

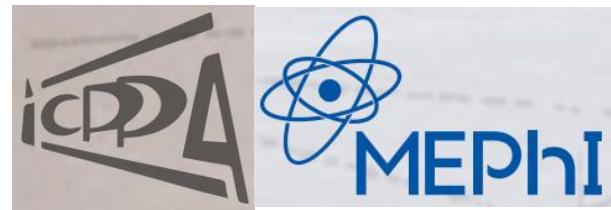
Photons as a signal of deconfinement in hadronic matter under extreme conditions

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

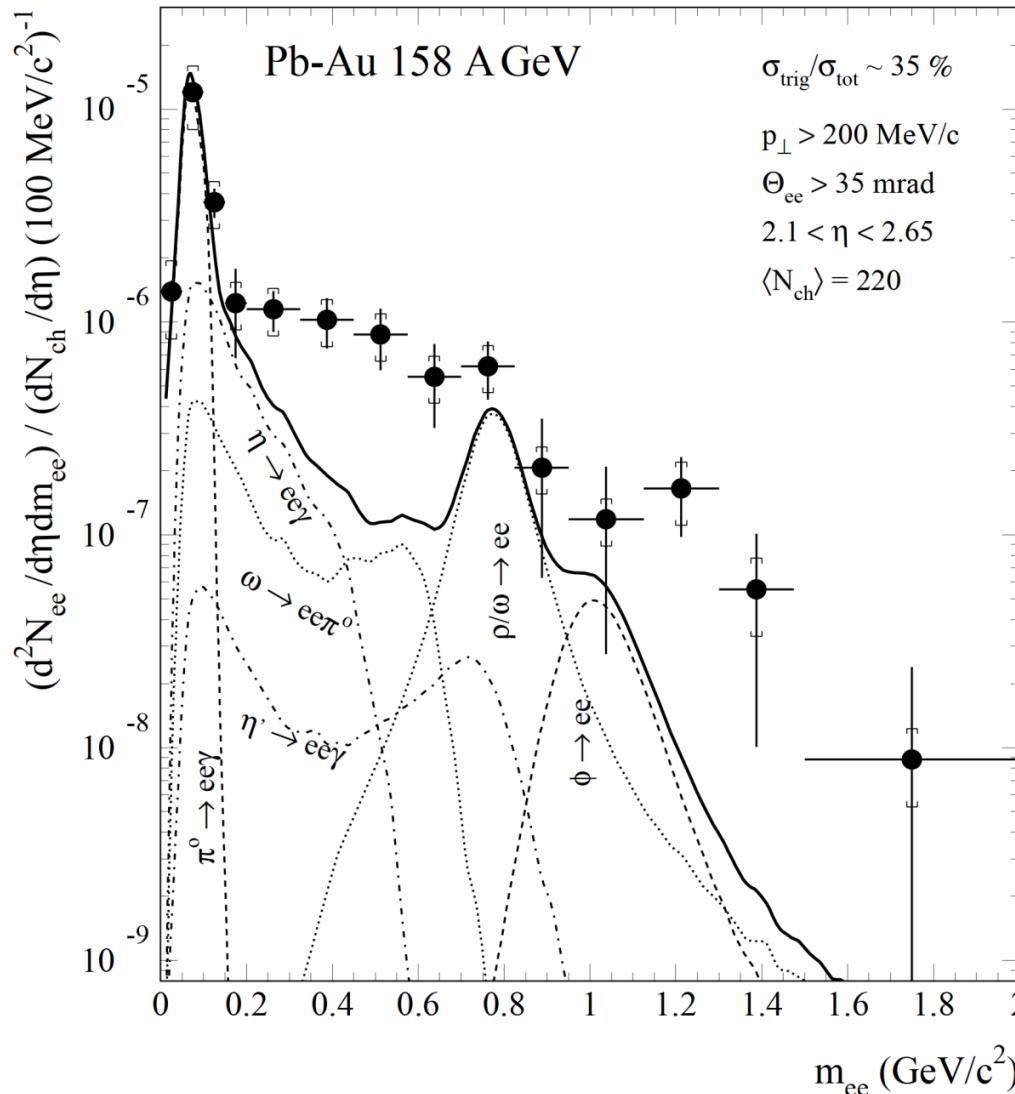
The 6th international conference on particle physics and astrophysics,
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Motivation

