Photons production in heavy ion collisions as a signal of deconfinement phase

Aleksei Nikolskii

Bogoliubov Laboratory of Theoretical Physics, JINR, Dubna

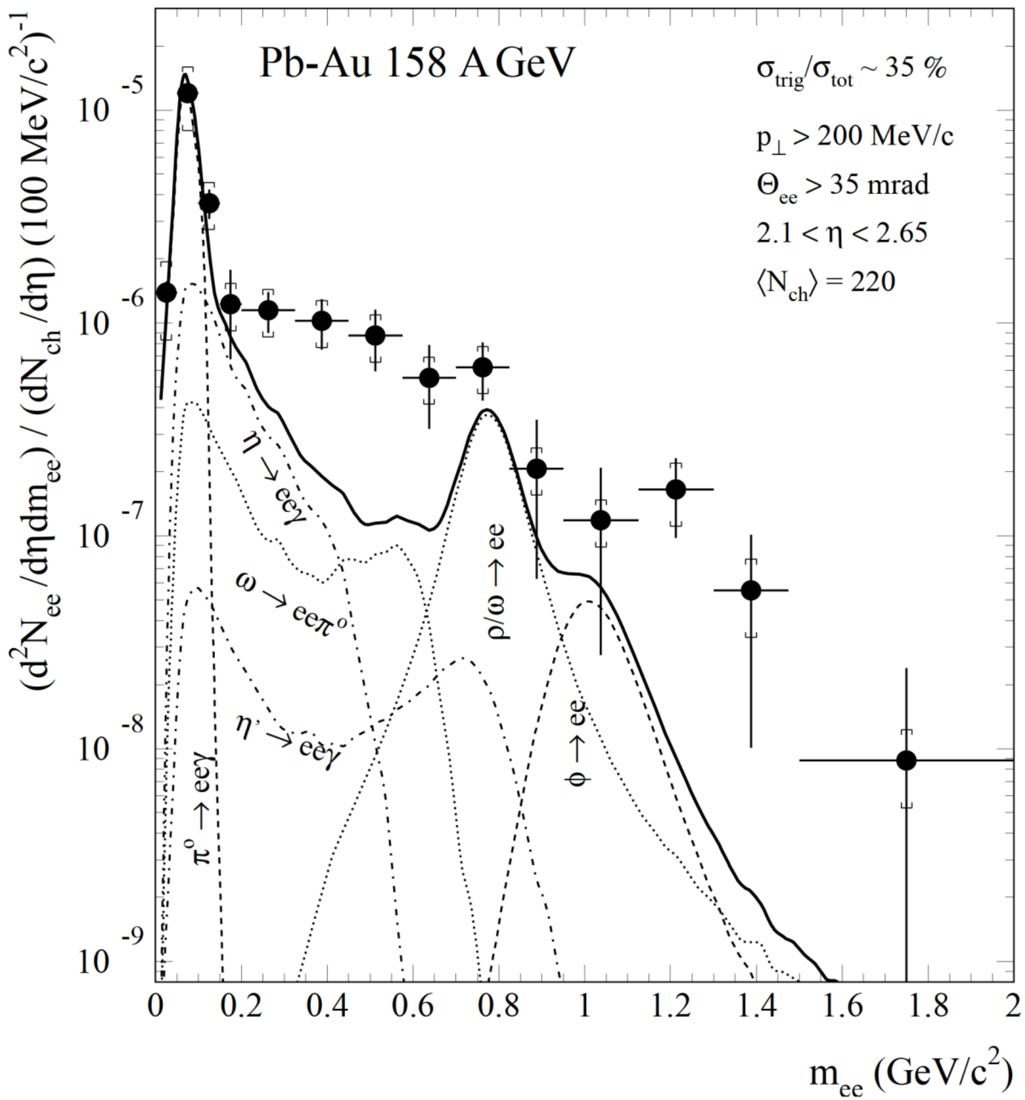
Based on arXiv: 2208.00842 [hep-ph]. Sergei Nedelko, Aleksei Nikolskii

The XXVI International Scientific Conference of Young Scientists and Specialists,
October 24-28, 2022

Motivation

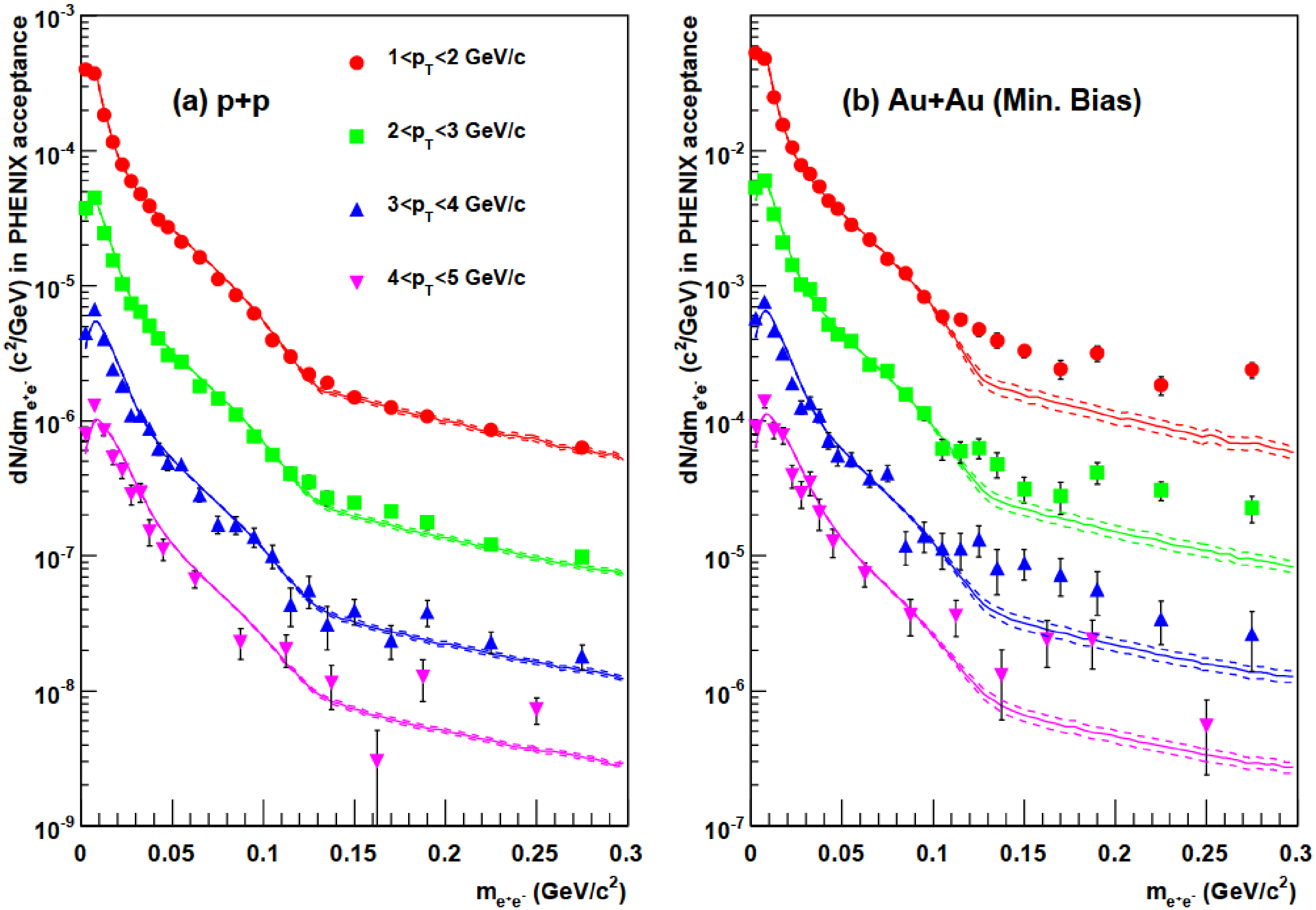
The photon excess in heavy ion collisions was first observed at CERN. **QGP- ?**

