BL-TP 94/48
IFT UWr/94
September 1994

DENSITY MODIFICATION OF DILEPTON PRODUCTION IN HOT HADRONIC MATTER*

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INTRODUCTION

It is well known that the thermodynamic properties of highly excited high density hadronic matter are qualitatively different from a dilute system. In particular, the results obtained from studies of the thermodynamics of the Statistical Bootstrap Model [1,2] have shown that there should exist a limiting temperature $0.13 < T_H < 0.18 \text{ GeV}$ above which hadronic matter, with hadrons as a fundamental constituents, does not anymore have a meaning. When approaching the limiting temperature the partition function was found to exhibit a singular structure suggesting the possibility of a phase transition. Indeed, recent developments in QCD and in particular in Lattice-Gauge Theory (LGT) have shown that there exists a phase transition from a confined hadron-phase to a deconfined quark-gluon plasma phase. The value of the phase transition temperature T_c found in LGT is compatible with the value of the Hagedorn temperature T_H [3].

One of the main objectives of ultrarelativistic heavy ion collisions is to study the properties and the nature of high density hadronic matter. It is expected, that for sufficiently large collision energies between heavy ions hadronic matter can be produced in a deconfined quark-gluon plasma phase [4]. The prime objective of the theoretical studies is to find the physical observables sensitive to a phase transition. It has been argued that dilepton production is one of the most promising signals for this [5].

In this work we will study the structure and the properties of low invariant mass dilepton production. Low mass dileptons can have both thermal and non-thermal origin. Non-thermal dilepton pairs are produced mainly by Dalitz decay of neutral mesons [6]. Thermal dileptons are mainly originating from virtual photon bremsstrahlung by pions and quarks. In a dense medium however, thermal production can be modified by various collective effects like dynamical or Debye screening or by multiple scattering. The last effect is known to be very important when discussing particle

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* This work is dedicated to Rolf Hagedorn's 75th birthday.

bremssstrahlung in a medium [7-10]. It is because it can lead to a substantial suppression of bremssstrahlung spectrum due to the Landau-Pomeranchuk effect (LPE) [8-10]. In this work we will discuss the origin and the influence of this effect on the production of dilepton pairs in hot hadronic matter. The cross-sections for real and virtual photon bremssstrahlung in a pion gas and in a quark-gluon plasma will be derived and analyzed. Finally, assuming initial conditions appropriate for U-U collisions at LHC energy, the bremssstrahlung contributions to the overall rate of low mass dilepton pairs will be calculated and compared with the non-thermal background.

MULTIPLE SCATTERING EFFECT AND COHERENCE LENGTH

Many electromagnetic processes, such as e.g. real or virtual photon emission, which occur in the interaction of fast charged particle in a medium, develop in a large region of space along the particle trajectory [7]. The length of this region is called the *coherence length* l_c . The concept of coherence length can be introduced on both the classical and quantum levels. For an illustrative purpose we shall interpret l_c as the length in which the *stripping* of a photon from the radiating charged particle occurs [11]. Suppose that a charged particle colliding with a nucleus radiates a photon of frequency ω . For an ultrarelativistic particle the typical radiation angle θ_B is very small and proportional to the ratio of the particle's mass to its energy E , that is $\theta_B \sim m/E$. Thus, since the radiation is almost parallel to the particle momentum it cannot separate immediately after a collision. The particle and the photon can be considered as independent objects only if they are separated by a distance of the order of the wave length of the radiation $\lambda \sim 1/\omega$. This separation requires some time (radiation formation time) $\Delta t = \lambda/v$, with $v_r = 1 - v \cos \theta_B$ being the relative velocity of particle and photon. The coherence length l_c is defined as the distance traveled by the charged particle in time Δt , that is:

$$l_c = v \Delta t \sim \frac{2E^2}{m^2 \omega} \quad (1)$$

where we have assumed that $\theta_B < 1$ and also replaced the quantity $(1 - v)$ by $m^2/2E^2$.