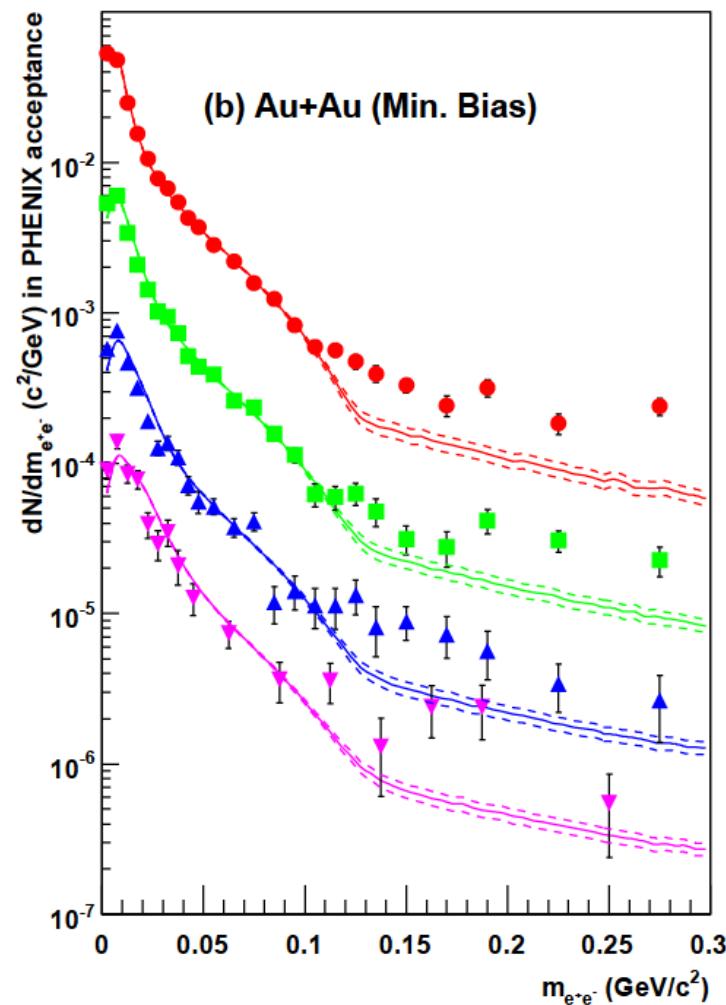
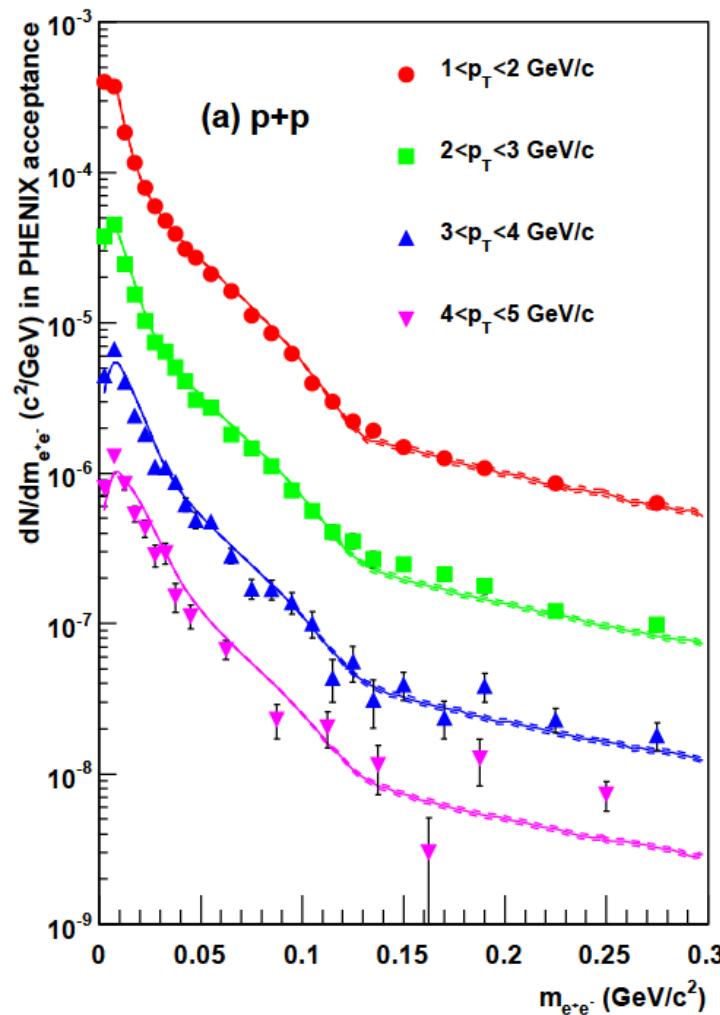
The photon excess in heavy ion collisions was first observed at CERN. **QGP- ?**

