in fact influenced in these two cases in different ways, both qualitatively and quantitatively. Specifically, we will show the possible manifestations of the modification of the two-nucleon scattering states in the nuclear matter on the photon bremsstrahlung emission.[1]

It should be mentioned that in the present work we re-

[1]A discussion of the in-medium effects on the processes of neutrino bremsstrahlung emission from neutron stars matter, processes which are similar to a certain extent to the photon bremsstrahlung discussed here, can be found in Ref. [24].

stricted ourselves to a certain class of medium effects that can be adequately described by the summing ladder diagrams in the particle-particle channel. We will not consider the radiation of photons from the intermediate scattering states. Note that another interesting related effect is the Landau-Pomeranchuk quenching of the bremsstrahlung. A strong in-medium suppression of the bremsstrahlung process due to the coherence effects has been predicted ([35] and reference therein).

The paper is organized as follows: In Sec. II, the expression for the photon production cross section for free nucleon-nucleon collisions is given and its relation to the in-medium $\mathcal{T}$ matrix in the nuclear case is shown. The NN cross section is calculated in the framework of the thermodynamic $\mathcal{T}$-matrix approach in the ladder approximation in Sec. III. The results for the bremsstrahlung cross section are presented in Sec. IV.

II. BREMSSTRAHLUNG OF A REAL PHOTON IN THE SOFT PHOTON APPROXIMATION

As a starting point for the discussion of the bremsstrahlung process, let us recall the case of free nucleon-nucleon bremsstrahlung $N_1N_2 \to N_3N_4\gamma$. By taking into account the relativistic kinematics, the respective cross section reads

$$d\sigma^{NN\gamma}(p_1,p_2,p_3,p_4,q_\gamma)=\frac{|\mathcal{T}^{NN\gamma}|^2}{2(2\pi)^5\sqrt{s(s-4m^2)}}\frac{d^3p_3}{2\varepsilon_3}\frac{d^3p_4}{2\varepsilon_4}\frac{d^3q_\gamma}{2\omega_\gamma}\delta^4(p_1+p_2-p_3-p_4-q_\gamma), \tag{1}$$

where p_i, $\varepsilon_i=\sqrt{p_i^2+m^2}$, and q_γ and ω_γ are the four-momentum and energy for the ith nucleon and photon, respectively, $s=(p_1+p_2)^2$, and $\mathcal{T}^{NN\gamma}$ is the amplitude of the process.

In further calculations of the bremsstrahlung cross sections, we shall use the soft photon approximation (SPA) [25,26], which usually contains the following three approximations: (a) Only the electric part of the electromagnetic current is kept, (b) the photon radiation from the internal charged meson exchange lines and nucleon-nucleon-meson vertices is neglected, and (c) an approximate integration is carried out over unobservable kinematic variables where the momentum, energy, and invariant mass of the virtual photon are assumed to be negligible as compared to other variables (e.g., the initial momentum and momentum transfer, etc.). The comparison with the exact diagrammatic calculation, for example, of the dilepton production [27–30] shows that the first two approximations do not change the final result by more than a few percent and, therefore, are justified. However, the third approximation appears to be crude. To improve the result, a phenomenological reduction factor has been introduced in Ref. [26], which is aimed at reducing the remaining phase-space volume of colliding hadrons in their final state. However, we would like to note that this factor cannot be derived explicitly from a multidimensional integral. Therefore, in the present paper we use a model which employs the first two approximations (a),(b) of the SPA, but takes into consideration the exact kinematic relations.

In the SPA the amplitude of the process can be represented in the form

$$\mathcal{T}^{NN\gamma}=\mathcal{T}^{\gamma}\mathcal{T}^{NN}, \tag{2}$$

where $\mathcal{T}^{NN}$ is the amplitude of the elastic $NN \to NN$ scattering and the electromagnetic matrix element is defined as

$$|\mathcal{T}^{\gamma}|^2=|\epsilon_\mu J^\mu|^2, \tag{3}$$

with the hadron current

$$J^\mu=-Q_a\frac{p_1^\mu}{(p_1\cdot q_\gamma)}-Q_b\frac{p_2^\mu}{(p_2\cdot q_\gamma)}+Q_a\frac{p_3^\mu}{(p_3\cdot q_\gamma)}+Q_b\frac{p_4^\mu}{(p_4\cdot q_\gamma)}\,. \tag{4}$$

Here Q_a and Q_b are the charges of the incoming nucleons and $(p_i\cdot q_\gamma)$ is the scalar product of the four-momenta p_i and q_γ.