The coherence length increases rapidly with increasing particle energy and with decreasing frequency of the radiation. Thus for very energetic particles and soft quanta, l_c can reach macroscopic dimensions. This simple analysis shows that particle bremssstrahlung is not really a point-like interaction, but appears over a large region of space in the longitudinal direction of particles motion [11]. Thus when discussing bremssstrahlung due to the interactions of a charged particle in a medium, one necessarily needs to consider two limiting cases:

(i) the coherence length is smaller than the average distance between medium constituents: in this case bremssstrahlung in a medium is undisturbed and equivalent to the radiation in N independent collisions.

(ii) the coherence length is large in comparison to the mean free path or to the average distance between the constituents in the medium: here it is necessary to take into account the interaction of the incident charged particle with the medium within the coherence length. Multiple scattering of the incident particle can disturb particle radiation, causing the standard vacuum formula for particle bremssstrahlung to fail. It is mostly due to the interference of the waves radiated in different collisions within the coherence length.

The result for the coherence length introduced in eq.(1) is very general and does not depend on the kind of medium in which the radiation occurs. Thus l_c itself does not determine yet the dynamics of the radiation. In particular, the knowledge of the

value of l_c is not sufficient to deduce whether interference within the coherence length leads to enhancement or suppression of particle bremssstrahlung in a medium. This is entirely determined by the medium's properties.

Landau and Pomeranchuk [8] have shown that in amorphous material the interference effect within the coherence length is destructive and consequently leads to a substantial decrease of bremssstrahlung. The above result is known in the literature as the Landau-Pomeranchuk effect. The LPE is not only related with electromagnetic radiation, but it was found to be important in a variety of problems, also those which are entirely non-electromagnetic. Recently LPE was discussed in the context of nuclear shadowing effect [12], gluon bremssstrahlung [13] and in the parton cascade model in order to account for the time delay of parton emissions [14]. In the following section we will discuss the influence of LPE on real and virtual photon bremssstrahlung in a hot QCD medium.

QUANTITATIVE DESCRIPTION OF BREMSSTRAHLUNG

In order to make quantitative the description of particle bremssstrahlung in a medium and discuss the role of multiple scattering leading to LPE we will start with a semi-classical approach which is valid in the soft photon approximation, i.e. when the frequency of the radiating photon ω is much smaller than the energy of the incident charged particle. In this approach the energy radiated per unit of momentum by the charged particle moving along the trajectory $\vec{r}(t)$ with velocity $\vec{v}(t)$ is given by the following equation [15]:

$$\frac{dI^\gamma}{d^3k} = \frac{\alpha}{4\pi^2} \left| \int dt \exp(i[\omega t - \vec{k}\vec{r}(t)]) [\vec{n} \times \vec{v}] \right|^2 \quad (2)$$

where $\vec{n} \equiv \vec{k}/|\vec{k}|$ is a unit vector in the direction of photon emission.

In the simplest case of scattering of a charged particle by static centers, the particle trajectory is a broken line. The semi-classical formula is known in this case to lead to the same result as the quantum description upon application of the soft photon theorem. In a medium however, the particle trajectory is quite complicated and due to its continuous interactions it is never on-shell. Thus when discussing bremssstrahlung in a medium it is not clear what set of Feynman diagrams corresponds to a given trajectory and what weight should be given to it. The advantage of the semi-classical approach over the quantum treatment is that it can easily be applied to radiation in a thermal medium. For a series of collisions one has to change first the time integral in eq.(2) to a sum of integrals corresponding to all the pieces of the trajectory. Assuming that the particle velocity is constant between two collisions and performing a time integration, the energy radiated by the charged particle undergoing N scatterings is given by:

$$\frac{dI^\gamma}{d^3k} = \frac{\alpha}{4\pi^2} \left| \sum_{j=1}^N \frac{\vec{n} \times \vec{v}_j}{\omega - \vec{k}\vec{v}_j} \exp(i t_{j-1}[\omega - \vec{k}\vec{v}_{j-1}]) \right. \\ \left. (1 - \exp(i \xi_j[\omega - \vec{k}\vec{v}_j])) \right|^2 \quad (3)$$

where $\xi_j \equiv t_j - t_{j-1}$ is the time between two collisions.

