Thermal photons as a measure for the rapidity dependence of the temperature

A. Dumitru, U. Katscher, J.A. Maruhn, H. Stöcker, W. Greiner

Institut für Theoretische Physik der J.W. Goethe Universität

Postfach 111932, D-60054 Frankfurt a.M., Germany

D.H. Rischke

Physics Department, Pupin Laboratories, Columbia University

New York, NY 10027, USA

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Abstract

The rapidity distribution of thermal photons produced in Pb+Pb collisions at CERN-SPS energies is calculated within scaling and three-fluid hydrodynamics. It is shown that these scenarios lead to very different rapidity spectra. A measurement of the rapidity dependence of photon radiation can give cleaner insight into the reaction dynamics than pion spectra, especially into the rapidity dependence of the temperature.

One of the goals of heavy-ion physics is the search for the so-called quark-gluon plasma (QGP), a novel phase of matter where quarks and gluons are deconfined. It is expected that such a state can be produced, e.g., when ordinary hadronic matter is strongly heated or compressed [1].

Real and virtual photons are promising probes of the QGP [2] (they may serve, e.g., as a thermometer [3]) since they do not suffer from strong interactions. Therefore, their mean free path is large enough [4] for them to leave the plasma volume without further interactions.

Transverse momentum distributions of thermal photons produced in $Pb+Pb$ collisions at 160 $AGeV$ were already presented in refs. [5, 6, 7]. Recently, also the rapidity distribution of electromagnetic radiation was investigated in more detail [8, 9, 10]. In these works, however, the collision dynamics was simplified by assuming a scaling hydrodynamics solution for the longitudinal motion, $v_z = z/t$ [11]. This limits the usefulness of the results of [8, 9, 10] to future collider experiments. In this letter we present rapidity distributions of thermal photons in $Pb(160\ AGeV) + Pb$ (at vanishing impact parameter, $b = 0$), calculated within the three-fluid hydrodynamical model, and show that the photons are very useful to constrain the reaction dynamics at these energies.

Let us first give a brief introduction into the three-fluid model. For a more detailed presentation, we refer the reader to refs. [7, 12, 13]. The original one-fluid hydrodynamic model [14] assumes instantaneous local thermodynamic equilibrium in the moment when the nuclei collide and thus is not appropriate to describe the initial stage of ultrarelativistic collisions, at least for $E_{Lab} \geq 10\ AGeV$. This problem is solved here by considering more than one fluid [15]. The three-fluid model divides the particles involved in a reaction into three separate fluids: the first two fluids correspond to the projectile and target nucleons, respectively, and the particles produced during the reaction are collected in the third fluid. Local thermodynamic equilibrium is maintained only in each fluid separately but not between the fluids. The fluids are able to penetrate and decelerate each other during the collision. Interactions between the fluids are due to binary collisions of the particles in the respective fluids. This allows for a treatment of non-equilibrium effects in the initial stage of the collision.

The equation of state (EOS) of the target and projectile fluids is that of an ideal nucleon gas plus compression energies. A linear ansatz for the compression energy with a compressibility of 250 MeV and a binding energy of 16 MeV is used [16].

The EOS of the third fluid is that of an ideal gas of massive π-, ρ-, ω-, and η-mesons. At $T_C = 160\ MeV$ we assume a first order phase transition into a QGP. For the (net baryon-free) QGP we use the bag-model EOS for (pointlike, massless, and noninteracting) u and d quarks. The bag constant is chosen such that the pressures of both phases coincide at $T = T_C$.

For comparison, we also perform calculations within one-dimensional scaling hydrodynamics [11] where we also use the latter EOS. Here, the compressional stage of the collision is not treated and thus two free parameters, the initial temperature T_i and (proper) time τ_i, where the scaling expansion starts, have to be fixed. We use the values given in ref. [17]: $T_i = 300\ MeV$, $\tau_i = 0.22\ fm$. At later times, the temperature in this model is given by ($T_i \equiv T_i^{(1)}$, $\tau_i \equiv \tau_i^{(1)}$)

$$T(\tau) = T_i^{(j)} \left(\frac{\tau_i^{(j)}}{\tau} \right)^{c_j^2} \quad , \tag{1}$$

where $j = 1, 2, 3$ labels the different phases of matter (QGP, mixed phase, and hadron gas) and c_j denotes the sound velocity in the corresponding phase.