The elastic $NN \to NN$ scattering amplitude squared is proportional to the differential elastic cross section [see Eq. (21) below]. To perform exploratory calculations, we further use the total elastic cross section $\sigma^{NN}(s)$ instead of the differential elastic cross section:

$$\frac{d\sigma^{NN}}{d\Omega}(s,\Omega)\simeq\frac{\sigma^{NN}(s)}{4\pi}\,. \tag{5}$$

In this case, having used Eqs. (1), (2), and (5), the expression for the photon production cross section becomes

$$\frac{d\sigma^{NN\gamma}}{d\omega_\gamma}(s,\omega_\gamma)=\frac{\alpha}{8\pi^2}\sqrt{s}\sigma^{NN}(s)\int|\mathcal{T}^{\gamma}|^2\frac{d\varepsilon_3}{|\vec{p}_3|}\,d\cos\theta_\gamma d\varphi_3, \tag{6}$$

where θ_γ is the photon polar angle and φ_3 is the azimuthal angle of the outgoing nucleon in the center-of-mass frame of the initial nucleon. This equation relates the bremsstrahlung cross section directly to the elastic NN cross section.

The application of Eq. (6) to collisions of nuclei includes two modifications. First, σ^{NN} is affected by the nuclear environment. This problem will be considered in Sec. III. Second, nuclear bremsstrahlung cross sections are obtained by integrating the $NN\gamma$ bremsstrahlung cross section (6) over the effective energy distribution $f(E)$ of colliding nucleons. In the general case, this distribution will depend on the reaction considered and should be found by solving appropriate kinetic equations. Such a complete treatment of the problem is beyond our consideration here. To estimate the importance of in-medium effects in nuclear collisions, we shall use certain ansatz for the form of $f(E)$ to simulate conditions specifying (6) for the given reaction.

Let us discuss the case of bremsstrahlung due to a collision of a nucleon of definite momentum with a nucleus in thermal equilibrium. Then, the spectrum of nucleons in the target nucleus can be described by a Fermi distribution characterized by the energy E_2:

$$f(E_2)=\left\{1+\exp\left[\frac{(E_2-\mu)}{T}\right]\right\}^{-1}, \tag{7}$$

where μ is the chemical potential, T is temperature, and E_i is the nonrelativistic single-quasiparticle spectrum. It is related to the relativistic quasiparticle energy ε_i as $E_i \simeq \varepsilon_i - m$.

Thus, for the nucleon-nucleus ($NA\gamma$) bremsstrahlung cross section, we find

$$\frac{d\sigma^{NA\gamma}}{d\omega_\gamma}(\vec{p}_1,\omega_\gamma)=\int \frac{d\sigma^{NN\gamma}}{d\omega_\gamma}(s,\omega_\gamma)f(E_2)\frac{d\vec{p}_2}{(2\pi)^3}, \tag{8}$$

with

$$s=(p_1+p_2)^2=2(m^2+\varepsilon_1\varepsilon_2-|\vec{p}_1||\vec{p}_2|\cos\theta_{12}).$$

Strictly speaking, this situation corresponds to the photon production in the first chance np collision. Note that the corresponding bremsstrahlung cross section is obtained by substituting the elastic in-medium nucleon-nucleon cross section $\sigma^{NN}(s)=\sigma(E_{\mathrm{lab}},K)$ in Eq. (6). It will be a function of the nonrelativistic nucleon energy in the laboratory frame $E_{\mathrm{lab}}=2E=2(\sqrt{s}/2-m)$ and a value of the total momentum K, where $K^2=\vec{p}_1^{\,2}+\vec{p}_2^{\,2}-2|\vec{p}_1||\vec{p}_2|\cos\theta_{12}$.

It is necessary also to take into account the increase in the kinematical limit for ω_γ as compared to the NN interaction:

$$\omega_{\gamma_{\max}}=\frac{s_{\max}-4m^2}{2\sqrt{s_{\max}}}, \tag{9}$$

$$s_{\max}=2(m^2+\varepsilon_1\varepsilon_{2\max}+|\vec{p}_1||\vec{p}_2|_{\max}). \tag{10}$$

In the NN case we would have $\vec{p}_2=0$, $\varepsilon_2=m$, and $s^{NN}_{\max}=2(m^2+\varepsilon_1 m)$, while in the NA case at low temperatures we get the Fermi momentum $p_F=|\vec{p}_2|_{\max}$, $\varepsilon_2=\varepsilon_F$, and $s^{NA}_{\max}=2(m^2+\varepsilon_1\varepsilon_F+|\vec{p}_1|p_F)$.