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



Motivation

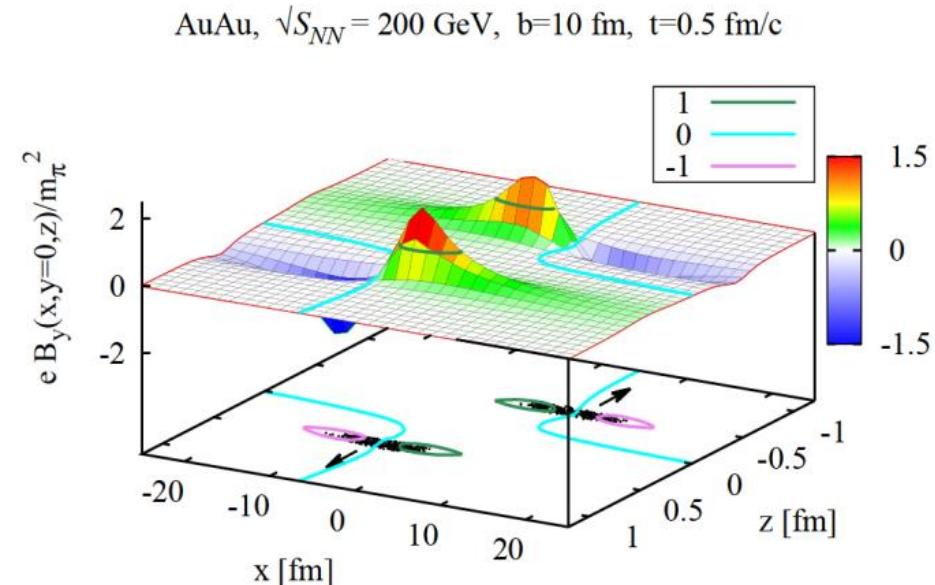
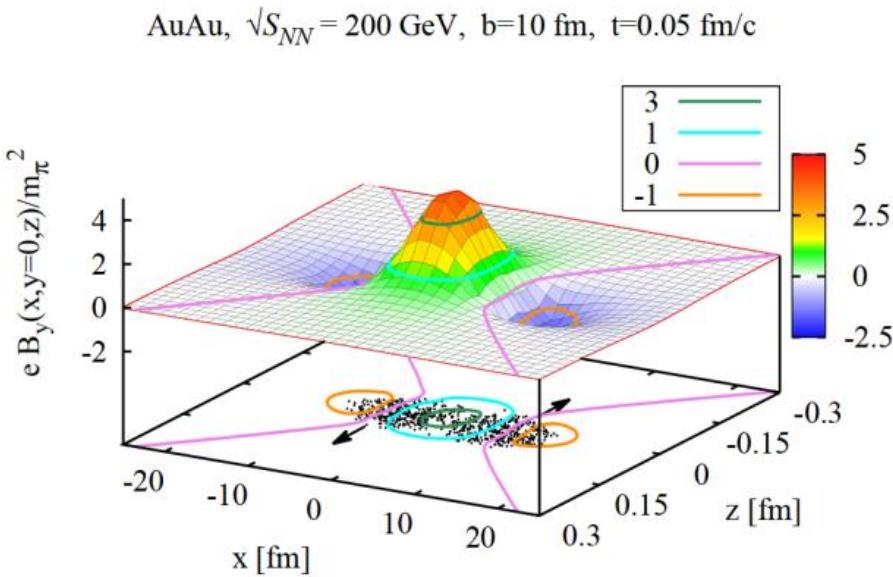
PHENIX Collaboration, Phys. Rev. Lett. 104, 132301, 2010