Special seminar: 10 February, 2000. CERES, Phys. Lett. B 422, 405, 1998



Motivation

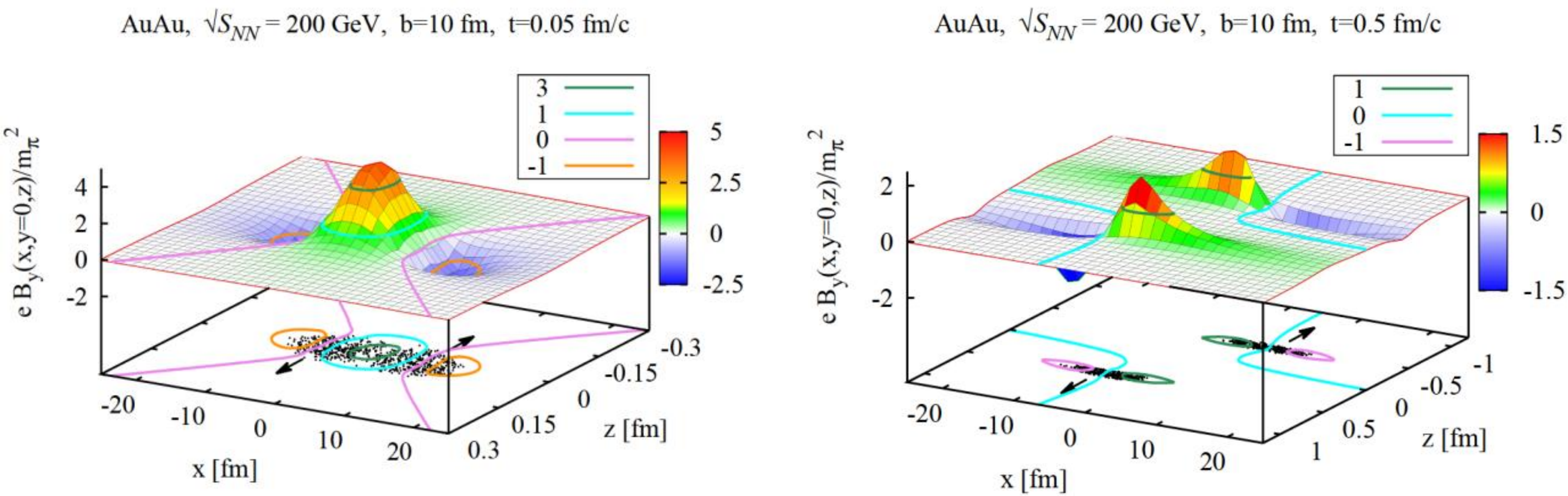
PHENIX Collaboration, Phys. Rev. Lett. 104, 132301, 2010



direct photon flow puzzle

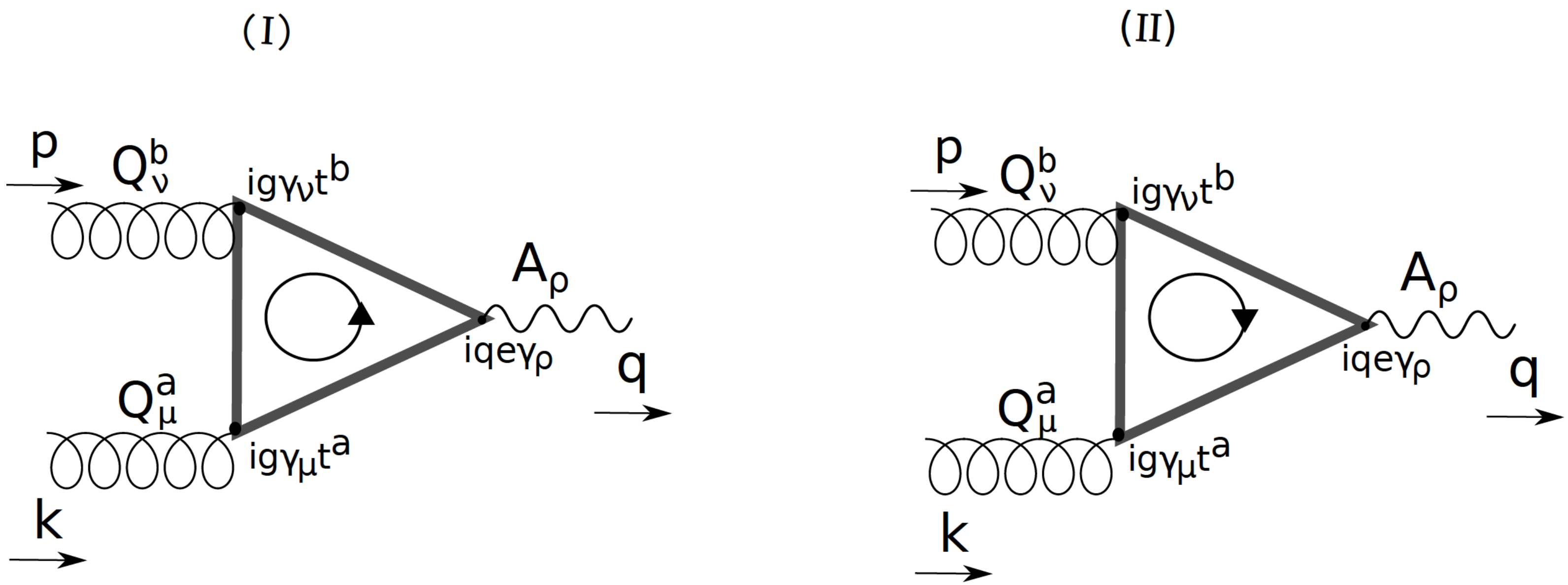
Motivation

- V. Skokov, A. Y. Illarionov, and V. Toneev, Int. J. Mod. Phys. A 24, 5925 (2009).
- V. Voronyuk, V. D. Toneev *et al.* Phys. Rev. C 83, 054911 (2011):