The nucleon-nucleus scattering case is of special interest if photons are created in the first few np collisions. However, the contribution of photons emitted during the subsequent collisions may be essential when a thermalized dense nuclear matter has been formed. This contribution becomes larger with increasing mass number of colliding nuclei. Thus, after the first few collisions, the formed nuclear fireball may be regarded as a thermalized system whose local thermodynamic equilibrium is characterized by certain temperature T and density n. Among the several possible modifications of two-nucleon scattering in the medium, the effects of the pion mode softening of the nucleon-nucleon interaction was applied to the photon bremsstrahlung problem in Ref. [10]. This effect emerges from the summation of the scattering amplitude in the particle-hole channel, which allows one to take into account medium polarization effects. In contrast, we shall concentrate on the medium effects described in terms of the in-medium $\mathcal{T}$ matrix in the particle-particle interaction channel, which is suitable for treating the short-range correlation and onset of superfluid pairing in the system. A simultaneous treatment of both the long-range polarization effects and the short-range correlations in the $\mathcal{T}$ channel, which requires diagram summations in the parquetlike manner, is not rigorously developed yet.

The only simply treatable case for nucleus-nucleus collisions is that when the formed hot and dense system is close to an equilibrium state. Then the bremsstrahlung cross section can be written as

$$d\sigma^{AA\gamma}(\omega_\gamma)=(2\pi)^3\int\left\{\prod_{i=1}^{4}\frac{d^3p_i}{(2\pi)^3}\right\}\int\frac{d^3q_\gamma}{2\omega_\gamma(2\pi)^3}|\mathcal{T}^{AA\gamma}|^2\delta(\vec{p}_4+\vec{p}_3+\vec{q}_\gamma-\vec{p}_1-\vec{p}_2)\,\delta(E_1+E_2-E_3-E_4-\omega_\gamma)f(E_1)f(E_2)$$
$$\times[1-f(E_3)][1-f_4(E_4)], \tag{11}$$

where the delta functions ensure the momentum-energy conservation laws and $f(E_i)$ is the Fermi distribution function.

The thermalized nuclear matter at low temperature can be regarded here as a highly degenerated Fermi system. Therefore, the momentum of nucleons can be assumed to be restricted to their Fermi momentum, i.e., $|p-p_F|\ll p_F$. This essentially simplifies the evaluation of the phase-space integrals appearing in Eq. (11). In particular, the integration over the energy variables and nucleon angles can be separated in this case. One finds

$$d\sigma^{AA\gamma}(\omega_\gamma)=2\pi I_\Omega I_E\int\frac{d^3q_\gamma}{2\omega_\gamma(2\pi)^3},$$

$$I_\Omega=\frac{(2m)^4}{(2\pi)^9}\int\left\{\prod_{i=1}^{4}d^3p_i\,\delta(p_i^2-p_F^2)\right\}\delta(\vec{p}_4+\vec{p}_3-\vec{p}_1-\vec{p}_2)|\mathcal{T}^\gamma|^2|\mathcal{T}^{NN}|^2, \tag{12}$$

$$I_E=\int\left\{\prod_{i=1}^{4}dE_i\right\}\delta(E_1+E_2-E_3-E_4-\omega_\gamma)f(E_1)f(E_2)[1-f(E_3)][1-f_4(E_4)].$$

Note that because of the SPA the reaction transition matrix element is factorized and photon momentum in the momentum conservation δ function is neglected. After straightforward calculations we get

$$I_\Omega=\frac{4m^4p_F^3}{\pi^3}\int_0^1 dy(1-y^2)\int_0^{\sqrt{1-y^2}}\frac{dx}{\sqrt{1-x^2-y^2}}\mathcal{T}(x,y), \tag{13}$$

$$I_E=\frac{\omega_\gamma}{T}\left[1+4\pi^2\left(\frac{\omega_\gamma}{T}\right)^2\right]g\left(\frac{\omega_\gamma}{T}\right). \tag{14}$$

In Eq. (13), the $\mathcal{T}$ matrix $\mathcal{T}(x,y)$ is a function only of new variables x,y, where $x=q/2p_F$ and $y=K/2p_F$ are the normalized momentum transfer ($\vec{q}=\vec{p}_1-\vec{p}_2$) and total momentum of a nucleon pair and $g(\omega_\gamma/T)=[\exp(\omega_\gamma/T-1)]^{-1}$ is the Planck distribution function for a photon. Using these relations and the explicit expression for the electromagnetic part of the transition probability, the photon emission cross section is reduced to

$$\frac{d^2\sigma^{AA\gamma}}{d\Omega d\omega_\gamma}=\frac{8e^2n}{3}[\nu(\mu)]^2T^2\left[1+\frac{1}{4\pi^2}\left(\frac{\omega_\gamma}{T}\right)^2\right]g\left(\frac{\omega_\gamma}{T}\right)\int_0^1 dy(1-y^2)\int_0^{\sqrt{1-y^2}}\frac{dx}{\sqrt{1-x^2-y^2}}\mathcal{T}(x,y), \tag{15}$$

with the density of the nucleon states at the Fermi surface $\nu(\mu)=mp_F/\pi^2$.