direct photon flow puzzle

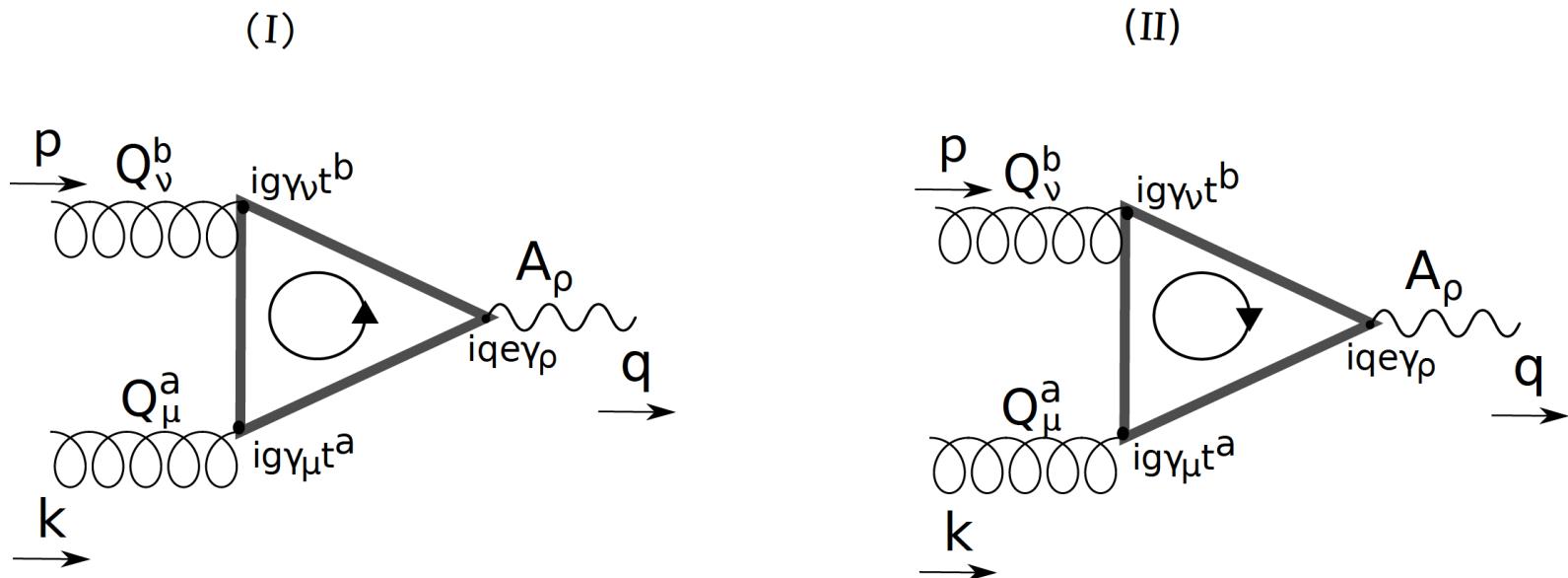
Motivation

- 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):**



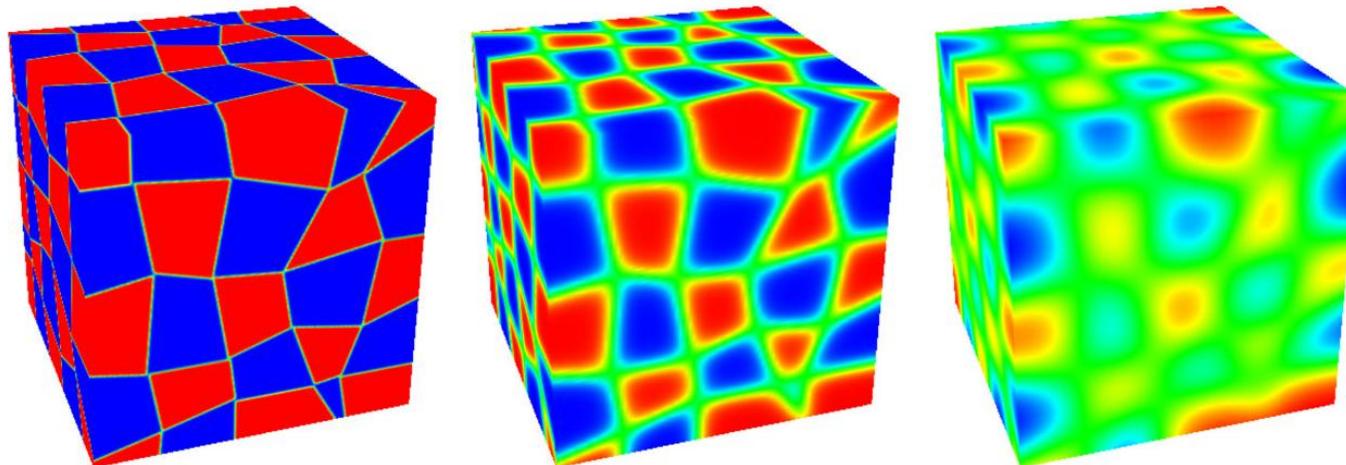
- A. Bzdak and V. Skokov, Anisotropy of photon production: initial eccentricity or magnetic field, Phys. Rev. Lett. 110, 192301 (2013);
- K. Tuchin, Particle production in strong electromagnetic fields in relativistic heavy-ion collisions, Adv. High Energy Phys. 2013, 490495 (2013);
- A. Ayala *et al.* 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

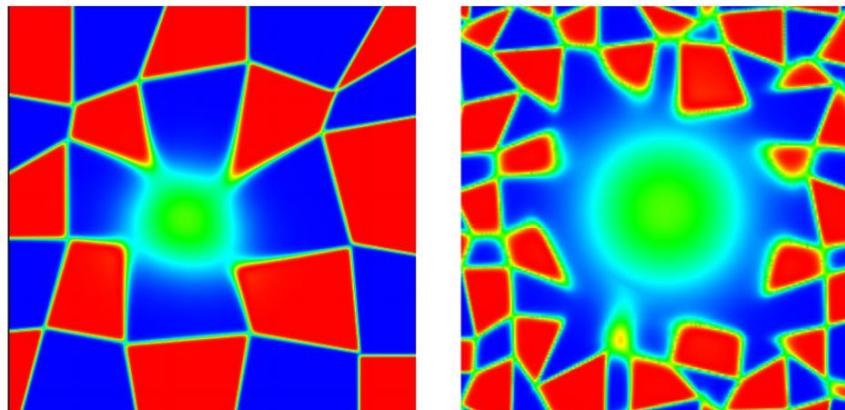


The diagrams of process $gg \rightarrow \gamma$ via triangle quark loop in the presence of homogeneous Abelian gauge field. p, k are the gluons momenta, q is the photon momentum. The arrows inside loop indicate the direction of loop momentum.

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 $gg \rightarrow \gamma$ (via quark loop) in the presence of a random ensemble of almost everywhere homogeneous Abelian (anti-)self-dual gluon field:

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

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

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

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

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

$$H_f(z) = \frac{vB}{8\pi^2} \int_0^1 \frac{ds}{s^2} \exp\left(-\frac{vB}{2s} z^2\right) \left(\frac{1-s}{1+s}\right)^{\frac{m_f^2}{4vB}} \\ \times \left[-i \frac{vB}{s} z_\mu \left(\gamma_\mu \pm i s \hat{f}_{\mu\nu} \gamma_\nu \gamma_5 \right) + m_f \left(P_\pm + \frac{1+s^2}{1-s^2} P_\mp + \frac{i}{2} \gamma_\mu \hat{f}_{\mu\nu} \gamma_\nu \frac{s}{1-s^2} \right) \right],$$