All quantities, except for $T_i^{(1)}$ and $\tau_i^{(1)}$, are determined by the equation of state, cf., e.g., ref. [5].

The thermal photon production rate from an equilibrated, baryon-free QGP is given (to first order in α and α_S) by [3]

$$E\frac{dR^\gamma}{d^3k} = \frac{5\alpha\alpha_S}{18\pi^2}T^2 e^{-E/T}\ln\left(\frac{2.912E}{g^2T}+1\right) \quad , \tag{2}$$

where E is the photon energy in the local rest frame of the fluid. In the following calculations we fix $\alpha_S = g^2/4\pi = 0.4$. As shown in ref. [3], the rate for a gas consisting of π-, ρ-, ω-, and η-mesons may also be parametrized by eq. (2). Other contributions, e.g. from the A_1 meson [18], as well as the effect of hadronic formfactors [3], are neglected. They are of the same order of magnitude as higher order corrections to eq. (2), which have also not been taken into account. Thus, eq. (2) is applied for all phases. The contributions from the first two fluids are neglected since, for the reactions considered here, these fluids are cooler. Also, they undergo a rapid longitudinal expansion and thus cool faster than the third fluid. This approximation must, of course, break down at large rapidities, where the temperature of the third fluid drops rapidly, as discussed below. As a check, eq. (2) has also been applied to the projectile and target fluids. This certainly overestimates their contribution because they contain only baryons. It turns out that, indeed, their contribution to thermal radiation is negligible up to photon rapidities

$\simeq 1.6$.

The thermal photon spectrum is obtained by an integration over space-time:

$$\frac{d^2N^\gamma}{k_T\,dk_T\,dy} = \int d^4x\,E\frac{dR^\gamma}{d^3k} \quad . \tag{3}$$

Fig. 1 shows our results. In the three-fluid model the temperature is strongly rapidity dependent and so is the spectrum of thermal photons. As already noted in ref. [7], there is much more high transverse-momentum radiation at midrapidity than in scaling hydrodynamics, which is mainly due to the different cooling law. In Bjorken's original model, the temperature is a function of proper time only and independent of (fluid-) rapidity, and thus the thermal photons show no rapidity dependence. This remains true even if transverse expansion [5, 19] is implemented into scaling hydrodynamics.

At this point we should comment on the different freeze-out procedures in the two models. In the three-fluid model the freeze-out of the third fluid is done instantaneously at some center-of-mass time t_f when its (average) temperature drops below $T_f = 100\ MeV$. On the other hand, in the Bjorken model, the freeze-out takes place at some proper time τ_f defined by $T(\tau_f) = T_f = 100\ MeV$. Thus, the space-time volume for which the photon spectrum (3) is determined, is different in both cases. However, at least as far as "hard" photons are concerned (i.e., $k_T \gg T_i$) this difference is irrelevant: the shape

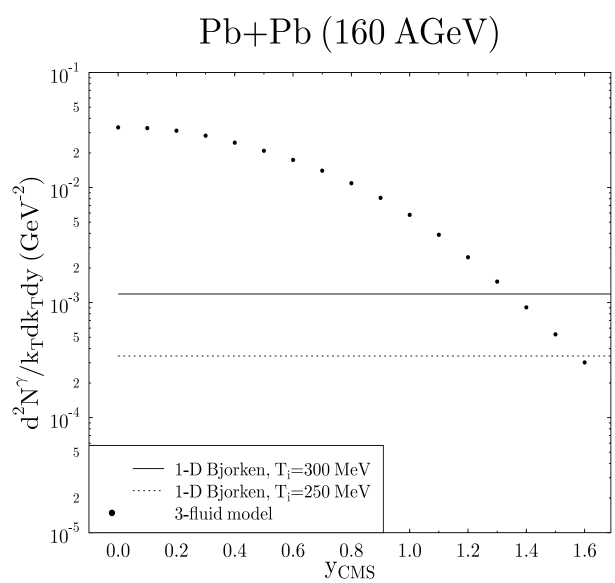

Figure 1: Rapidity distribution of thermal photons (for $k_T = 2\ GeV$) calculated within the three-fluid model (dots) and longitudinal scaling hydrodynamics.