Thus the photon emission cross section depends on the medium temperature through the explicit factor $T^2(1+\omega_\gamma^2/4\pi T^2)g(\omega_\gamma/T)$ and also implicitly through the temperature-dependent $\mathcal{T}$ matrix and effective mass of a nucleon. It is noteworthy that in addition to the direct proportionality of the reaction probability to the nucleus density, n, and to the density of nucleon states $\nu(\mu)$, the medium density is entering into the $\mathcal{T}$ matrix and the nucleon effective mass too.

In order to reveal the medium influence on the $np\rightarrow np\gamma$ reaction, the ratio of the in-medium photon emission cross section to that of the free-space scattering must be considered:

$$R_{AA\gamma}\equiv\left(\frac{d^2\sigma^{AA\gamma}}{d\Omega d\omega_\gamma}\right)\Bigg/\left(\frac{d^2\sigma^{AA\gamma}}{d\Omega d\omega_\gamma}\right)_{\text{free}}=\frac{\int_0^1 dy(1-y^2)\int_0^{\sqrt{1-y^2}}\frac{dx}{\sqrt{1-x^2-y^2}}\mathcal{T}(x,y)}{\int_0^1 dy(1-y^2)\int_0^{\sqrt{1-y^2}}\frac{dx}{\sqrt{1-x^2-y^2}}\mathcal{T}^{\text{free}}(x,y)}. \tag{16}$$

The results of the numerical evaluation of this ratio will be given in Sec. IV followed by the discussion of the in-medium $\mathcal{T}$ matrix in the next section.

To proceed further to the medium modifications of the photon bremsstrahlung process, we will need the in-medium nucleon-nucleon cross. It will be used as an input in the expressions for different cases of nucleon-nucleus and nucleus-nucleus collisions derived above.

III. IN-MEDIUM *NN* CROSS SECTION

To show the physical approximations made in calculations of the photon bremsstrahlung, we shall briefly discuss, following Ref. [14], the derivation of the in-medium scattering cross section by means of the thermodynamic Green function technique. In this approach, the two-nucleon scattering in the medium is described by the thermodynamic $\mathcal{T}$ matrix, which is governed by the corresponding Bethe-Salpeter equation

$$\mathcal{T}(121'2')=K(121'2')+\int d3\, d3'\, d4\, d4'\, \mathcal{K}(1234)G_1(33')\times G_1(44')\mathcal{T}(3'4'1'2'), \tag{17}$$

where 1 denotes the wave number p_1, spin σ_1, and isospin τ_1 of the nucleons. The ladder approximation for the thermodynamic $\mathcal{T}$ matrix is obtained by the substitution of the four-point interaction $\mathcal{K}$ by the bare nucleon-nucleon interaction V in Eq. (17). Within the quasiparticle approximation [22], the product of the two one-particle Greens functions G_1 in (17) results in

$$G_2^0(p_1,p_2,z)=\frac{1-f[E(p_1)]-f[E(p_2)]}{z-E(p_1)-E(p_2)}, \tag{18}$$

where z is the two-particle energy and $E(p_1)$ and $E(p_2)$ are the quasiparticle energies. The last are defined in terms of the self-energy as

$$E(p_1)=\frac{\hbar^2p_1^2}{2m}+v(p_1),$$

with

$$v(p_1)=\mathrm{Re}\Sigma(p_1,\omega)|_{\omega=E(p_1)}, \tag{19}$$

where the self-energy $\Sigma(p_1,\omega)$ was calculated in the $\mathcal{T}$-matrix approximation (for details, see Ref. [13]).

Within these approximations two medium effects are included in the quantity G_2^0. One of them is the phase-space occupancy Q (Pauli blocking) of the surrounding nucleons given by $Q(p_1,p_2)=1-f[E(p_1)]-f[E(p_2)]$ in Eq. (18). This form of the Pauli operator takes into account hole-hole scattering, which is neglected in the usual Brueckner theory where the Pauli operator has the factorized form $Q_B(p_1,p_2)=\{1-f[E(p_1)]\}\{1-f[E(p_2)]\}$ [22]. The second medium contribution is coming from the renormalization of the quasiparticle energies (19) entering into Eq. (18).