In a thermal medium many particle trajectories are possible. Thus in order to calculate the total energy radiated by the incident particle one needs to take into account the stochastic nature of the particle interactions and perform an average in the

above equation over all possible particle trajectories. This can be done analytically in two limiting cases: (i) assuming that after a collision the particle velocities are isotropically distributed (isotropic scattering) and (ii) assuming small angle scattering. These two limiting cases correspond to the two distinct phases of QCD matter at finite temperature. The isotropic scattering is approximately valid for pions as the fundamental constituents. Pion-pion elastic scattering is mainly due to s-wave interactions with nearly isotropic final momentum distribution. The small-angle scattering holds in the deconfined phase where the quark interactions are mainly determined by the t-channel Feynman diagrams. This means that small-angle scattering dominates and particle velocities are strongly correlated in successive collisions.

The average over trajectories in a pion gas and in a quark-gluon plasma require separate investigations as the analytical procedure is very different [16-19].

Bremsstrahlung in a pion gas

As was explained in the previous section the calculation of the cross-section due to particle bremsstrahlung is reduced to finding the average of eq.(3) over all possible particle trajectories in a medium. In a pionic gas the average over trajectories is equivalent to the average over all particle velocities after the collision and over all possible times between two collisions. This is because a given trajectory of a charged pion is parameterized by the scattering angle and the time between collisions. As the pionic gas is assumed to be in thermal equilibrium, many possible velocities can result in between two collisions. Because of this, upon taking the square in eq.(3), all non-diagonal terms will give zero contribution. Indeed, a typical non-diagonal term in eq.(3) has the form:

$$\frac{[\vec{n} \times \vec{v}_j] \cdot [\vec{n} \times \vec{v}_l]}{(\omega - \vec{k} \cdot \vec{v}_j) \cdot (\omega - \vec{k} \cdot \vec{v}_l)} \exp[i(\omega(t_j - t_l) - \vec{k} \cdot (\vec{R}_j - \vec{R}_l))] \times (1 - e^{-i(\omega - \vec{k} \cdot \vec{v}_j)\xi_j})(1 - e^{-i(\omega - \vec{k} \cdot \vec{v}_l)\xi_l}) \quad (4)$$

where $t_j, t_l, \vec{R}_j, \vec{R}_l$ are the time and the position of the corresponding scattering event; $\vec{v}_j$ is the velocity of the test particle after j -th collisions and $\vec{n}$ is the unit vector in the direction of the photon emission.

The above expression contains a product of two main factors: the term $[\vec{n} \times \vec{v}_j]$ and the rest. These two terms have a different dependence on $\vec{v}_j$. Indeed, the term $[\vec{n} \times \vec{v}_j]$ depends only on transversal (with respect to the direction of photon emission) component of $\vec{v}_j$, whereas all the rest depends only on its longitudinal component.

To find the final result for the intensity of particle radiation one needs to average the probability of the photon emission over all possible trajectories of the test particle. Our main assumption in the pionic medium is that the scattering keeps no correlation between final and the initial particle velocity. Therefore, one can always find (with equal probability) two trajectories which coincide between points $\vec{R}_l$ and $\vec{R}_j$, but have opposite transversal components of the particles velocity after j -th collision. Thus the off-diagonal terms have different signs for the above trajectories. Consequently, after averaging over velocities these will cancel, leading in general to the vanishing of all non-diagonal terms in eq.(3). We have to point out however that these cancellations are only due to our strict assumption about complete isotropy of particle velocities in each scattering. The deviation from this assumption will naturally imply non-vanishing contributions of some non-diagonal terms in eq.(3). Nevertheless, the pionic medium is very close to the isotropic limit and we assume that only the diagonal terms remain significant.

The average over time between two successive collisions ξ is performed assuming an exponential distribution

$$\frac{dW}{d\xi} = a \exp(-\xi a) \quad (5)$$

with a being the average number of collisions per unit time (inverse mean free time)

$$a(p, T) \equiv g \int \frac{d^3 q}{(2\pi)^3} e^{-\sqrt{p^2 + m_\pi^2}/T} \sigma_{\pi\pi}(s) v \quad (6)$$

where $\sigma_{\pi\pi}$ is the elastic pion-pion cross-section, $v = \sqrt{s(s - 4m_\pi^2)}/2E_1 E_2$ the relative particle velocity, $s = (p + q)^2$ and g is the degeneracy factor (3 for the pion gas). With the ξ -distribution function (5) the time average in (3) can be done analytically. When relating the intensity of the radiation $dI/d^3 k$ with the production cross-section