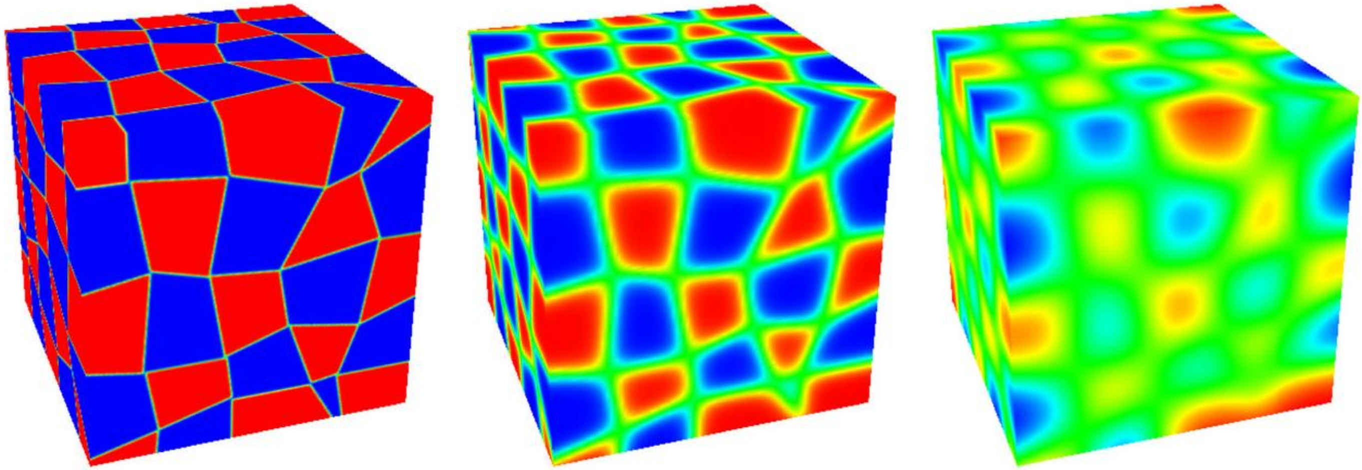
- A. Bzdak and V. Skokov, Anisotropy of photon production: initial eccentricity or magnetic field, Phys. Rev. Lett. 110, 192301 (2013);
- K. Tuchin, Particle production in strong electromagnetic fields in relativistic heavy-ion collisions, Adv. High Energy Phys. 2013, 490495 (2013);
- A. Ayala *et al.* Prompt photon yield and elliptic flow from gluon fusion induced by magnetic fields in relativistic heavy-ion collisions, Phys. Rev. D 96, 014023 (2017);

Investigated process of photon production

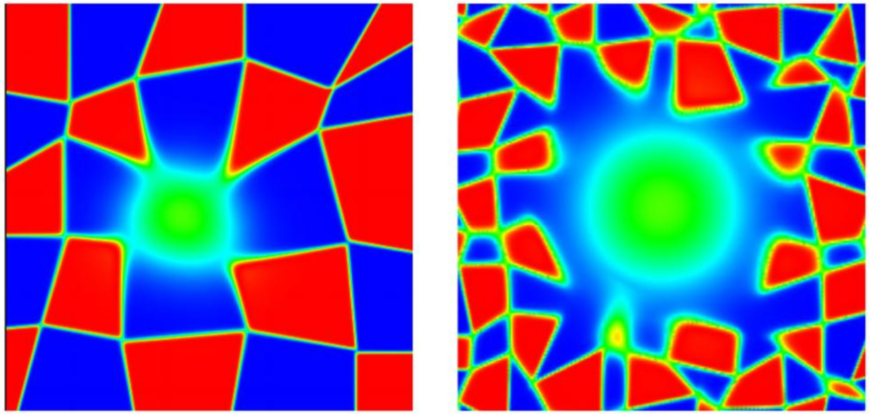


The diagrams of process $gg \to \gamma$ ***via* triangle quark loop** in the presence of homogeneous Abelian gauge field. *p, k* are the gluons momenta, *q* is the photon momentum. The arrows inside loop indicate the direction of loop momentum.

Domain model of QCD vacuum and hadronization



The blue and red areas correspond to confining almost everywhere homogeneous Abelian (anti-)self-dual gluon field.
Confinement.



The green areas correspond to the chromomagnetic field , in which quasiparticles with the color charge can be excited.
Deconfinement.

The strong magnetic field generated in relativistic heavy ion collisions is a catalyst for deconfinement.

- S. N. Nedelko and V. E. Voronin, Eur. Phys. J. A51, 45 (2015);
- B. V. Galilo and S. N. Nedelko, Phys. Rev. D 84, 094017 (2011).

The confinement phase

Let us consider the process $gg \to \gamma$ (*via* quark loop) in the presence of a random ensemble of almost everywhere homogeneous Abelian (anti-)self-dual gluon field:

$$\hat{B}_{\mu} = \frac{1}{2}\hat{B}_{\mu\nu}x_{\nu},\ \hat{B}_{\mu\nu} = \hat{n}B_{\mu\nu},\ \hat{n} = t^{8},\ \tilde{B}_{\mu\nu} = \frac{1}{2}\varepsilon_{\mu\nu\alpha\beta}B_{\alpha\beta} = \pm B_{\alpha\beta},\ \hat{B}_{\rho\mu}\hat{B}_{\rho\nu} = 4v^{2}B^{2}\delta_{\mu\nu},$$

$$\hat{f}_{\alpha\beta} = \frac{\hat{n}}{2vB}B_{\alpha\beta},\ v = \mathrm{diag}\left(\frac{1}{6},\frac{1}{6},\frac{1}{3}\right),\ \hat{f}^{ik}_{\mu\nu}\hat{f}^{kj}_{\nu\alpha} = \delta^{ij}\delta_{\mu\nu},$$

where field strength B sets the scale related to the value of the scalar gluon condensate.

The propagator of the quark* with mass m_f in the vacuum field presence has the form

$$S_f(x,y) = \exp\left(\frac{i}{2}x_{\mu}\hat{B}_{\mu\nu}y_{\nu}\right)H_f(x-y), \qquad (1)$$

The propagator is an entire function of p^2.
Quarks do not exist as particles.

$$H_f(z) = \frac{vB}{8\pi^2}\int_0^1 \frac{ds}{s^2}\exp\left(-\frac{vB}{2s}z^2\right)\left(\frac{1-s}{1+s}\right)^{\frac{m_f^2}{4vB}}$$

$$\times\left[-i\frac{vB}{s}z_{\mu}\left(\gamma_{\mu} \pm is\,\hat{f}_{\mu\nu}\gamma_{\nu}\gamma_5\right) + m_f\left(P_{\pm} + \frac{1+s^2}{1-s^2}P_{\mp} + \frac{i}{2}\gamma_{\mu}\hat{f}_{\mu\nu}\gamma_{\nu}\frac{s}{1-s^2}\right)\right],$$

* B. V. Galilo and S. N. Nedelko, Impact of the strong electromagnetic field on the QCD effective potential for homogeneous Abelian gluon field configurations, Phys. Rev. D 84, 094017 (2011), arXiv:1107.4737 [hep-ph].