In our calculation of in-medium cross sections, the Paris potential is used for a bare nucleon-nucleon interaction. The Paris potential is derived from meson theory and gives a quantitatively reliable description of the on- and off-shell properties of the nucleon-nucleon interaction in the vacuum [31]. In particular, the nucleon-nucleon scattering phase shifts, which are well known from experiment, are reproduced with high accuracy. It was successfully applied to the calculation of the equilibrium properties of the nuclear matter as well (see Ref. [32]).

Using the partial wave decomposition, which becomes possible after angle averaging of the Pauli operator and the quasiparticle energies, the total thermodynamic $\mathcal{T}$ matrix may be constructed from the partial $\mathcal{T}$ matrices. In the Matsubara-Fourier representation, the partial $\mathcal{T}$ matrix for the channel $\alpha=(S,L,J)$ reads

$$\mathcal{T}_\alpha^{LL'}(p,p',K,z)=V_\alpha^{LL'}(p,p')+\sum_{p''L''} V_\alpha^{LL''}(p,p'')G_2^0(p'',K,z)\mathcal{T}_\alpha^{L''L'}(p'',p',K,z), \tag{20}$$

where the relative momentum $\vec{p}=(\vec{p}_1-\vec{p}_2)/2$ is introduced. In evaluating the quasiparticle energies, a p-dependent effective two-particle effective mass $m_{12}^*(p,K)$ is introduced. This allows one to define an effective chemical potential $\mu_{\rm rel}=\mu-\Delta\epsilon$ with respect to the continuum boundary which includes the averaged single-particle self-energy shift $\Delta\epsilon$.

The in-medium differential cross section for an unpolarized system is defined via the on-shell $\mathcal{T}$ matrix ($|\vec{p}|=|\vec{p}|=p$) as

$$\frac{d\sigma}{d\Omega}(p)=\frac{N(p)^2}{(2s_1+1)(2s_2+1)}\sum_{S,M_S,S',M_S'}\frac{(2\pi)^4}{p^2}|\mathcal{T}(\vec{p}SM_S,\vec{p}'S'M_S')|^2, \tag{21}$$

with the generalized density of states,

$$N(p)=\frac{pm_{12}^*(p,K)}{2(2\pi)^3\hbar^2}\,. \tag{22}$$

The in-medium differential cross section, Eq. (21), depends on the temperature T and chemical potential μ of the medium as well as on the total momentum K of a pair of nucleons. (The evident notation of some variables is skipped.) In the low-density limit $\mu/T\to-\infty$, the thermodynamic $\mathcal{T}$ matrix, Eq. (20), approaches the scattering $\mathcal{T}$ matrix describing the isolated elastic NN scattering. In this limit, the in-medium differential cross section, Eq. (21), approaches to the free NN differential cross section. By integrating Eq. (21) over the angles, one arrives at the total cross section in the medium:

$$\sigma(p)=\sum_{J,L,L'}\frac{(2J+1)2\pi^3N(p)^2}{(2s_1+1)(2s_2+1)p^2}|\mathcal{T}_\alpha^{LL'}(p,p)|^2. \tag{23}$$

Note that Eq. (23) gives the NN cross section with taking Pauli blocking into account in the intermediate states only, i.e., without correction for Pauli blocking in the outgoing channel [17]. Numerical results will be given in the next section.

IV. RESULTS AND DISCUSSION

In order to evaluate the photon bremsstrahlung cross section in the soft photon approximation, we first discuss the in-medium nucleon-nucleon cross section.

For the numerical evaluation of $\mathcal{T}_\alpha^{LL'}$ and the cross section, Eqs. (21) and (23), we use a separable approximation of the Paris nucleon-nucleon potential [33] (see [13] for details). The features of the Paris interaction mentioned above, especially the reliable description of the empirical two-nucleon scattering data, are preserved in the separable approximation by Haidenbauer and Plessas [33]. This separable approximation was applied in nuclear matter calculations too [13,32].

