Off-equilibrium photon production in the chiral phase transition

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Introduction and motivation

In the early stage of heavy-ion collisions, chiral symmetry is temporarily restored. During this so called chiral phase transition, the quark masses change from their constituent value, m_c, to their bare value, m_b. This mass shift leads to the spontaneous pair production of quarks and antiquarks [1, 2]. We investigate the photon production arising from this pair creation process. We provide an ansatz, which eliminates the unphysical contribution from the vacuum polarization and renders the resulting photon spectra UV-finite if the time evolution of the quark masses is modeled in a suitable manner.

Chiral photon production

The change of the quark mass is modeled by coupling a scalar background field in a Yukawa-like manner to Dirac fermions in the QED-Lagrangian,

$$\mathcal{L}(x) = \mathcal{L}_{\text{QED}}(x) - g\phi(t)\hat{\bar{\psi}}(x)\hat{\psi}(x) \ . \qquad (1)$$

As the background field is assumed to be time dependent only, the fermions effectively acquire a time dependent mass

$$m(t) = m_c + g\phi(t) \ .$$

For our investigations on photon production induced by this mass change, we specify the boundary conditions for our system at $t_0 = -\infty$ and assume that it does not contain any quarks/antiquarks or photons initially. The coupling of the Dirac field to the external source field is treated exactly as to properly take into account the non-perturbative nature of the pair creation process. The photon yield is obtained by a standard perturbative quantum field theoretical (pQFT) calculation, where we restrict ourselves to first order QED processes. The contribution from these processes becomes possible since the quarks and antiquarks are rendered off-shell by the coupling to the external source field, $\phi(t)$. This leads to the following expression for the photon yield

$$2k\frac{\mathrm{d}^6 n_\gamma(t)}{\mathrm{d}^3x\mathrm{d}^3k} = \frac{1}{(2\pi)^3}\int_{-\infty}^{t}\mathrm{d}t_1\int_{-\infty}^{t}\mathrm{d}t_2\ \mathrm{i}\Pi_T^{<}(\vec{k},t_1,t_2) \cdot \mathrm{e}^{\mathrm{i}k(t_1-t_2)} \ . \qquad (2)$$

The photon self energy is given by the one-loop approximation

$$\mathrm{i}\Pi_{\mu\nu}^{<}(\vec{k},t_1,t_2) = e^2\int\frac{\mathrm{d}^3p}{(2\pi)^3}\mathrm{Tr}\left\{\gamma_\mu S_{\text{F}}^{<}(\vec{p}+\vec{k},t_1,t_2)\,\gamma_\nu S_{\text{F}}^{>}(\vec{p},t_2,t_1)\right\} \ , \qquad (3)$$

with the fermion propagators fulfilling the equations of motion

$$\left(\mathrm{i}\gamma^0\partial_{t_1} + \gamma^i p_i - m(t_1)\right)S_{\text{F}}^{\lessgtr}(\vec{p},t_1,t_2) = 0 \ , \qquad (4a)$$
$$\left(\mathrm{i}\gamma^0\partial_{t_2} - \gamma^i p_i + m(t_2)\right)S_{\text{F}}^{\lessgtr}(\vec{p},t_1,t_2) = 0 \ . \qquad (4b)$$

The photon self energy reduces to the vacuum polarization when t_1 and t_2 are taken from the domain where the quark mass is still at its initial constituent value, m_c. This contribution does not vanish in the limit $t_1, t_2 \to -\infty$ so that the time integrals entering (2) require the introduction of a regulator, $f_\varepsilon(t)$.

As suggested in [3], we follow a standard In/Out description, where the initial state is specified at $t_0 = -\infty$ and the photon yield is considered in the asymptotic limit $t \to \infty$. The regulator is hence chosen as

$$f_\varepsilon(t) = \mathrm{e}^{-\varepsilon|t|} \ , \qquad (5)$$

which serves as an adiabatic switching of the electromagnetic interaction. The photon numbers are then extracted from (2) for asymptotically free particles by taking the successive limits $t \to \infty$ and $\varepsilon \to 0$. Because of this, the resulting photon numbers do not contain unphysical contributions from the vacuum polarization. It must be emphasized that keeping the exact sequence of limits is crucial for that [4].

Numerical investigations and results

For our numerical investigations on photon production, we consider different mass parameterizations, $m(t)$, which are depicted in figure 1. As in [1, 2], we have chosen $m_c = 0.35$ GeV and $m_b = 0.01$ GeV. Furthermore, we assumed a change duration of $\tau = 1.0$ fm/c for both $m_2(t)$ and $m_3(t)$.