The confinement phase

The terms with **odd** powers $f_{\mu\nu}$ **violate** the conditions of the Furry theorem

$$M^{(I)} = ieg^2 (2\pi)^4 \delta^{(4)}(p+k-q)\left(\frac{vB}{8\pi^2}\right)\int_0^1\int_0^1\int_0^1 \frac{ds_1}{s_1^2}\frac{ds_2}{s_2^2}\frac{ds_3}{s_3^2}\frac{(-ivB)^3}{s_1 s_2 s_3}$$

$$\left(\frac{1-s_1}{1+s_1}\right)^{\frac{m_f^2}{4vB}}\left(\frac{1-s_2}{1+s_2}\right)^{\frac{m_f^2}{4vB}}\left(\frac{1-s_3}{1+s_3}\right)^{\frac{m_f^2}{4vB}}\int d^4x d^4y\, e^{-i(px-ky)}$$

$$\left\langle \mathrm{Tr}\left[e^{ivBx^\mu \hat{f}_{\mu\nu} y^\nu - \frac{v}{2s_1}x^2 - \frac{v}{2s_2}y^2 - \frac{v}{2s_3}(y-x)^2}\ \hat{f}_{\alpha\omega}\hat{f}_{\beta\chi}\hat{f}_{\lambda\eta}\mathrm{K}^{\mu\nu\rho}_{\alpha\omega\beta\chi\lambda\eta} + \hat{f}_{\alpha\eta}\hat{f}_{\beta\omega}\Pi^{\mu\nu\rho}_{\alpha\eta\beta\omega} + \hat{f}_{\alpha\omega}\Gamma^{\mu\nu\rho}_{\alpha\omega}\right]\right\rangle \epsilon^a_\mu(k)\epsilon^b_\nu(p)\epsilon_\rho(q),$$

The amplitudes differ by the sign of the phase factor

$$M^{(II)} = ieg^2 (2\pi)^4 \delta^{(4)}(p+k-q)\left(\frac{vB}{8\pi^2}\right)\int_0^1\int_0^1\int_0^1 \frac{ds_1}{s_1^2}\frac{ds_2}{s_2^2}\frac{ds_3}{s_3^2}\frac{(-ivB)^3}{s_1 s_2 s_3}$$

$$\left(\frac{1-s_1}{1+s_1}\right)^{\frac{m_f^2}{4vB}}\left(\frac{1-s_2}{1+s_2}\right)^{\frac{m_f^2}{4vB}}\left(\frac{1-s_3}{1+s_3}\right)^{\frac{m_f^2}{4vB}}\int d^4x d^4y\, e^{-i(px-ky)}$$

$$\left\langle \mathrm{Tr}\left[e^{-ivBx^\mu \hat{f}_{\mu\nu} y^\nu - \frac{v}{2s_1}(x-y)^2 - \frac{v}{2s_2}y^2 - \frac{v}{2s_3}x^2}\ \hat{f}_{\alpha\omega}\hat{f}_{\beta\chi}\hat{f}_{\lambda\eta}\mathrm{K}^{\mu\nu\rho}_{\alpha\omega\beta\chi\lambda\eta} - \hat{f}_{\alpha\eta}\hat{f}_{\beta\omega}\Pi^{\mu\nu\rho}_{\alpha\eta\beta\omega} + \hat{f}_{\alpha\omega}\Gamma^{\mu\nu\rho}_{\alpha\omega}\right]\right\rangle \epsilon^a_\mu(k)\epsilon^b_\nu(p)\epsilon_\rho(q).$$

The confinement phase

The sign of the phase factor is reflected in the result of averaging over the spacial orientation of the background field*

$$\left\langle \prod_{j=1}^{n} \hat{f}_{\alpha_j\beta_j} e^{\pm i f_{\mu\nu} J_{\mu\nu}} \right\rangle = \frac{(\pm 1)^n}{(2i)^n} \prod_{j=1}^{n} \frac{\partial}{\partial J_{\alpha_j\beta_j}} \frac{\sin\sqrt{2\left(J_{\mu\nu}J_{\mu\nu} \pm J_{\mu\nu}\tilde{J}_{\mu\nu}\right)}}{\sqrt{2\left(J_{\mu\nu}J_{\mu\nu} \pm J_{\mu\nu}\tilde{J}_{\mu\nu}\right)}}, \qquad (2)$$

and

$$\left\langle \prod_{j=1}^{n} \hat{f}_{\alpha_j\beta_j} e^{-if_{\mu\nu} J_{\mu\nu}} \right\rangle = (-1)^n \left\langle \prod_{j=1}^{n} \hat{f}_{\alpha_j\beta_j} e^{if_{\mu\nu} J_{\mu\nu}} \right\rangle.$$

The terms in $M = M^{(I)} + M^{(II)}$ with the product of an even number of the field tensor $f_{\mu\nu}$ cancel each other out identically just as in the case of the "usual" Furry theorem.
The terms with the product of an odd number – cancel each other upon averaging of the spacial orientation of the background field.

* S. N. Nedelko and V. E. Voronin, Influence of confining gluon configurations on the $P \to \gamma^* \gamma$ transition form factors, Phys. Rev. D95, 074038 (2017), arXiv:1612.02621 [hep-ph].

The deconfinement phase

The deconfinement phase is characterized by the presence of the chromomagnetic field with the singled direction*

$$\hat{B}_{\mu\nu} = \hat{n} B_{\mu\nu} = \hat{n} B \, f_{\mu\nu}, \;\; f_{12} = -f_{21} = 1.$$

The coordinates and momenta (in Euclidean space-time)

$$x_{\perp} = \left(x_1, x_2, 0, 0\right), \; x_{\parallel} = \left(0, 0, x_3, x_4\right),$$
$$p_{\perp} = \left(p_1, p_2, 0, 0\right), p_{\parallel} = \left(0, 0, p_3, p_4\right).$$

The quark propagator with mass m_f in the presence of an external chromomagnetic field takes the form

$$S\left(x, y\right) = \exp\left\{-\frac{i}{2} x_{\perp}^{\mu} \hat{B}_{\mu\nu} y_{\perp}^{\nu}\right\} H_f\left(x - y\right), \qquad (3)$$

The propagator is complete, i.e. accounting for contribution of all Landau levels μ_n.

$$H_f\left(z\right) = \frac{B|\hat{n}|}{16\pi^2} \int_0^{\infty} \frac{ds}{s} \left(\coth\left(B|\hat{n}|s\right) - \sigma_{\rho\lambda} f_{\rho\lambda}\right) \exp\left\{-m_f^2 s - \frac{1}{4s} z_{\parallel}^2 - \frac{1}{8s}\left(B|\hat{n}|s \coth\left(B|\hat{n}|s\right) + 1\right) z_{\perp}^2\right\}$$
$$\left\{m_f - \frac{i}{2s}\gamma_{\mu} z_{\parallel}^2 - \frac{1}{2}\gamma_{\mu} \hat{B}_{\mu\nu} z_{\perp}^{\nu} - \frac{i}{4s}\left(B|\hat{n}|s \coth\left(B|\hat{n}|s\right) + 1\right)\gamma_{\mu} z_{\perp}^{\mu}\right\}, \sigma_{\rho\lambda} = \frac{i}{2}\left[\gamma_{\rho}, \gamma_{\lambda}\right].$$

* B. V. Galilo and S. N. Nedelko, Impact of the strong electromagnetic field on the QCD effective potential for homogeneous Abelian gluon field configurations, Phys. Rev. D 84, 094017 (2011), arXiv:1107.4737 [hep-ph].