As has been shown in the previous section, the in-medium total nucleon-nucleon cross section depends on the relative energy $E=\hbar^2p^2/m_{12}^*(p,K)$ (or the energy in the laboratory frame, $E_{\rm lab}=2E$), on the total momentum K of the pair (via

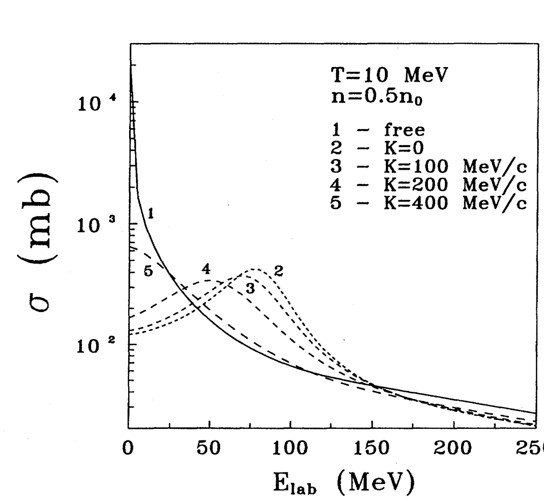

FIG. 1. In-medium pn cross section σ as a function of $E_{\rm lab}$ at density $n=0.5n_0$ and temperature $T=10$ MeV for several values of the total momentum K. The solid line gives the free cross section.

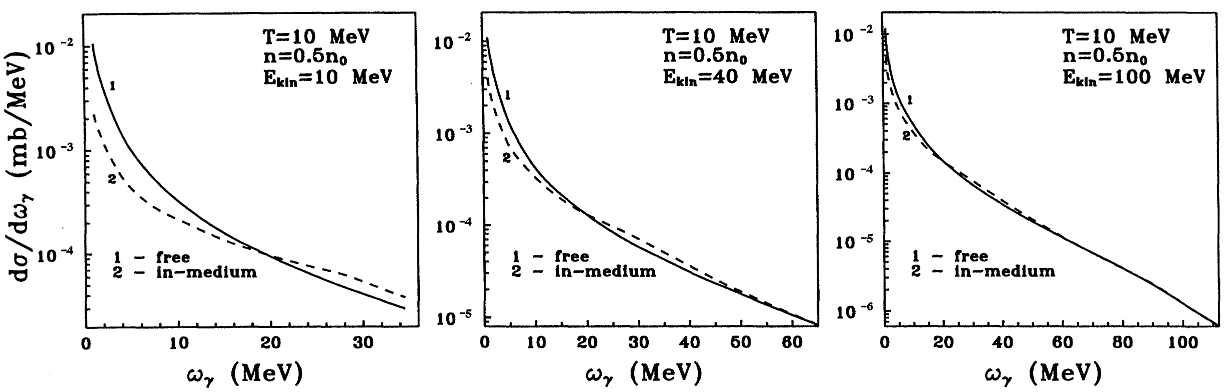

FIG. 2. *pn* bremsstrahlung cross section for the nucleon-nucleus collisions at 10, 40, and 100 MeV energies of a bombarding nucleon. Solid lines (curves 1) correspond to the calculation of the bremsstrahlung cross section with the total free cross section. The dashes lines (curves 2) correspond to the calculation with in-medium total nucleon-nucleon cross section at temperature $T=10$ MeV and density $n=0.5n_0$.

the Pauli blocking and the self-energy shifts), and also on the thermodynamic variables characterizing the medium, namely, density n and temperature T. Figure 1 shows the energy dependence of the average in-medium cross section σ_{np} at a fixed density $n=0.5n_0$ (where $n_0=0.17$ fm^{-3} is the saturation density) and temperature $T=10$ MeV for several values of the total momentum K of the pair. For comparison, the free total NN cross section is also plotted (solid curve 1), which is calculated using Eq. (23) and neglecting all medium effects. Taking into account partial waves up to $L=2$, the experimental free NN cross section is reproduced with a high accuracy.

At finite density, the strongest deviations from the free cross section occur in the zero total momentum ($K=0$, curve 2). While at low energies $E_{\mathrm{lab}}<50$ MeV, the in-medium cross section is strongly suppressed as compared to the free cross section, it is enhanced in the energy range 50 MeV $<E_{\mathrm{lab}}<150$ MeV and then again suppressed at still higher energies. Whereas the strong suppression at very low energies was expected, since the available phase space is reduced as a result of Pauli blocking, the enhancement at higher energies is less obvious. The maximum at $E_{\mathrm{lab}}\approx 90$ MeV [$E\simeq 2\mu_{\mathrm{rel}}$, with μ_{rel} being the effective chemical potential, as defined after Eq. (20)] can be understood as a precursor effect of the superfluid phase transition in the symmetrical nuclear matter, which shows up when the critical temperature is approached for this phase transition from above (see, for details, [14]). At still higher energies the self-energy effects (effective mass) are dominating. In this limit the T-matrix approximation goes over to the Born approximation and leads to a suppression which is proportional to the square of the effective mass [13].

