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Photon production from baryon rich quark-gluon plasma

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Abstract. We study high energy photon production from a quark-gluon plasma at finite baryon density. We find that the photon production spectrum from the quark-gluon plasma maintained at constant temperature is only mildly dependent on the quark chemical potential.

Keywords. Phase transition; quark-gluon plasma; photon production; heavy ion collisions.

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1. Introduction

The ongoing and planned experiments involving ultrarelativistic heavy ion collisions are believed to unravel the exciting possibility of a transition from normal nuclear matter to a deconfined phase of quarks and gluons (QGP). Among the various signals of QGP formation are the electromagnetic 'probes', both virtual and real photons, which are particularly suitable to carry information of the QGP state since they are able to leave the reaction volume with little interaction and hence act as ideal markers of the evolutionary history of the QGP state. The production rates of dileptons (from virtual photons) and of real high-energy photons from QGP have been calculated in the literature [1, 2] and have been compared with their production rates from the hot hadronic matter in order to arrive at the definitive conclusion regarding the formation of QGP. A detailed treatment of the photon production at $T = 200\,\mathrm{MeV}$ led Kapusta, Lichard and Seibert [3] to conclude that 'the hadron gas shines just as brightly as the quark-gluon plasma' and that high momentum photons may not be able to distinguish between a quark-gluon plasma and hadronic gas but may provide an excellent thermometer for ultrarelativistic heavy ion collisions. It was pointed out by Sinha and collaborators [4] that this comparison has been made by assuming the two sources to be static at the same temperature. When account was taken of the space-time evolution in a relativistic hydrodynamical model over the entire evolutionary history, it was found to provide a window in the high P_T region ($P_T >$2–3 GeV) where the contribution from QGP outshines that from the hadronic sector.

The dilepton and photon emission rates discussed above have been calculated for baryon free plasma, by neglecting the baryon chemical potential μ_{B} of the plasma and depend only on its temperature. However experiments at Brookhaven's AGS with energies of the order of 10–15 GeV per nucleon have indicated nearly complete stopping of the nuclei during collisions and sizeable stopping at CERN's existing SPS [5]. This is also confirmed in numerical simulations of nuclear collisions employing relativistic quantum molecular dynamics [6] and a relativistic cascade model [7]. Not only this, even at RHIC and at CERN's anticipated LHC, recent calculations [8] hint that even for energies $\sim$ 100–200 GeV per nucleon, the colliding heavy ions may not be fully transparent and the baryonic chemical potential may not be small compared to the temperature, at least in the experimentally accessible region of rapidity $|y| > 1$, as the proposed experiments are unable to measure the electromagnetic probes in the almost baryon free mid rapidity ($y = 0$) region.

Thus with the possibility that baryon density may be large in the QGP, the thermodynamic equilibrium is a function of both the temperature and the chemical potential μ_B. This would lead to modification of the dilepton and photon emission rates. Recent calculations by Dumitru *et al* [9] indeed demonstrated such an effect. They found that the dilepton producton rate is a strongly decreasing function of μ_B. For a value of μ/T estimated from heavy ion collisions lying between 1 and 2, they found the dilepton production suppressed by a factor 3–10, depending on the value of μ/T, the dilepton mass and momentum, the plasma energy density etc. In this context it is important to estimate the effect of finite chemical potential on the production of direct photons coming from the $q\bar{q}$ annihilation process $q\bar{q} \to g\gamma$ and from the QCD Compton processes $qg \to q\gamma, \bar{q}g \to \bar{q}\gamma$. Direct photon production in the hadronic sector coming mainly from the pseudo scalar and vector meson decays however, will not be affected by a finite value of the chemical potential.