The deconfinement phase

The amplitude of the diagrams (I) and (II) has the form

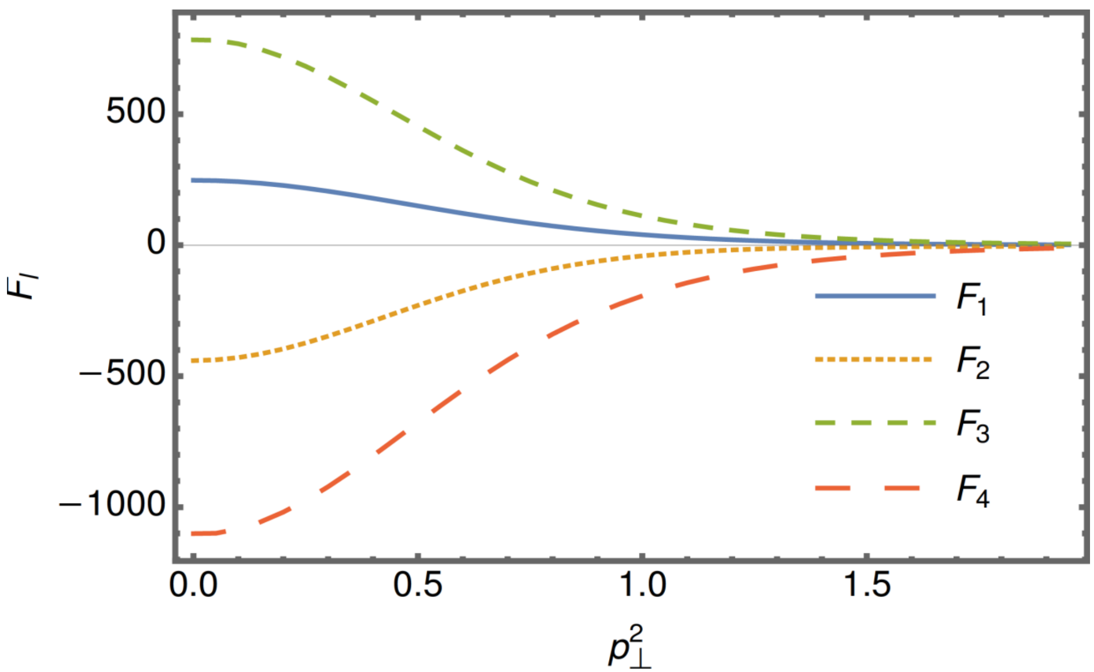
$$M = M^{(I)} + M^{(II)} = i(2\pi)^4 \delta^{(4)}(p+k-q) g^2 e \sum_l \mathcal{F}^l_{\mu\nu\rho}(p,k) F_l(p,k) \delta^{a8}\delta^{b8}\epsilon^a_\mu(k)\epsilon^b_\nu(p)\epsilon_\rho(q)$$

Form factors F_l have the structure

$$F_l(p,k) = \sum_f q_f \mathrm{Tr}_{\hat{n}} \int_0^\infty ds_1 ds_2 ds_3 \left[\psi_l^{(I)}(s_1,s_2,s_3 \,|\, \hat{n}, m_f) + \psi_l^{(II)}(s_1,s_2,s_3 \,|\, \hat{n}, m_f)\right]$$

$$\times \exp\left\{-p_\parallel^2 \phi_1(s_1,s_2,s_3) - p_\parallel k_\parallel \phi_2(s_1,s_2,s_3) - k_\parallel^2 \phi_3(s_1,s_2,s_3)\right.$$

$$\left. - p_\perp^2 \phi_4(s_1,s_2,s_3) - p_\perp k_\perp \phi_5(s_1,s_2,s_3) - k_\perp^2 \phi_6(s_1,s_2,s_3) - m_f^2(s_1+s_2+s_3)\right\} \qquad (4)$$



Some of the form factors $M = M^{(I)} + M^{(II)}$ as functions of transverse gluon momenta $p_\perp^2 = k_\perp^2$ for fixed longitudinal momenta $p_\parallel^2 = k_\parallel^2 = 1$ in Euclidean kinematics. Dimensionless notation $p^2 = p^2/B$ is used. Form factors are dimensionless.

arXiv: 2208.00842 [hep-ph]

The deconfinement phase

To calculate $T=|M|^2$ one has to continue representation to Minkowsky kinematics

$$p_{\|}^2 \to -p_{\|}^2,\ k_{\|}^2 \to -k_{\|}^2,\ p_{\|}k_{\|} \to -p_{\|}k_{\|}. \tag{5}$$

In Minkowski space-time, on-shell conditions for gluons and photon $p^2=0, k^2=0,$ $(p+k)^2=0$ impose the following relations

$$p_{\|}^2 = p_{\perp}^2,\ \ k_{\|}^2 = k_{\perp}^2,\ \ p_{\|}k_{\|} = p_{\perp}k_{\perp}.$$

The probability of photon production is given by the squared amplitude averaged over the initial gluon polarization states and summed over the final polarizations of photon

$$\bar{T}(p,k,q) = \Delta v \Delta \tau (2\pi)^4 \delta(p+k-q) T(p,k), \tag{6}$$

$$T(p,k) = \frac{2\alpha\alpha_s}{\pi} \int ds_1 ds_2 ds_3 dr_1 dr_2 dr_3\, F(s_1,s_2,s_3,r_1,r_2,r_3 \,|\, p,k)$$

$$\times \exp\left\{ p_{\perp}^2 \Phi_1(s_1,s_2,s_3,r_1,r_2,r_3) + p_{\perp}k_{\perp}\Phi_2(s_1,s_2,s_3,r_1,r_2,r_3) + \right.$$

$$\left. k_{\perp}^2 \Phi_3(s_1,s_2,s_3,r_1,r_2,r_3) - m_f^2(s_1+s_2+s_3+r_1+r_2+r_3) \right\},$$

The deconfinement phase

The functions

$$\begin{aligned}
\Phi_1 &= \phi_1(s_1,s_2,s_3)+\phi_1(r_1,r_2,r_3)-\phi_4(s_1,s_2,s_3)-\phi_4(r_1,r_2,r_3),\\
\Phi_2 &= \phi_2(s_1,s_2,s_3)+\phi_2(r_1,r_2,r_3)-\phi_5(s_1,s_2,s_3)-\phi_5(r_1,r_2,r_3),\\
\Phi_3 &= \phi_3(s_1,s_2,s_3)+\phi_3(r_1,r_2,r_3)-\phi_6(s_1,s_2,s_3)-\phi_6(r_1,r_2,r_3).
\end{aligned} \tag{7}$$

are positive in the whole region of integration and grow linearly for $s_j \longrightarrow \infty$, and the proper time integrals in Eq. (6) «converge» only for the limited range of gluon momenta p and k.

For the regime $p_\perp^2 = k_\perp^2$ the Integral (6) is converges if

$$\boxed{p_\perp^2 < \frac{3}{2} m_f^2,}$$

as it can be seen from Eqs. (4) и (7).