When the total momentum K of a pair increases, the Pauli blocking is reduced. Because of this effect, the suppression at low energies as well as the enhancement at energies $E\simeq 2\mu_{\mathrm{rel}}$ becomes less pronounced. For the value of $K\approx 400$ MeV/c, the in-medium cross section is already quite close to the free cross section except for very low energies. Thus, at finite density and temperature, considered here, the NN cross section considerably deviates from its free space value, exhibiting a nontrivial behavior around the energy $E_{\mathrm{lab}}\approx 90$ MeV.

The results presented in Figs. 2 and 3 show the in-medium influence on photon spectra from nucleon-nucleus collisions. It can be seen that at moderate temperature $T\approx 10$ MeV a low-energy domain in the photon spectrum is reduced especially at the low nucleon energy, while an increase of the bremsstrahlung cross section is observed for hard photons ($\omega_\gamma>\mu_{\mathrm{rel}}$). At sufficiently high energies and density $n=0.5n_0$ (Fig. 2), the difference between in-medium and free space bremsstrahlung cross sections vanishes, because the NN scattering cross section tends to the free scattering one with increasing total momenta K (see Fig. 1). Note that the bremsstrahlung photon emission from internal (as a result of charge exchange currents) is neglected in our consideration, but can contribute to photon production when the collision energy increases.

The influence of the nuclear matter on the photon emission rate increases when we proceed to the regime of higher densities and smaller temperatures: The in-medium result deviates from the vacuum scattering by a factor of 5–10 at density $n=n_0$ and temperature $T=5$ MeV (Fig. 3) and only by a factor of 2–3 (in the low-energy domain) at $n=0.5n_0$ and $T=10$ MeV (Fig. 2). This result gives a hint that the discussed effect seems to be manifested more clearly in central proton-nucleus collisions rather than at the impinging stage of heavy-ion collisions where conditions may be somewhat close to that corresponding to Fig. 2.

In the case of a nucleus-nucleon collision, we have derived the ratio $R_{AA\gamma}$ [see Eq. (16)], which shows the relevance of the in-medium effects at low temperatures. In Fig. 4, this quantity is displayed as a function of the nuclear matter density for three different temperatures. We find that, generally, the emission probability is essentially suppressed as compared to the free scattering case. For high-density region $n>n_0=0.17$ fm^3, it amounts to $\sim 15\%$ and is weakly dependent on temperature. At low densities the temperature dependence becomes more pronounced, showing an increase in the $R_{AA\gamma}$ ratio with decreasing temperature. This can be understood by examining the effective in-medium cross sections.

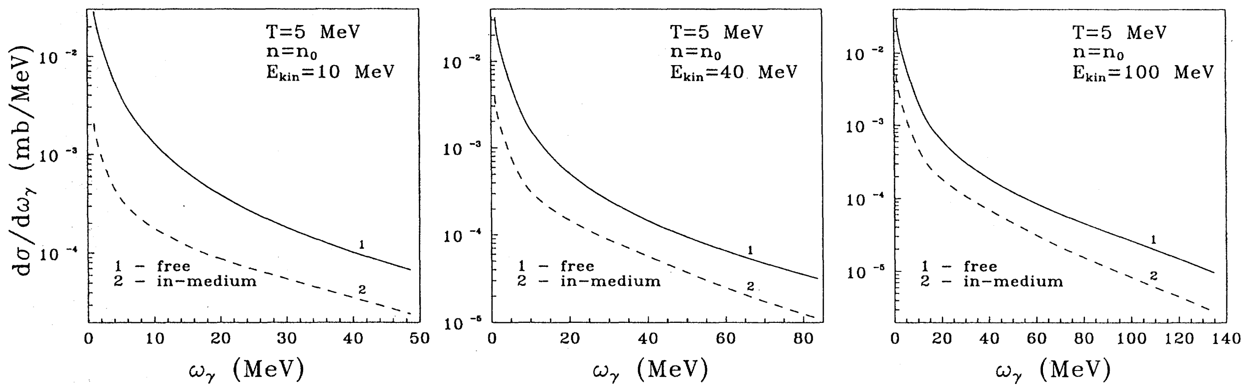

FIG. 3. Same as in Fig. 2 for density $n = n_0$ and temperature $T = 5$ MeV.