of the space-time volume differs only for the very late stage ($t \approx t_f$), when the temperature is too low to give a sizeable contribution to the "hard" photon yield.

In contrast to the three-fluid model, effects of finite baryon-chemical potential (which appears due to the expected large amount of baryon stopping [12, 20]) on the temperature of the mesons are not accounted for in the Bjorken model as presented so far. A finite μ_B will mainly manifest itself in lower initial temperatures whereas the *net*[1] effect of $\mu_B \neq 0$ on "hard" (i.e., $k_T \gg T_i$) photon radiation from the QGP phase was shown to be small [21]. We assume that this remains true also in the hadronic phase. Using the standard argument that relates the final pion multiplicity to the initial entropy through the assumption of entropy conservation during the whole expansion stage [17, 22], and the expression

$$s = T^3 \left(\frac{4}{3} \frac{37\pi^2}{30} + \frac{2\mu_B^2}{9T^2} \right) \tag{4}$$

for the entropy density in the QGP phase, it turns out that values $3 \leq \mu_B/T \leq 9$ [23] diminish T_i by $4 - 21\%$. For simplicity let us assume that the ratio of pions to baryons and thus μ_B/T is independent of rapidity. The effect on the thermal photon radiation is shown in fig. 1: the dotted curve was calculated using an initial temperature of 250 MeV, all other parameters

[1]That is, the change of the photon rate with μ_B at fixed T.

being the same as before. If μ_B/T is taken to be rapidity dependent this may introduce some rapidity dependence into the Bjorken model. However, the resulting photon rapidity spectrum would still lie between the dotted and the full curve in fig. 1.

In principle it would be possible to introduce a (fluid-) rapidity dependent initial temperature into the Bjorken model by fitting the pion rapidity distribution *for all rapidities* [8, 9, 10]. Then, eq. (1) is applied in each rapidity slice separately, i.e. $T_i \to T_i(\eta)$, $T(\tau) \to T(\tau, \eta)$, and no longer globally. We will, however, not adopt this procedure here since it introduces an infinite number of parameters and contradicts the original philosophy of scaling hydrodynamics as a simple model with only two free parameters (T_i, τ_i). Also, this procedure leads to a violation of the conservation laws expressed by the hydrodynamic equations of motion unless a finite (and rapidity dependent) baryon-chemical potential is introduced at the same time [9]. It can be shown that for longitudinal scaling expansion (even including a cylindrically symmetric transverse expansion) $\partial p(T, \mu)/\partial\eta = 0$, or, equivalently,

$$s\frac{\partial T}{\partial \eta} + n\frac{\partial \mu}{\partial \eta} = 0 \qquad (5)$$

has to hold (the partial derivatives with respect to η being performed at constant τ), where s and n denote the entropy and particle number, respectively.

From simple arguments one may expect the local photon density to be

Pb+Pb (160 AGeV)

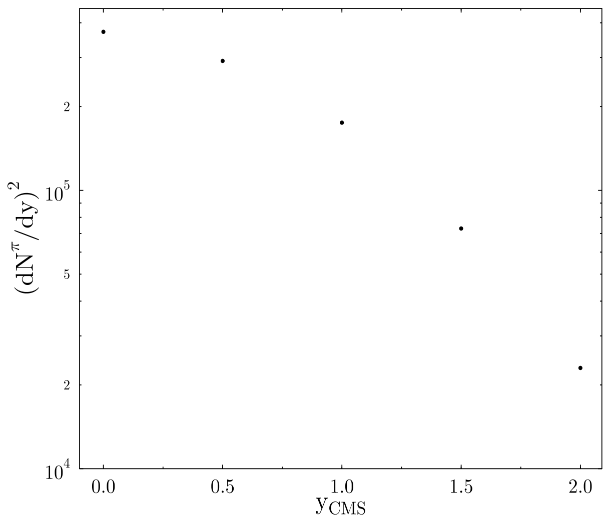

Figure 2: Squared pion rapidity density.