$$\frac{dI^\gamma}{d^3 k} = N \omega \frac{1}{\sigma_{\pi\pi}} \frac{d\sigma^\gamma}{d^3 k} \quad (7)$$

the final result for the bremsstrahlung cross-section, after averaging over all possible particle trajectories reads [15]:

$$\langle \frac{d\sigma^\gamma}{d\omega d\Omega} \rangle = \frac{2\alpha\sigma_{\pi\pi}|\vec{v}|^2}{4\pi^2} \omega \frac{\sin^2 \theta}{a^2 + \omega^2(1 - |\vec{v}| \cos \theta)^2} \quad (8)$$

with $d\Omega$ being the unit element of the solid angle.

The medium modification of the bremsstrahlung cross-section due to LPE is in the denominator in the above expression through the term $a(p, T)$. If this term is absent then we are recovering the standard formula for charged particle bremsstrahlung. How important the modification of the cross-section due to multiple scattering effect is can be seen by considering two limiting solutions for the cross-section in eq.(8). The mean free path λ of incident pions in a thermal medium is given by $\lambda = v a^{-1}$. The coherence length according to (1) is determined from $l_c = v[\omega(1 - v \cos(\theta))]^{-1}$. With the above definitions one can find that

$$\langle \frac{d\sigma^\gamma}{d\omega d\Omega} \rangle = \left(\frac{d\sigma^\gamma}{d\omega d\Omega} \right)^{(1-\text{sc.})} \times \begin{cases} 1 & \text{if } \lambda \gg l_c \\ (\frac{\lambda}{l_c})^2 & \text{if } l_c \gg \lambda \end{cases} \quad (9)$$

Thus, when the mean free path exceeds the coherence length the medium acts as the sum of N independent scatterers. In the opposite limit the interference effect within the coherence length starts to be important. The cross-section is modified by the factor $(\lambda/l_c)^2$ leading to the suppression due to LPE. The destructive interference of waves within the coherence length is also seen when analyzing frequency dependence of the bremsstrahlung cross-section. From (8) one can find that

$$\left\langle \frac{d\sigma^\gamma}{d\omega d\Omega} \right\rangle \sim \begin{cases} \frac{1}{\omega} & \text{if } \lambda \gg l_c \\ \omega & \text{if } l_c \gg \lambda \end{cases} \quad (10)$$

In the limit of very soft photon, the cross-section decreases linearly with the frequency of the radiation which is in contrast to the standard result which shows $1/\omega$ behaviour. With increasing ω the coherence length is decreasing. Thus, for the value of ω where $\lambda \sim l_c$ the LPE is not important any more and typical $1/\omega$ behaviour for the bremsstrahlung cross-section is again recovered.

Until now we have discussed real photon bremsstrahlung in a pionic medium and derived the cross-section for this electromagnetic process. Assuming that both the momenta and the invariant mass M of virtual photon are small in comparison with the particle energy then under some approximation the formula (8) can be also applied for virtual photon bremsstrahlung by simple replacement of ω by $\sqrt{k^2 + M^2}$.

Bremsstrahlung in a quark-gluon plasma

The averaging over trajectories in the case of quark bremsstrahlung requires a separate investigation. This is mainly because the dynamics in a quark-gluon plasma is different from a pionic medium. The quark cross-section in a plasma is dominated by small angle scattering. Thus, the quark velocities in two subsequent collisions are strongly correlated. In this case non-diagonal terms in (3) cannot be neglected as they will have important contributions to the total cross-section. In order to calculate the energy radiated by a quark moving in a thermal plasma it is convenient to rewrite (2) into the following form [9,16]:

$$\frac{dI}{d^3k} = \frac{e_q^2 \alpha}{(2\pi)^2} 2\text{Re} \int_0^\infty dt \int_0^\infty d\tau K(t, \tau) \quad (11)$$

where the kernel K is defined as:

$$K(t, \tau) \equiv \exp[-i\omega\tau - \vec{k} \cdot (\vec{r}[t + \tau] - \vec{r}[t])] (\vec{v}[t + \tau] \cdot \vec{v}[t] - (\vec{n} \cdot \vec{v}[t + \tau])(\vec{n} \cdot \vec{v}[t])) \quad (12)$$