* 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

$$M^{(I)} = ie g^2 (2\pi)^4 \delta^{(4)}(p+k-q) \left(\frac{vB}{8\pi^2} \right) \int_0^1 \int_0^1 \int_0^1 \frac{ds_1}{s_1^2} \frac{ds_2}{s_2^2} \frac{ds_3}{s_3^2} \frac{(-ivB)^3}{s_1 s_2 s_3} \\ \left(\frac{1-s_1}{1+s_1} \right)^{\frac{m_f^2}{4vB}} \left(\frac{1-s_2}{1+s_2} \right)^{\frac{m_f^2}{4vB}} \left(\frac{1-s_3}{1+s_3} \right)^{\frac{m_f^2}{4vB}} \int d^4x d^4y e^{-i(px-ky)} \\ \left\langle \text{Tr} \left[e^{ivBx^\mu \hat{f}_{\mu\nu} y^\nu - \frac{v}{2s_1}x^2 - \frac{v}{2s_2}y^2 - \frac{v}{2s_3}(y-x)^2} \hat{f}_{\alpha\omega} \hat{f}_{\beta\chi} \hat{f}_{\lambda\eta} K_{\alpha\omega\beta\chi\lambda\eta}^{\mu\nu\rho} + \hat{f}_{\alpha\eta} \hat{f}_{\beta\omega} \Pi_{\alpha\eta\beta\omega}^{\mu\nu\rho} + \hat{f}_{\alpha\omega} \Gamma_{\alpha\omega}^{\mu\nu\rho} \right] \right\rangle \epsilon_\mu^a(k) \epsilon_\nu^b(p) \epsilon_\rho(q),$$

The amplitudes differ by the **sign** of the phase factor

$$M^{(II)} = ie g^2 (2\pi)^4 \delta^{(4)}(p+k-q) \left(\frac{vB}{8\pi^2} \right) \int_0^1 \int_0^1 \int_0^1 \frac{ds_1}{s_1^2} \frac{ds_2}{s_2^2} \frac{ds_3}{s_3^2} \frac{(-ivB)^3}{s_1 s_2 s_3} \\ \left(\frac{1-s_1}{1+s_1} \right)^{\frac{m_f^2}{4vB}} \left(\frac{1-s_2}{1+s_2} \right)^{\frac{m_f^2}{4vB}} \left(\frac{1-s_3}{1+s_3} \right)^{\frac{m_f^2}{4vB}} \int d^4x d^4y e^{-i(px-ky)} \\ \left\langle \text{Tr} \left[e^{-ivBx^\mu \hat{f}_{\mu\nu} y^\nu - \frac{v}{2s_1}(x-y)^2 - \frac{v}{2s_2}y^2 - \frac{v}{2s_3}x^2} \hat{f}_{\alpha\omega} \hat{f}_{\beta\chi} \hat{f}_{\lambda\eta} K_{\alpha\omega\beta\chi\lambda\eta}^{\mu\nu\rho} - \hat{f}_{\alpha\eta} \hat{f}_{\beta\omega} \Pi_{\alpha\eta\beta\omega}^{\mu\nu\rho} + \hat{f}_{\alpha\omega} \Gamma_{\alpha\omega}^{\mu\nu\rho} \right] \right\rangle \epsilon_\mu^a(k) \epsilon_\nu^b(p) \epsilon_\rho(q).$$

The confinement phase

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

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

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 $P \rightarrow \gamma^* \gamma$ 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*

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

The coordinates and momenta (in Euclidean space-time)

$$x_\perp = (x_1, x_2, 0, 0), \quad x_\parallel = (0, 0, x_3, x_4), \\ p_\perp = (p_1, p_2, 0, 0), \quad p_\parallel = (0, 0, p_3, p_4).$$

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

$$S(x, y) = \exp \left\{ -\frac{i}{2} x_\perp^\mu \hat{B}_{\mu\nu} y_\perp^\nu \right\} H_f(x - y), \quad \text{The propagator is complete, i.e. accounting for contribution of all Landau levels } \mu_n. \quad (3)$$

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

* 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].