Note that at the temperatures of interest the in-medium cross section is strongly suppressed for collisions at small relative energies, and consequently $R_{AA\gamma}$ is smaller than unity. This effect results from the suppression of the available phase for intermediate scattering states by the Pauli-blocking operator $Q = \langle 1 - f - f \rangle$ in the two-nucleon propagator, Eq. (18). However, for lower temperatures, an increase in the effective cross section is observed at relative energies around $E_{\mathrm{rel}} = 2\mu_{\mathrm{rel}}$, associated with the neutron-proton pairing attraction in the ${}^1S_3$-${}^1D_3$ interaction channel. It is compensated partially by the reduction of the cross section at small relative energies, and therefore the ratio $R_{AA\gamma}$ is enhanced. In the high-density region, this compensation is not operative because the effective chemical potential is shifted to higher values at which the attractive pairing interaction becomes gradually weaker. The real picture is more complicated since the ratio $R_{AA\gamma}$ involves a two-dimensional integral over the scattering $\mathscr{T}$ matrix as a function of the relative and total momenta and certain weight functions; however, the numerical evaluation confirms the qualitative picture given above.

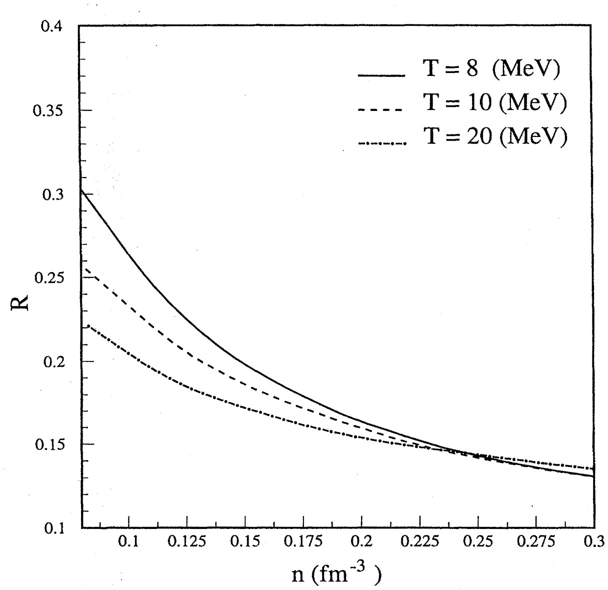

FIG. 4. Ratio $R_{AA\gamma}$ for the photon bremsstrahlung in the case of nucleus-nucleus collisions as a function of the density. With decreasing temperature this ratio increases as a result of the nontrivial behavior of the in-medium NN cross section at low temperatures (see text).

V. CONCLUSIONS

A comparison of the nucleon-nucleus and nucleus-nucleus cases shows that in nucleus-nucleus collisions the medium effects are much more pronounced for the later time moments when complete thermalization is achieved in the system. As has been shown in Ref. [2], the bremsstrahlung radiation in heavy-ion collisions at energies <100 MeV/nucleon is coming mainly from two sources separated in time: from the initial stage of the np collisions and from the thermalized zone at the nucleon density close to the freeze-out density. This picture is in agreement with the recent photon-photon interferometry measurements [6]. As follows from our consideration, the bremsstrahlung photons from these two sources will be suppressed in a different manner. It is desirable to check this prediction for the bremsstrahlung suppression in the dynamical transport model with in-medium modifications of NN cross sections incorporated in the low-temperature limit, similarly to the approach at higher temperatures given in Ref. [34].

In conclusion, let us summarize the results.

(i) Expressions for the differential cross section of photon emission from the nucleon-nucleon and nucleon-nucleus bremsstrahlung processes at low temperatures in terms of the in-medium $\mathscr{T}$ matrix are derived. Important in-medium modifications arise as a result of the Pauli blocking of intermediate scattering states and quasiparticle energy shifts due to the mean field interaction in the in-medium $\mathscr{T}$ matrix.

(ii) In contrast to the reduction found in a simple models which incorporate only the nucleon effective mass, our consideration in the T-matrix approach shows a more involved behavior. It is found that the photon emission rate for nucleon-nucleus scattering can be enhanced for photon energies lower than the chemical potential of the medium (i.e., $\omega_\gamma < 20$ MeV for $n = 0.5n_0$, $T = 10$ MeV) and suppressed at larger energies in the relatively hot and dilute medium.

(iii) Using the decoupling procedure for the energy and momentum integration variables, valid for the degenerate thermalized medium, the phase-space integral is evaluated for the nucleus-nucleus scattering case. The ratio of the in-medium modified emission rate to that of the free space is calculated and a reduction by about 15% is found. Thus the largest modification of in-medium *np* bremsstrahlung cross sections is expected to occur in the case of central proton-nucleus collisions and for heavy-ion collisions when photons are emitted from the thermalized zone at comparatively low nucleon density.

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