The problem consists now in averaging the above expression over all possible quark trajectories in a plasma. This can be done using functional integral techniques [20] or the method originally proposed by Migdal [9]. In the following we will apply Migdal's prescription to find the indicated averaging. This is done introducing two probability functions: $w_1(\vec{r}, \vec{v}; t)$ - the probability to find at the time t a particle with coordinate $\vec{r}$ and velocity $\vec{v}$ and $w_2(\vec{r}, \vec{v}; t | \vec{r}', \vec{v}'; t')$ - the conditional probability of value $\vec{r}$ and $\vec{v}$ at the time t under the condition that at the moment of time t' these quantities have had the values $\vec{r}'$ and $\vec{v}'$. Then the averaging over trajectories is equivalent to finding the integral

$$\langle K \rangle = \int d\vec{r} d\vec{r}' d\vec{v} d\vec{v}' w_1 w_2 K \quad (13)$$

The probability functions w_1, w_2 satisfy the kinetic equation

$$\frac{\partial w}{\partial t} + \vec{v} \frac{\partial w}{\partial \vec{r}} = n(T) \int d\vec{v}' \sigma(\vec{v}' - \vec{v}) [w(\vec{r}, \vec{v}'; t) - w(\vec{r}, \vec{v}; t)] \quad (14)$$

where σ is the elastic quark cross-section changing the velocity from $\vec{v}'$ to $\vec{v}$ and $n(T)$ is the density of the medium. The solution of the above differential equation can be found using a Fokker-Planck approximation [9,17] with the initial conditions

$$w_1(\vec{r}, \vec{v}; 0) = \delta(\vec{r}) \delta(\vec{v} - \vec{v}_0) \quad (15)$$

$$w_2(\vec{r}, \vec{v}; t | \vec{r}', \vec{v}'; t) = \delta(\vec{r} - \vec{r}') \delta(\vec{v} - \vec{v}') \quad (16)$$

where $\vec{v}_0$ denotes the initial quark velocity.

Using the solutions for the probability function in eqs.(11,13) the integrations can be done analytically with the following final result for the energy radiated by a quark due to virtual photon bremsstrahlung [17]:

$$\left\langle \frac{dI^\gamma}{d\omega dt} \right\rangle \simeq \frac{8e_q^2 \alpha}{3\pi} \frac{k}{\omega} q(p, T) \left[\left(\frac{m_q^2}{p^2} + \frac{M^2}{\omega^2} \right)^2 + \frac{16q(p, T)}{9\omega} \right]^{-1/2} \quad (17)$$

where $q(p, T) \equiv \frac{1}{4} v \langle \theta^2 \rangle$ is the characteristic parameter of small angle scattering and

$$\langle \theta^2 \rangle \equiv n(T) v \int \theta^2 \frac{d\sigma}{d\Omega} d\Omega \quad (18)$$

is the mean square scattering angle per unit path.

The influence of the Landau-Pomeranchuk effect on the virtual photon emission by a moving quark in a medium is contained in the last term in the square bracket in eq.(17). If this term is negligibly small compared to the first one, then the effect is absent. It is easy to see that the above expression remains finite in the limit of massless quarks. Thus, the scattering of particles in a medium regularizes the infrared behaviour of the corresponding amplitudes. For on-shell photons formula (17) coincides with the one previously obtained by Migdal [9].

The mean square of the angle of multiple scattering for the coherence length l_c is defined as

$$\theta_{MS} \equiv l_c \sqrt{\langle \theta^2 \rangle} \quad (19)$$

whereas a typical emission angle of the bremsstrahlung photon $\theta_B \sim m_q/E$. With these two parameters the energy spectrum of real photon (17) can be written as:

$$\left\langle \frac{dI^\gamma}{d\omega dt} \right\rangle \simeq \frac{8e_q^2 \alpha}{3\pi} \left(\frac{p}{m_q} \right)^2 q(p, T) \frac{1}{[1 + \frac{2}{9} (\frac{\theta_{MS}}{\theta_B})^2]^{1/2}} \quad (20)$$