2. Photon production rate

Consider the annihilation process

$$q(p_1) + \bar{q}(p_2) \to g(p_3) + \gamma(p) \tag{1}$$

and the Compton processes

$$q(p_1)(\bar{q}(p_1)) + g(p_2) \to q(p_3)(\bar{q}(p_3)) + \gamma(p) \tag{2}$$

and let $\sum |M_i|^2$ be the matrix element squared and summed over final spins and colour degrees of freedom. The photon production rate per unit volume is then given by

$$\begin{aligned} R_i =& N_i \int \prod_{i=1}^{3} \frac{\mathrm{d}^3 p_i}{(2\pi)^3 2E_i} f_1(E_1) f_2(E_2)(1 \pm f_3(E_3))(2\pi)^4 \\ & \times \delta^4(p_1 + p_2 - p_3 - p) \sum |M_i|^2 \frac{\mathrm{d}^3 p}{(2\pi)^3 2E}, \end{aligned} \tag{3}$$

where the f_is are the appropriate Fermi–Dirac or Bose–Einstein distribution functions and $(1 \pm f_3(E_3))$ accounts for the Pauli blocking or Bose enhancement in the final state. N_i is the overall degeneracy factor and depends upon the electric and colour charge of initial state particles (quarks and gluons) in the specific reaction. The rate can be considerably simplified by noting that for relativisitic particles the final state particles namely, photon and gluon for annihilation and photon and quark (antiquark) for Compton process come out with approximately the same energy and momentum as the incoming particles.

Defining $s = (p_1 + p_2)^2$ and $t = (p_1 - p)^2$ and converting the total rate into a differential rate using the δ function, all the integrals can be done except those over s and t and we get

$$E\frac{\mathrm{d}R_i}{\mathrm{d}^3p} = \frac{N_i}{2(2\pi)^6}\frac{T}{32E}\left[\frac{1}{\mathrm{e}^{(E-\mu)/T}}\frac{1}{1+\mathrm{e}^{\mu/T}} \times \int\frac{\mathrm{d}s}{s}\{\ln(1+\mathrm{e}^{-((s/4E)+\mu)/T}) - \ln(1-\mathrm{e}^{-s/4ET})\} \times \int \mathrm{d}t\sum|M_i|^2 + \mu \leftrightarrow -\mu\right]. \tag{4}$$

For relativistic particles the matrix element square is related to the differential cross-section by [3]

$$\frac{\mathrm{d}\sigma}{\mathrm{d}t} = \frac{\Sigma|M|^2}{16\pi s^2}. \tag{5}$$

For the annihilation process, in this limit

$$\frac{\mathrm{d}\sigma}{\mathrm{d}t} = \frac{8\pi\alpha\alpha_c}{9s^2}\frac{u^2+t^2}{ut} \tag{6}$$

and $N = 20$ when we sum over u and d quarks.

For the two Compton processes

$$\frac{\mathrm{d}\sigma}{\mathrm{d}t} = \frac{-\pi\alpha\alpha_c}{3s^2}\frac{u^2+s^2}{us} \tag{7}$$

and $N = 320/3$.

The integral over t gives the total cross-section and if the quark masses are neglected, the cross-section would diverge. A non-vanishing quark mass regulates the divergence. Although the bare quark mass is very small, the effective mass of the quark in an interacting thermal quark-gluon plasma is considerably modified due to the many body effects of the medium at high temperatures and densities. Braaten and Pisarski [10] found that a quark acquires an effective thermal mass $m_{\mathrm{th}}^2 = 2\pi\alpha_c/3(T^2 + \mu^2/\pi^2)$ and Kapusta *et al* [3] showed that the net effect of the medium on photon production is to replace m^2 by $k_c^2 = 2m_{\mathrm{th}}^2$. We would thus replace m^2 by k_c^2.