- V.O. Papanyan and V.I. Ritus, Three-photon interaction in an intense field and scaling invariance, Zh. Eksp. Teor. Fiz. 65, 1756 (1973).
- S.L. Adler, J.N. Bahcall, C.G. Callan, and M.N. Rosenbluth, Photon spliting in a strong magnetic field, Phys. Rev. Lett. 25, 1061 (1970).
- V.O. Papanyan and V.I. Ritus, Vacuum polarization and photon splitting in an intense fiield, Zh. Eksp. Teor. Fiz. 61, 2231 (1971).

The squared amplitude / massless quarks

- A. Ayala *et al.* Phys. Rev. D 96, 014023 (2017).
- A. Ayala *et al.* Eur. Phys. J. A 56, 53 (2020).

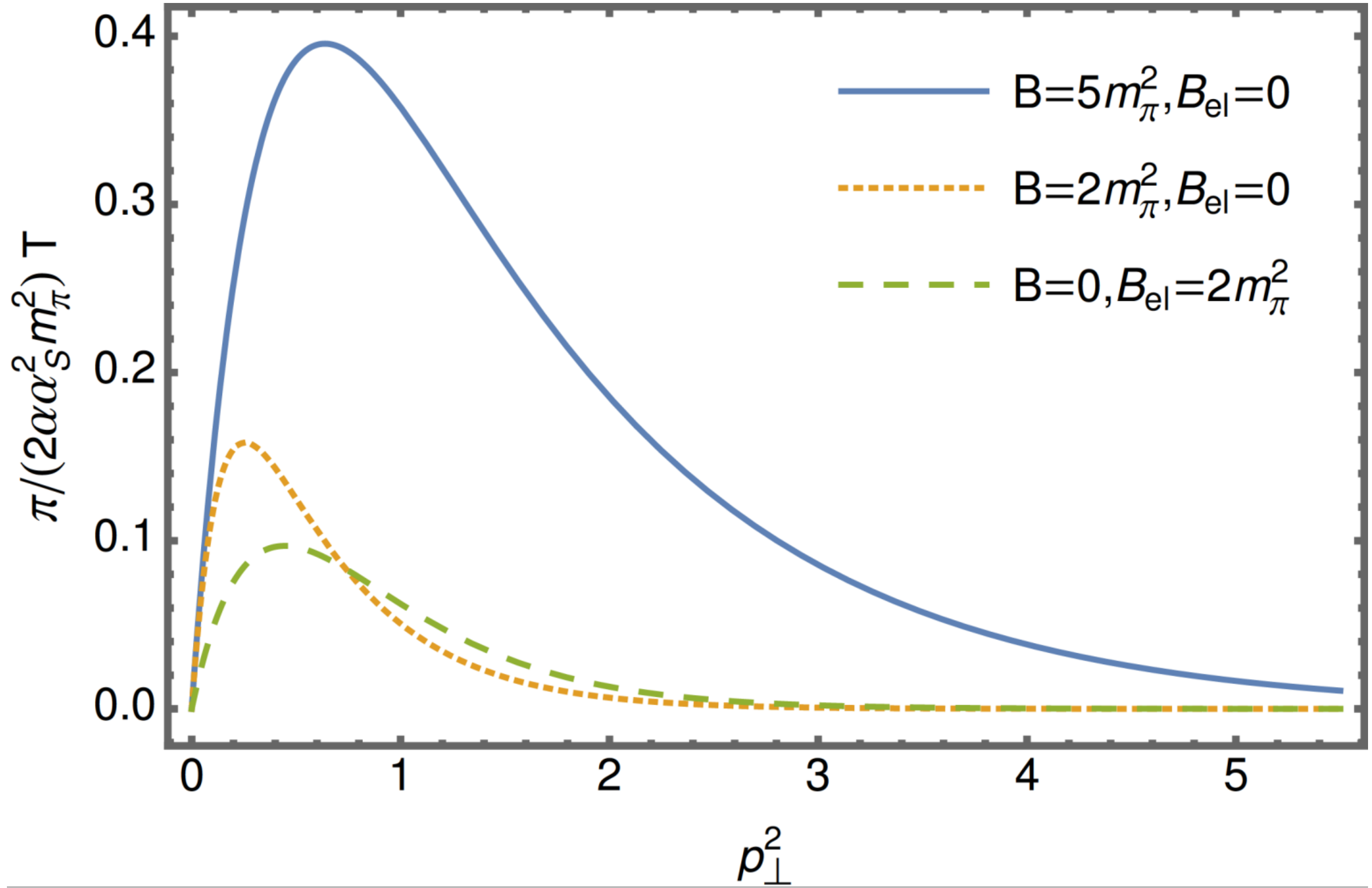
$$T(p,k)=\frac{2\alpha\alpha_s^2}{\pi}q_f^2\left(2p_\perp^2+k_\perp^2+p_\perp k_\perp\right)\exp\left\{-\frac{1}{|q_f B_{el}|}\left(p_\perp^2+k_\perp^2+p_\perp k_\perp\right)\right\}. \quad (8)$$

- the limit of strong field (pure magnetic): $eB_{el} \gg m_f^2$;
- massless quarks;
- the quark propagator – the lowest (LLL) and the first (1LL) Landau levels.

$$|q_f B_{el}| \rightarrow |q_f B_{el}+\hat{n}B|$$

$$T(p,k)=\frac{2\alpha\alpha_s^2}{N_c\pi}q_f^2\,\mathrm{Tr}_{\hat{n}}\left(2p_\perp^2+k_\perp^2+p_\perp k_\perp\right)\exp\left\{-\frac{1}{|q_f B_{el}+\hat{n}B|}\left(p_\perp^2+k_\perp^2+p_\perp k_\perp\right)\right\}. \quad (9)$$