Thus, there are two distinct results in the limit when $\theta_{MS} < \theta_B$ and $\theta_{MS} > \theta_B$:

$$\left\langle \frac{d\Gamma}{d\omega dt} \right\rangle = \left[\frac{d\Gamma}{d\omega dt} \right]_{\text{Bethe-Heitler}} \times \begin{cases} 1 & \text{if } \theta_{MS} < \theta_B \\ \sqrt{\frac{g}{2}} \frac{\theta_B}{\theta_{MS}} & \text{if } \theta_{MS} > \theta_B \end{cases} \quad (21)$$

From the above equation one can conclude that if the multiple scattering angle θ_{MS} is smaller than a typical bremsstrahlung angle then there is no medium modification of quark bremsstrahlung cross-section as it is determined by the well known Bethe-Heitler formula. However, the radiation spectrum is suppressed if at the distance l_c the scattering deflects the test quark through an angle θ_{MS} larger than the typical emission angle θ_B .

In the case of small angle scattering the relevant parameter which measures the LPE is the mean square of scattering angle for the coherence length. This is in contrast to the isotropic case where the important parameter was the mean free path of the test particle. Thus, particle bremsstrahlung in a quark-gluon plasma and in a pion gas required different conditions to see Landau-Pomeranchuk suppression effect. From (21) one can also find the following ω dependence of the cross-section

$$\left\langle \frac{d\Gamma}{d\omega dt} \right\rangle(\omega) \sim \begin{cases} \text{const.} & \text{Bethe - Heitler} \\ \sqrt{\omega} & \text{Landau - Pomeranchuk} \end{cases} \quad (22)$$

The photon energy spectrum described by the Bethe-Heitler formula is known to be independent of the frequency of the radiation. The Landau-Pomeranchuk-Migdal effect changes qualitatively this result leading to $\sqrt{\omega}$ -behavior. The $\sqrt{\omega}$ -behavior has been recently verified experimentally in SLAC experiment E-146 [21] where the production of photons due to bremsstrahlung of electron beams traversing uranium target was measured.

As it was explained above the role of the LPE in a quark-gluon plasma is entirely determined by the parameter $\langle \theta^2 \rangle$. To lowest order in α_s , the small-angle elastic cross-sections in a quark gluon plasma are described by the t -channel exchange Feynman diagrams which give:

$$\frac{d\sigma_{ab \rightarrow ab}}{dt} = \frac{2\pi\alpha_s^2}{t^2} C_{ab} \quad (23)$$

$$C_{ab} = \begin{cases} 4/9 & \text{for } qq \rightarrow qq \\ 1 & \text{for } qg \rightarrow qg \end{cases} \quad (24)$$

When integrating the cross-sections in eq.(18) over $d\Omega$ with a weight factor θ^2 , one finds a typical logarithmically divergent integral

$$\int_{\theta_{min}}^{\theta_{max}} \frac{d\theta}{\theta} = \ln \left[\frac{\theta_{max}}{\theta_{min}} \right] \quad (25)$$

To estimate this integral one chooses the upper limit θ_{max} to be close to unity: this is the upper boundary for the validity of the small-angle approximation. The

lower limit θ_{min} in a Coulomb-like medium is usually cut-off by many-body effects like charge screening. In the plasma a typical screening scale is a Debye mass m_D proportional to gT with g being the coupling constant and T the temperature of the medium [22]. Therefore, $\theta_{max}/\theta_{min} \sim p/gT$, where p is the momentum of the charged particle. In a thermal system, however, temperature gives a natural scale for the particle momentum. Thus,

$$\langle \theta^2 \rangle \sim \frac{T^3}{p^2} \alpha_s^2 \ln \alpha_s, \quad (26)$$

The above parameter is shown to be an increasing function of the temperature. Thus, also the suppression of bremsstrahlung is more effective with increasing temperature as seen in (21). This is to be expected since with increasing temperature the density of the medium is also increasing and consequently multiple scattering is more important.