proportional to the square of the local pion density. However, in the three-fluid model this proportionality is not maintained in the finally observed pion spectra. As one can see in fig. 2, the squared pion multiplicity obviously decreases much slower with rapidity. Also, the transverse momentum distribution of pions is softer than that of photons [7]. This can be readily explained with the help of fig. 3 which shows the rapidity distribution of the temperature of the third fluid. Note that the highest temperatures, which dominate "hard" photon and high-mass dilepton production, prevail only for short times where the rapidity distribution of produced particles is narrow (in the present model). This is due to the fact that the third fluid is

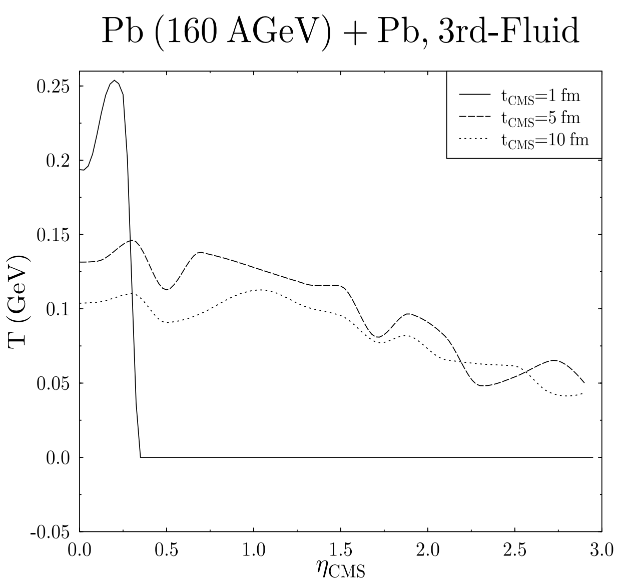

Figure 3: The (mean) temperature of the third fluid as a function of fluid-rapidity.

initially produced at midrapidity and broadens in rapidity-space during the expansion. On the other hand, the final pion distribution emerges at freeze-out where the temperature rapidity distribution is much broader. Therefore, the final pion (rapidity) distribution is also much broader (in rapidity space) than that of the photons. A measurement of the thermal photon rapidity spectrum would help to answer if this picture is correct (at these energies), or if, instead, the temperature distribution is broad in rapidity space from the very beginning of the expansion (as in scaling hydrodynamics).

In conclusion, the rapidity distribution of thermal photons in $Pb + Pb$

reactions at the SPS has been studied within three-fluid and scaling hydrodynamics. It has been demonstrated that thermal photons provide a powerful tool to constrain the reaction dynamics, i.e., the time and fluid-rapidity dependence of the temperature of the hot and dense reaction zone.

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References

[1] see, e.g.,

G. Baym in: “Quark Matter Formation in Heavy Ion Collisions“, eds. M. Jacob and H. Satz, World Scientific, Singapore, 1982

M. Creutz: “Quarks, gluons and lattices“, Cambridge Univ. Press, Cambridge, 1983

B. Müller, “The physics of the quark-gluon plasma”, Lecture Notes in Physics Vol. 225, Springer, New York, 1985

L. McLerran: Rev. Mod. Phys. 58 (1986) 1021

[2] E.V. Shuryak: Phys. Lett. B78 (1978) 150

L. McLerran, T. Toimela: Phys. Rev. D31 (1985) 545

B. Sinha: Phys. Lett. B128 (1983) 91; B157 (1985) 221

K. Kajantie, J. Kapusta, L. McLerran, A. Mekijan: Phys. Rev. D34 (1986) 2746

P.V. Ruuskanen: Nucl. Phys. A544 (1992) 169c

[3] J. Kapusta, P. Lichard, D. Seibert: Phys. Rev. D44 (1991) 2774

[4] M.H. Thoma: preprint UGI-94-04, Univ. Gießen, Germany (to be published in Phys. Rev. D)