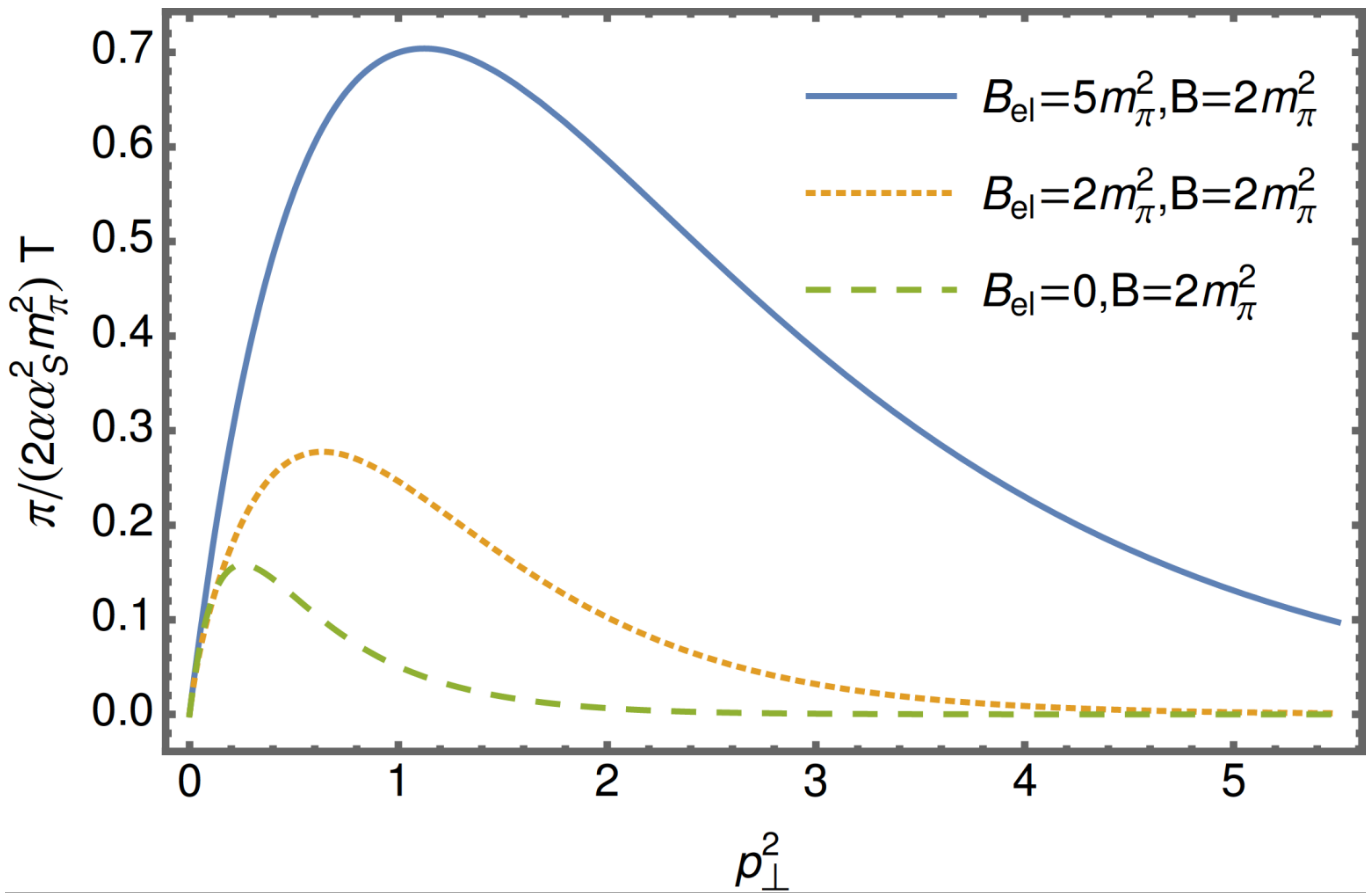
The squared amplitude / massless quarks



Dependence T(p,k) on the gluon momenta (9) in the regime $p_\perp^2 = k_\perp^2$.

The green long-dashed line corresponds to the purely magnetic field B_{el} , dotted and solid lines represent the case of pure chromomagnetic field B with different strength. The mass of the pion is chosen as the scale. Dimensionless notation $p_\perp^2 = p_\perp^2 / B$ is used.

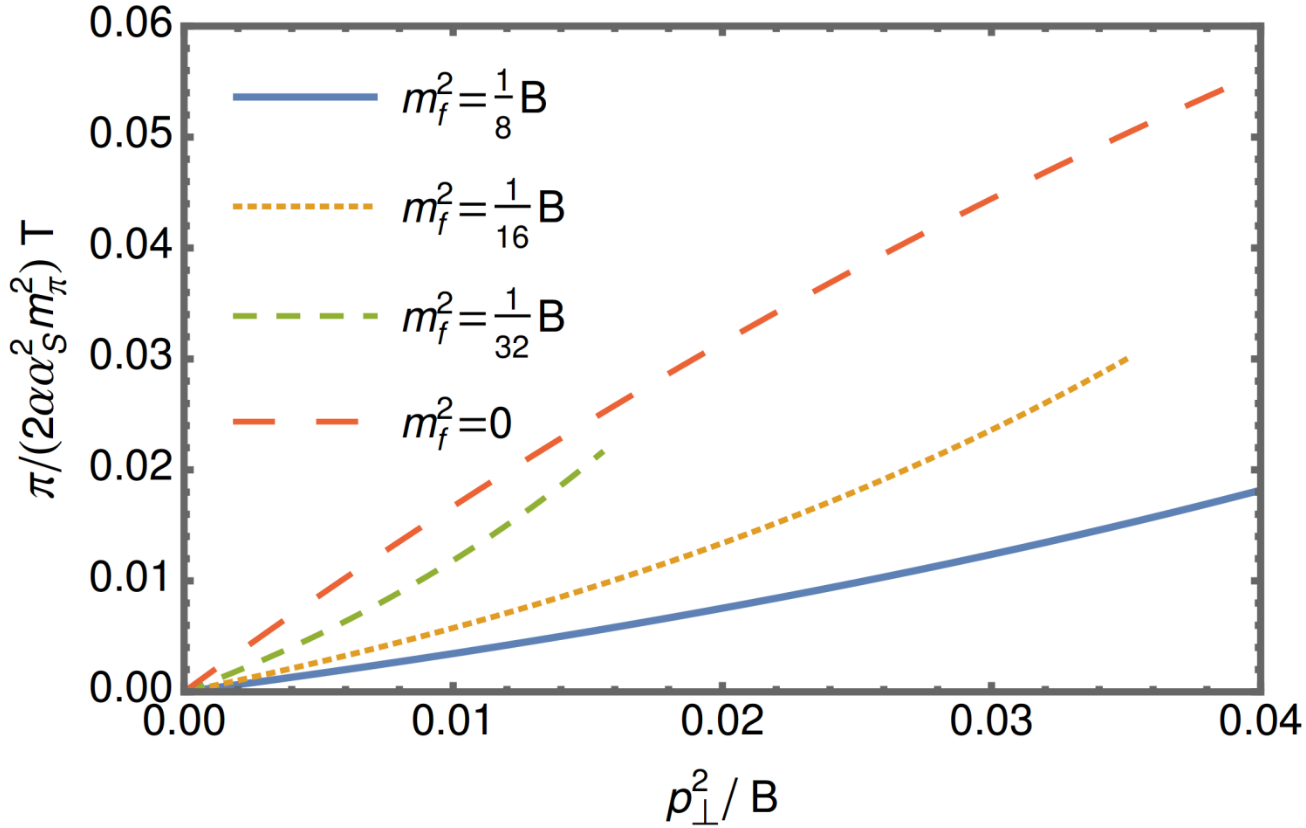
The amplitude squared / massless quarks



Dependence T(p,k) in the regime $p_\perp^2 = k_\perp^2$ on the different strength of magnetic field B_{el}.

The green long-dashed line corresponds to the chromomagnetic field $B=2m_\pi^2$ alone, and the dotted and solid curves represent the cases of both fields with different magnetic field strengths $B_{el}=2m_\pi^2$, $B_{el}=5m_\pi^2$ respectively. Dimensionless notation $p_\perp^2 = p_\perp^2 / B$ is used.