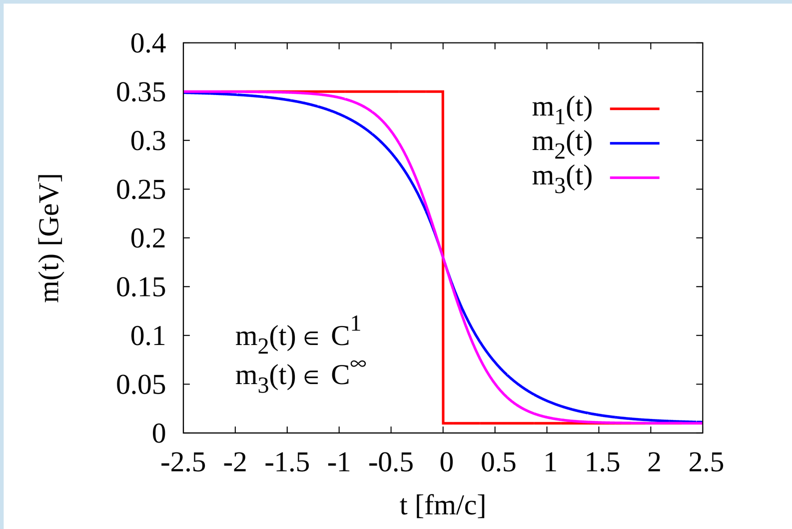

Figure 1: The change of the quark mass is modeled by different parameterizations, $m(t)$.

For the case of an instantaneous mass shift, the individual contributions to the photon numbers can be associated with first order QED-processes as well as with their interference among each other. The loop integral features a linear divergence resulting from the quark/antiquark occupation numbers decaying as $\sim 1/p^2$ for large p. As this decay behavior is an artifact of the instantaneous mass shift, we regulate this divergence by cutting the loop integral at $p = \Lambda_C$. Figure 2 shows the resulting photon spectra for different values of Λ_C.

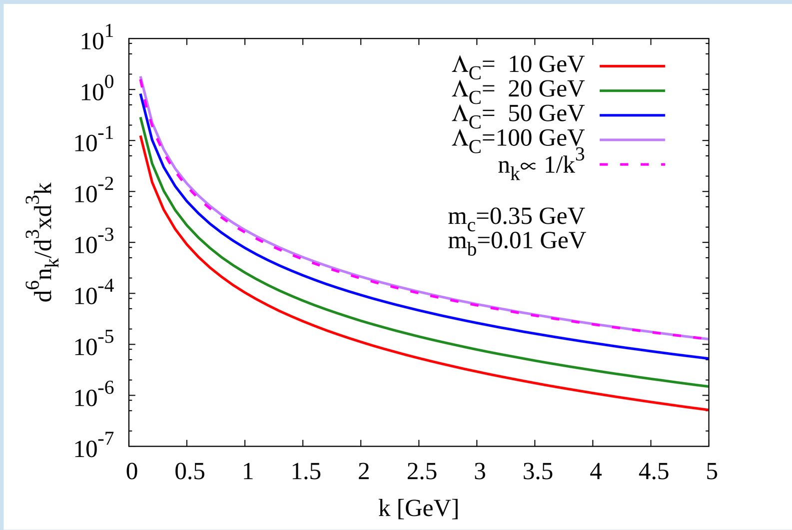

Figure 2: Asymptotic photon spectra for an instant mass shift, $m_1(t)$, and different values of Λ_C.

The photon spectra decay $\sim 1/k^3$ for all values of Λ_C, which means that the total photon number density and the total energy density are both UV-divergent. If we turn from an instantaneous mass shift to a mass shift over a finite time interval, τ, which represents a more realistic scenario, the mentioned divergence in the loop integral is cured. Furthermore, one can infer from figure 3 that asymptotic photon spectra are highly sensitive to the order of differentiability of the mass parameterization, $m(t)$, considered.

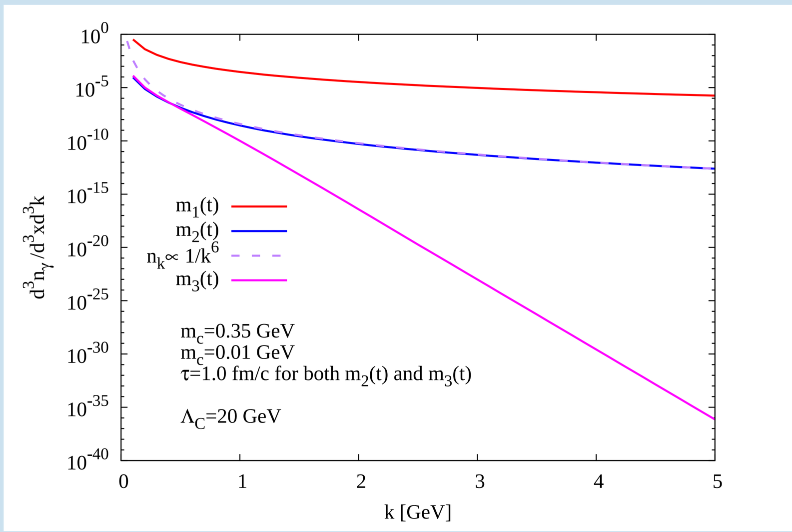

Figure 3: Asymptotic photon spectra for different mass parameterizations, $m(t)$.