[5] J. Alam, D.K. Srivastava, B. Sinha, D.N. Basu: Phys. Rev. D48 (1993) 1117

[6] N. Arbex, U. Ornik, M. Plümer, A. Timmermann, R.M. Weiner: preprint GSI-94-91 (to appear in Phys. Lett. B)

[7] A. Dumitru, U. Katscher, J.A. Maruhn, H. Stöcker, W. Greiner, D.H. Rischke: preprint UFTP-375/1994, Univ. Frankfurt, Germany (to appear in Phys. Rev. C)

[8] R. Vogt, B.V. Jacak, P.L. McGaughey, P.V. Ruuskanen: Phys. Rev. D49 (1994) 3345

[9] B. Kämpfer, O.P. Pavlenko, M.I. Gorenstein, A. Peshier, G. Soff: preprint FZR-50/1994, FZ Rossendorf, Germany

[10] S. Sarkar, D.K. Srivastava, B. Sinha: Phys. Rev. C51 (1995) 318

[11] J.D. Bjorken: Phys. Rev. D27 (1983) 140

[12] U. Katscher, J.A. Maruhn, W. Greiner: preprint UFTP-378/1994, Univ. Frankfurt, Germany

[13] U. Katscher, D.H. Rischke, J.A. Maruhn, W. Greiner, I.N. Mishustin, L.M. Satarov: Z. Phys. A346 (1993) 209

[14] see, e.g.,

J.A. Maruhn, W. Greiner: in “Treatise on Heavy–Ion Science“, Vol. 4, ed. D.A. Bromley,

Plenum Press, New York/London 1985, p. 565

R.B. Clare, D. Strottman: Phys. Rep. 141 (1986) 177

[15] A.A. Amsden, A.S. Goldhaber, F.H. Harlow, J.R. Nix: Phys. Rev. C17 (1978) 2080

A. Rosenhauer, J.A. Maruhn, W. Greiner, L.P. Csernai: Z. Phys. A326 (1987) 213

I.N. Mishustin, V.N. Russkikh, L.M. Satarov: Nucl. Phys. A494 (1989) 595

[16] W. Scheid, W. Greiner: Z. Phys. 226 (1969) 364

D.H. Youngblood, C.M. Rozsa, J.M. Moss, D.R. Brown, J.D. Bronson: Phys. Rev. Lett. 39 (1977) 1188

M.M. Sharma et al: Phys. Rev. C38 (1988) 2562

[17] J. Kapusta, L. McLerran, D.K. Srivastava: Phys. Lett. B283 (1992) 145

[18] L. Xiong, E. Shuryak, G.E. Brown: Phys. Rev. D46 (1992) 3798

C. Song: Phys. Rev. C47 (1993) 2861

[19] H. von Gersdorff, M. Kataja, L. McLerran, P.V. Ruuskanen: Phys. Rev. D34 (1986) 794

[20] A. von Keitz, L.A. Winckelmann, A. Jahns, H. Sorge, H. Stöcker, W. Greiner: Phys. Lett. B263 (1991) 355
D. Röhrich (NA35-collab.): Nucl. Phys. A566 (1994) 35c

[21] A. Dumitru, D.H. Rischke, H. Stöcker, W. Greiner: Mod. Phys. Lett. A8 (1993) 1291
C.T. Traxler, H. Vija, M.H. Thoma: preprint UGI-94-14, Univ. Gießen, Germany (subm. to Phys. Lett.)

[22] R.C. Hwa, K. Kajantie: Phys. Rev. D32 (1985) 1109

[23] M. Hofmann et al: Nucl. Phys. A566 (1994) 15c
Th. Schönfeld: PhD-thesis, Frankfurt 1993 (unpublished)