Putting all the factors together and performing the integrals over t and s we get

$$E\frac{\mathrm{d}R_{\mathrm{ann.}}}{\mathrm{d}^3p} = \frac{5}{9}\frac{T^2}{\pi^4}\alpha\alpha_c\left[\frac{\pi^2}{6}\left\{\ln\left(\frac{4ET}{k_c^2}\right)+C_B\right\}\left\{\frac{1}{\mathrm{e}^{\mu/T}+1}\frac{1}{\mathrm{e}^{(E-\mu)/T}+1}+(\mu\to-\mu)\right\} + \frac{1}{1+\mathrm{e}^{\mu/T}}\frac{1}{\mathrm{e}^{(E-\mu)/T}+1}\sum_{n=1}^{\infty}(-1)^{n+1}\frac{\mathrm{e}^{-n\mu/T}}{n^2} \times \left(\ln\left(\frac{4ET}{k_c^2}\right) - C_{\mathrm{Euler}} - 1 - \ln n\right) + \frac{1}{1+\mathrm{e}^{-\mu/T}}\frac{1}{\mathrm{e}^{(E+\mu)/T}+1}I_a\right], \tag{8}$$

where

$$I_a = \int_{k_c^2/ET}^{\infty} \ln(1 + \mathrm{e}^{-x+\mu/T}) \left\{ \ln\left(\frac{4ETx}{k_c^2}\right) - 1 \right\} \mathrm{d}x \text{ and } C_B = -2.1472, \tag{9}$$

$$\begin{aligned} E\frac{\mathrm{d}R_{\text{comp.}}}{\mathrm{d}^3p} = \frac{5}{9}\frac{T^2}{\pi^4}\alpha\alpha_c \Bigg[& \frac{\pi^2}{6}\left\{ \ln\left(\frac{4ET}{k_c^2}\right) + C_B' \right\} \\ & \times \left\{ \frac{1}{\mathrm{e}^{-\mu/T}+1} \frac{1}{\mathrm{e}^{(E+\mu)/T}+1} + (\mu \to -\mu) \right\} \\ & + \frac{1}{1+\mathrm{e}^{-\mu/T}} \frac{1}{\mathrm{e}^{(E+\mu)/T}+1} \sum_{n=1}^{\infty} \frac{(-1)^{n+1}\mathrm{e}^{-n\mu/T}}{n^2} \\ & \times \left(\ln\left(\frac{4ET}{k_c^2}\right) - C_{\text{Euler}} + \frac{1}{2} - \ln n \right) \\ & + \frac{1}{1+\mathrm{e}^{\mu/T}} \frac{1}{\mathrm{e}^{(E-\mu)/T}+1} I_c \Bigg], \end{aligned} \tag{10}$$

where

$$I_c = \int_{k_c^2/4ET}^{\infty} \ln(1 + \mathrm{e}^{-x+\mu/T}) \left\{ \ln\left(\frac{4ETx}{k_c^2}\right) + \frac{1}{2} \right\} \mathrm{d}x,$$

$$C_B' = -C_{\text{Euler}} + \frac{1}{2} - \frac{6}{\pi^2} \sum \frac{\ln n}{n^2} = -0.6472. \tag{11}$$

For baryon free plasma $\mu = 0$ and the integrals I_a and I_c can be performed analytically and we recover the well-known results [3] namely,

$$E\frac{\mathrm{d}R_{\text{comp.}}}{\mathrm{d}^3p} = \frac{5}{9}\frac{\alpha\alpha_c}{6\pi^2} T^2 \mathrm{e}^{-E/T} \left\{ \ln\left(\frac{4ET}{k_c^2}\right) + C_F \right\}, \tag{12}$$

$$E\frac{\mathrm{d}R_{\text{ann.}}}{\mathrm{d}^3p} = \frac{5}{9}\frac{\alpha\alpha_c}{3\pi^2} T^2 \mathrm{e}^{-E/T} \left\{ \ln\left(\frac{4ET}{k_c^2}\right) + C_B \right\}. \tag{13}$$

3. Results

We have calculated the effect of finite chemical potential of the quarks on the thermal production rate of photons arising from the quark-gluon plasma. In table 1 we have displayed the rate as a function of μ/T for three values of photon energy at two different quark-gluon temperatures 200 and 150 MeV respectively. We find from the table that the photon production rate is very mildly dependent on the quark chemical potential. Since the direct photon production in the hadronic sector is unaffected by a finite value of chemical potential, the quark-gluon plasma and the hadron gas maintained at a constant temperature produce the high energy thermal photons at roughly the same rate even for the baryon rich matter. If however the energy density of the plasma created in the collisions is known, one needs to employ the equation of state governing the plasma to relate the temperature and chemical potential [11]. Using the phenomenological bag

Table 1. Photon thermal production rate $E(dR/d^3p)$ in $fm^{-4}\,GeV^{-2}$ as a function of μ/T for $E = 1, 2$ and 3 GeV at two different quark-gluon plasma temperatures $T = 200$ MeV and 150 MeV respectively.