If we turn from $m_1(t)$ to $m_2(t)$, which is continuously differentiable once, we see that the photon spectra are suppressed from $\sim 1/k^3$ to $\sim 1/k^6$ in the ultraviolet domain and are thus rendered UV-finite. If we moreover turn from $m_2(t)$ to $m_3(t)$, which is continuously differentiable arbitrarily many times, the photon spectra are suppressed even further to an exponential decay in the ultraviolet domain. We hence encounter a very similar sensitivity as for the asymptotic particle spectra investigated in [1, 2, 4].

Figures 4 and 5 furthermore show that the suppression of the photon numbers with respect to the instantaneous case is the stronger the more slowly the mass shift is assumed to take place.

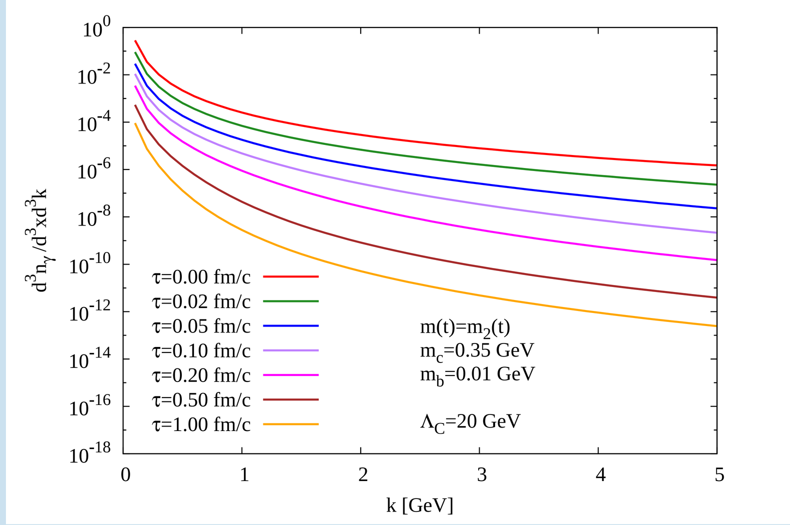

Figure 4: Asymptotic photon spectra for $m_2(t)$ and different change durations, τ.

As it must be, both $m_2(t)$ and $m_3(t)$ reproduce the photon spectrum for an instantaneous mass shift in the limit $\tau \to 0$.

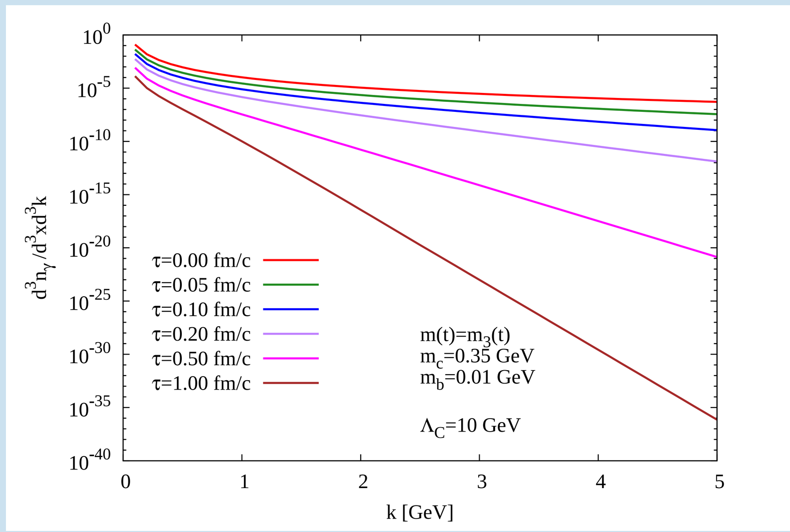

Figure 5: Asymptotic photon spectra for $m_3(t)$ and different change durations, τ.

To summarize, our investigations have shown that our approach on chiral photon production can also get the UV-behavior of the resulting photon spectra under control for physical mass parameterizations, $m(t)$.

Summary and Outlook

We have presented an ansatz for the description of chiral photon production, which eliminates possible unphysical vacuum contributions and renders the resulting photon spectra integrable in the ultraviolet domain.

As the quark pair production induced by the chiral mass shift contributes to the formation of the quark-gluon plasma (QGP) during a heavy-ion collision, our investigations are relevant in the context of finite lifetime effects on the photon emission from a QGP. It is of particular interest if the problems related with the vacuum contribution and/or the UV-behavior of the photon spectra within previous approaches [5, 6, 7] can be cured by extracting the photon numbers in the asymptotic limit. In this context, the role of the Ward-Takahashi identities, which are violated therein but conserved by (1) requires special consideration.

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