Multiple scattering and charge screening are not the only collective effects which influence the bremsstrahlung cross-section. Additionally there is the dynamical generation of a thermal mass for fermionic degrees of freedom in the medium. According to [22] fermions in a quark-gluon plasma acquire a thermal mass $m_f^2 \sim 2\pi\alpha_s/3T$ induced by many-body dynamical effects. From eq.(17), however, one can deduce, that when considering the process of real and virtual photon bremsstrahlung in a plasma, the thermal mass generation effect can be small compared with to the multiple scattering one. Indeed, taking the limit $M \rightarrow 0$ in (17), one can see, that the second term in the square brackets is of the order of $\alpha_s^2 \ln \alpha_s$, while the first one is only of the order of α_s^2 . For this reason the LPE is mainly of importance when discussing the soft electromagnetic signal from a quark-gluon plasma.

SOFT DILEPTON PRODUCTION IN HEAVY ION COLLISIONS

In the last section we have derived the cross-section for real and virtual photon bremsstrahlung in a pion gas and in a quark-gluon plasma under the soft photon approximation. We have discussed qualitatively the modifications of these cross-sections due to the Landau-Pomeranchuk suppression effect. Having in mind phenomenological application of these results in heavy ion collisions we will discuss in this section dilepton production due to bremsstrahlung. The cross-section for dilepton production can be related with the one for virtual photon as follows:

$$\frac{d\sigma^{e^+e^-}}{dM} = \frac{2\alpha}{3\pi M} \sqrt{1 - \frac{4m_l^2}{M^2}} \left(1 + \frac{2m_l^2}{M^2} \right) \sigma^*, \quad (27)$$

To obtain $d\sigma^{e^+e^-}/dM$ one needs to integrate the cross-section in a pion gas and in a quark gluon plasma, derived in the previous section, over the virtual photon frequency. The difficulty here, however, is that this procedure is out of the range of the soft photon approximation. We have imposed the upper integration limits for the frequency integral with the help of the energy-momentum conservation laws. [17-19]

In heavy ion collisions one is more interested in the dilepton production rate rather than the cross-section. The expression for the rate is given by kinetic theory*

* In the quark-gluon plasma case one averages over collisions instead of integrating over the scatterer's momentum distribution [17].

$$\frac{dN^{e^+e^-}}{d^4x} = \int d^3p_1 f(\vec{p}_1) \int d^3p_2 f(\vec{p}_2) |\bar{v}_{12}| \frac{d\sigma^{e^+e^-}}{dM^2} \quad (28)$$

where the momentum distribution function $f(p)$ of the incident particle and of the particles in the medium are assumed to be thermal and described by a Boltzman distribution function.

Figure 1: Invariant mass distribution of dilepton pair at $T = 0.2 \text{ GeV}$.

In fig.1 we show the dilepton production rate in a pion gas (*full line*) and in a quark gluon plasma (*dashed-line*) due to particle bremsstrahlung with and without LPE taken into account. The calculations are performed at fixed temperature $T=0.2 \text{ GeV}$. As can be seen in this figure, the Landau-Pomeranchuk effect leads to a significant reduction of dilepton production due to bremsstrahlung of virtual photons from both quark and pion moving in a thermal medium. The reductions of the rates are particularly important in the low invariant mass region. With increasing M the Landau-Pomeranchuk suppression is reduced. This behaviour can be understood intuitively as follows. The time necessary for e^+e^- pair production is proportional to its inverse mass. Increasing invariant mass means also decreasing the production time and the coherence length. Thus, the amount of collisions of the test particle with a surrounding medium constituents within l_c is decreasing as well. Consequently, the medium effect is less important. We have to stress however, that when going beyond the soft photon approximation the suppression can be seen for larger invariant mass values. Our approximations should be valid in the part of the spectrum where the invariant mass is less than $M \sim 0.3 - 0.5 \text{ GeV}$. The results presented in fig.1. show that in the soft part of $dN/d^4x dM$ spectrum quark bremsstrahlung is of the crucial importance. At fixed temperature the medium suppression of virtual photons radiated from a charged particle is larger in a plasma than in a hadron gas. This behaviour is certainly not

only related with differences between the cross-sections as discussed in the previous section. At a fixed temperature the number of degrees of freedom and consequently the density of a plasma is higher than in a hadron gas. This naturally implies higher suppression in a quark-gluon plasma.