	$T = 200$ MeV			$T = 150$ MeV		
Photon energy E μ/T	1 GeV Rate $\times 10^5$	2 GeV Rate $\times 10^7$	3 GeV Rate $\times 10^9$	1 GeV Rate $\times 10^6$	2 GeV Rate $\times 10^9$	3 GeV Rate $\times 10^{12}$
0.0	1.734	1.817	1.478	2.270	4.182	6.286
0.5	1.631	1.732	1.414	2.151	3.998	6.026
1.0	1.585	1.707	1.398	2.109	3.95	5.967
1.5	1.575	1.725	1.416	2.121	3.998	6.051
2.0	1.591	1.785	1.468	2.181	4.141	6.276
2.5	1.624	1.882	1.55	2.282	4.372	6.631

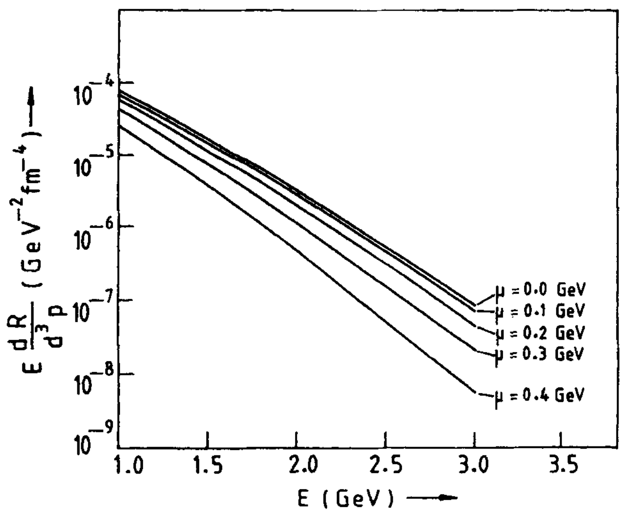

Figure 1. Thermal photon production spectrum $E\,dR/d^3p$ in $GeV^{-2}fm^{-4}$ for a fixed energy density of the plasma ($\varepsilon = GeV\,fm^{-3}$) for various values of μ.

model equation of state

$$\varepsilon = \left(\frac{37\pi^2}{30} - \frac{11\pi\alpha_s}{3}\right)T^4 + 3\left(1 - \frac{2\alpha_s}{\pi}\right)T^2\mu^2 + \frac{3}{2\pi^2}\left(1 - \frac{2\alpha_s}{\pi}\right)\mu^4 + B, \quad (14)$$

we have plotted in figure 1 the photon thermal production rate $E\,dR/d^3p$ as a function of photon energy at a fixed energy density ($\varepsilon = 5\,GeV/fm^3$) for different values of μ. We now find the photon spectrum to be strongly dependent on μ and the production rate to be highly suppressed as in ref. [11] essentially because of a decrease in the temperature of a fixed energy baryon rich plasma in comparison to its baryon free counterpart. In conclusion, if the quark-gluon plasma is formed in a collision and expands and hadronizes in a first order phase transition, the two phases maintain thermal equilibrium and the high energy photons can still be considered as good 'thermometer' for the hot dense matter created in high energy heavy-ion collision experiment at SPS and RHIC.

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Note

After the completion of the paper our attention was drawn to a recent work by Traxler *et al* (*Phys. Lett.* **B346**, 329 (1995)) who used the Bratten–Pisarski resummation method to the baryon rich plasma and employed numerical methods to evaluate the four dimensional integrals. Our results agree with them.

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