The comparison: massless and massive quarks



The comparison of the squared amplitude (9) obtained by the Landau level decomposition for a massless quark (red long-dashed line) with the squared amplitude (6) taking into account all Landau levels for different quark masses m_f at low gluon momenta in the regime $p_\perp^2 = k_\perp^2 < 3m_f^2/2.$

The chromomagnetic field strength $B = 4m_\pi^2$, and magnetic field $B_{el} = 0$. Dimensionless notation $p_\perp^2 = p_\perp^2 / B$ is used.

The photon spectrum / massless quarks

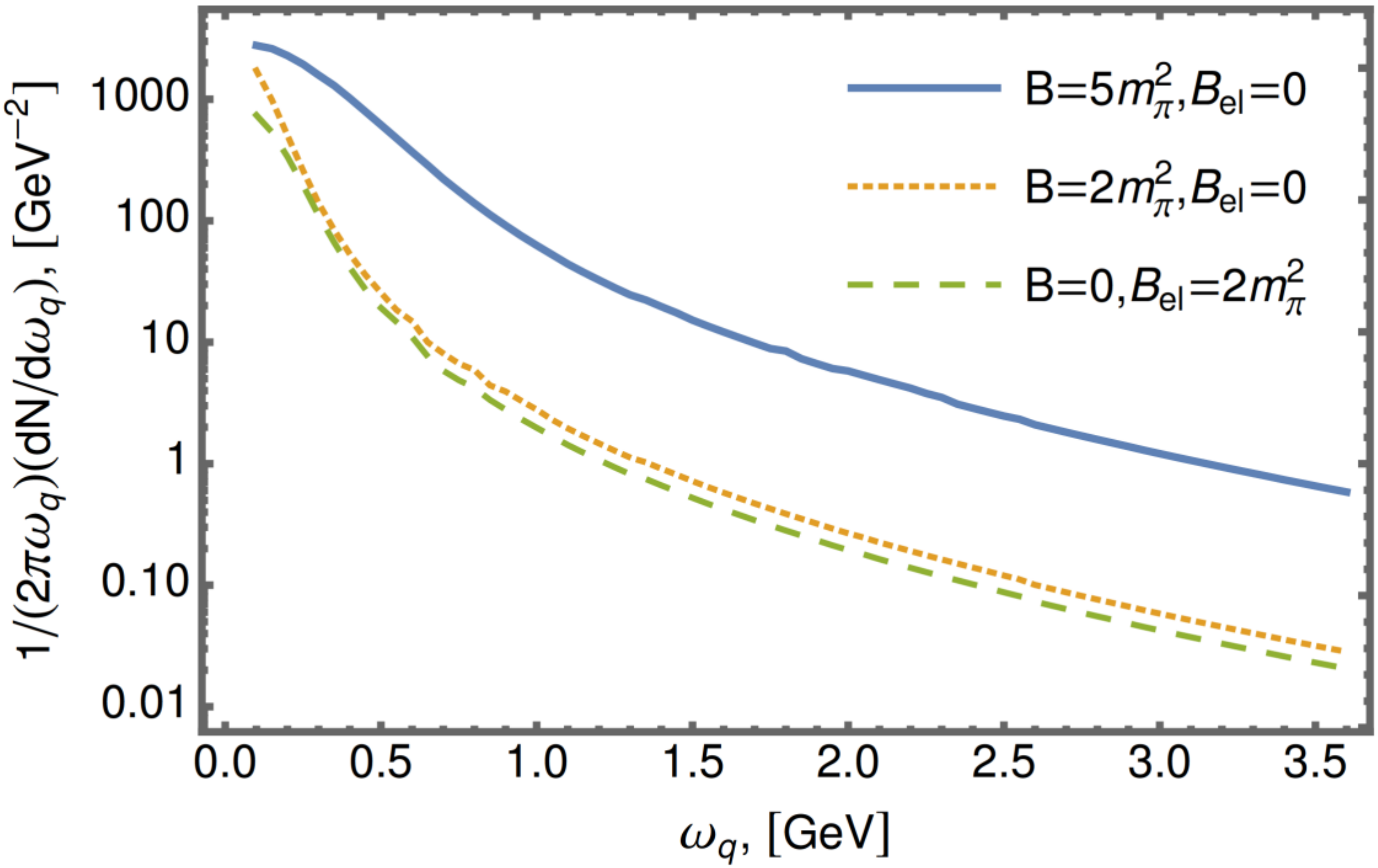
The invariant photon distribution in the presence of chromomagnetic and magnetic fields can be represented in the form

$$\frac{1}{2\pi\omega_q}\frac{dN}{d\omega_q} = \nu\Delta\tau\,\frac{\alpha\alpha_s^2\pi}{2N_c\left(2\pi\right)^6\omega_q}\,q_f^2\,\mathrm{Tr}_{\hat{n}}\int\limits_0^{\omega_q} d\omega_p\left(2\omega_p^2+\omega_q^2-\omega_p\omega_q\right)e^{-g_f^B\left(\omega_p,\omega_q\right)}$$
$$\times\left[I_0\left(g_f^B\left(\omega_p,\omega_q\right)\right)-I_1\left(g_f^B\left(\omega_p,\omega_q\right)\right)\right]\left(n\left(\omega_p\right)n\left(\left|\omega_q-\omega_p\right|\right)\right), \tag{10}$$

$I_0\left(g_f^B\left(\omega_p,\omega_q\right)\right)$, $I_1\left(g_f^B\left(\omega_p,\omega_q\right)\right)$ are the modified Bessel function of the first kind,

$$g_f^B\left(\omega_p,\omega_q\right)=\frac{\omega_p^2+\omega_q^2-\omega_p\omega_q}{2\left|\hat{n}B+q_fB_{el}\right|},\ \ n\left(\omega\right)=\frac{\eta}{e^{\omega/\Lambda_s}-1}.$$

The photon spectrum / massless quarks



Differential energy distribution (10) of the generated photons for a pure magnetic field B_{el} (green long-dashed line) and pure chromomagnetic field B (the dotted and solid curves). The pion mass $m_\pi = 0.135\ \mathrm{GeV}$ is chosen as the scale.

Conclusions

- Photon production in the process $gg \rightarrow \gamma$ (*via* triangle quark loop) may serve as a signal of the transition between of the confinement – deconfinement phases in heavy ion collisions;

- The confinement phase: the conversion probability of two gluons into a photon vanishes due to the random nature of the statistical ensemble of confining vacuum fields;

- The deconfinement phase: the «short-living» strong magnetic field with singled direction is generated by relativistic heavy ion collisions and plays the role of a catalyst for the deconfinement phase transition. This transition is accompanied by the appearance of the «long-living» chromomagnetic field with the same direction as the magnetic field – as a consequence, the conditions of Furry's theorem are not satisfied – the conversion probability of two gluons into a photon is nonzero;

- The case of massless quarks – the chromomagnetic field leads to an increase in the amplitude of photon production due to a longer lifetime than the magnetic one, and the photon signal is increased at the small gluon momenta in comparison with the pure magnetic field with the same strength;

- The case of massive quarks – the contribution of strange quark is important. In this case, the expansion of the quark propagator in Landau levels is not applicable;

- The generated photons have a strong angular anisotropy because the production amplitude is «tied» to the singled direction of the magnetic and chromomagnetic fields. *Direct photon flow puzzle.*

Thanks for your attention!