The suppression of dilepton production rate due to LPE is not only increasing with decreasing energy of bremsstrahlung photon but also with increasing temperature. This is illustrated in fig.2 where the rate in a plasma is calculated with and without LPE at fixed invariant mass $M = 0.1 \text{ GeV}$ as a function of temperature.

Figure 2: Rate in a plasma at fixed M as a function of T .

Until now we have studied dilepton production in a medium at fixed temperature. We have to point out, however, that hadronic matter produced in heavy ion collisions is not a static object. It rather undergoes hydrodynamical expansion leading to a time dependence of temperature and of other thermodynamical quantities. The experimentally measured distributions are integrated over the space-time history of heavy ion collisions. In fig.3 we show the results for soft dilepton production rate assuming Bjorken expansion for the space-time evolution of hot hadronic matter produced in a heavy ion collisions [19]. The initial formation time $\tau_0 \sim 0.07 \text{ fm}$ and initial temperature $T \sim 0.9 \text{ GeV}$ are relevant for U-U collisions at LHC energy [23]. The comparison of quark (*short-dashed line*) and pion (*full line*) bremsstrahlung with non thermal background due to η (*long dashed line*) and π_0 (*dashed-dotted line*) Dalitz decay shows the importance of bremsstrahlung process as a source of soft dilepton production. In the window of invariant mass $0.1 \leq M \leq 0.3$ pion bremsstrahlung acids Dalitz eta decay, quark bremsstrahlung and basic quark annihilation contribution (*circles*). For the invariant mass $M > 0.5 \text{ GeV}$ the pion annihilations (*dotted line*) is a dominating process for dilepton productions. Nevertheless, one might expect to observe a quark-gluon plasma signal when measuring double differential spectra.

$dN^{e^+e^-}/dydMdp_t$, at high transversal virtual photon momenta, p_t and at low invariant mass. The reason is that a typical momentum distribution is related with the temperature of the emitting system. The plasma phase exists at higher temperatures than the pion gas phase, therefore, the momentum distribution of the pairs emitted from the plasma phase is expected to be broader than that from the hadronic phase. Hence, at some high enough value of the transversal momentum of the pair, the plasma bremsstrahlung should dominate $dN^{e^+e^-}/dydMdp_t$ spectrum. This is illustrated in fig.4 where the double differential production rate is calculated at fixed invariant mass $M = 0.15\text{GeV}$. At large transverse momentum quark bremsstrahlung (dashed line) is indeed dominating, both Dalitz eta decay contribution (dashed-dotted line) and pion bremsstrahlung (full line). One should, however, keep in mind that the above results are obtained assuming that there is no transverse expansion or that it is very slow. The hydrodynamical model with fast transverse expansion in the late stage of the collision could influence the results in fig.4. This is because transverse expansion would imply an additional boost

Figure 3: Dilepton rate in the expanding medium.

of the dilepton pair in the transverse direction, increasing in this way the rate of pion bremsstrahlung for the pair with large transverse momentum. Also the preequilibrium background could still modify the production rate in fig.4 in the kinematical window of large p_t .

CONCLUSIONS

In this work soft dilepton production due to real and virtual photon bremsstrahlung in hot QCD medium was studied. We have shown that the bremsstrahlung process can be substantially suppressed in a thermal medium due to the Landau-Pomeranchuk effect. This effect is particularly important at high temperature and at the low energy of the bremsstrahlung photons. The analysis of soft dilepton production

in an expanding medium shows that particle bremsstrahlung is the main source of thermal pairs.

Figure 4: Dilepton rate at fixed value of M as a function of p_t .

Assuming optimistic initial conditions for the LHC experiment the thermal bremsstrahlung can exceed the contribution due to the Dalitz eta decay in the kinematical window where the invariant mass $m_\pi < M < 2m_\pi$. These results indicate that thermal dilepton production could be experimentally observed by measuring the soft part of the dilepton production rate at LHC energy.

Acknowledgments

One of us (KR) acknowledges stimulating discussions with R. Baier, E. L. Feinberg, H. Satz and H.J. Specht. This work was partly supported by the Ministry for Research and Technology (BMFT) and the Committee of Scientific Research (KBN).

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