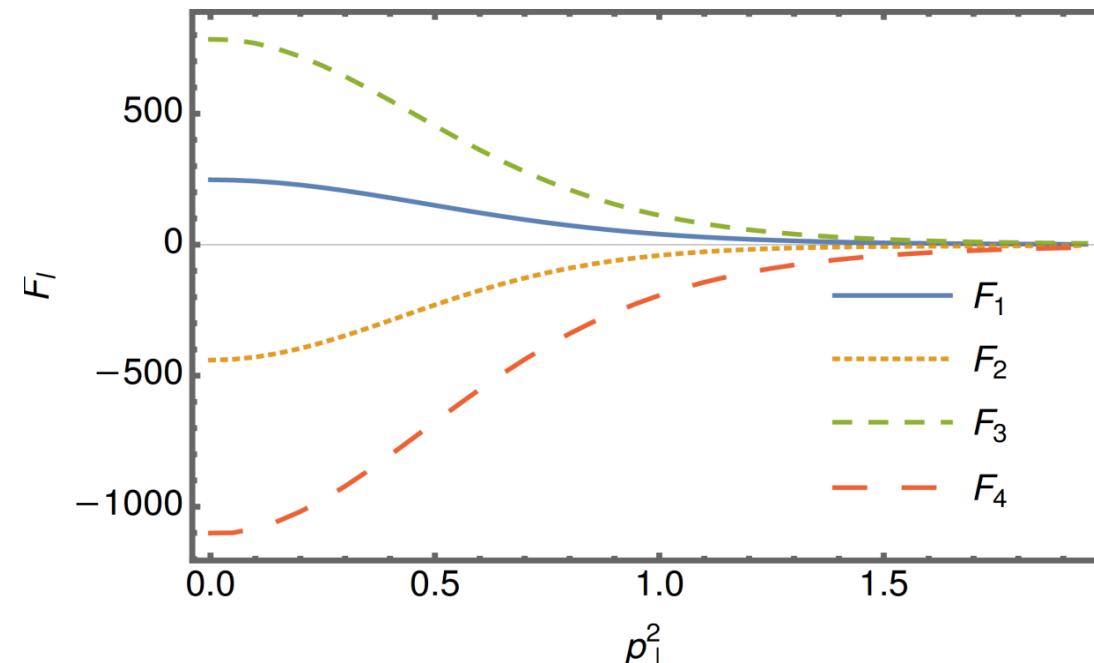
The deconfinement phase

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

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

Form factors F_l have the structure

$$\begin{aligned} F_l(p, k) = & \sum_f q_f \text{Tr}_{\hat{n}} \int_0^\infty ds_1 ds_2 ds_3 \left[\psi_l^{(I)}(s_1, s_2, s_3 | \hat{n}, m_f) + \psi_l^{(II)}(s_1, s_2, s_3 | \hat{n}, m_f) \right] \\ & \times \exp \left\{ -p_{\parallel}^2 \phi_1(s_1, s_2, s_3) - p_{\parallel} k_{\parallel} \phi_2(s_1, s_2, s_3) - k_{\parallel}^2 \phi_3(s_1, s_2, s_3) \right. \\ & \left. - p_{\perp}^2 \phi_4(s_1, s_2, s_3) - p_{\perp} k_{\perp} \phi_5(s_1, s_2, s_3) - k_{\perp}^2 \phi_6(s_1, s_2, s_3) - m_f^2(s_1 + s_2 + s_3) \right\} \end{aligned} \quad (4)$$



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.

The deconfinement phase

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

$$p_{\parallel}^2 \rightarrow -p_{\parallel}^2, k_{\parallel}^2 \rightarrow -k_{\parallel}^2, p_{\parallel}k_{\parallel} \rightarrow -p_{\parallel}k_{\parallel}. \quad (5)$$

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

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

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

$$\begin{aligned} \bar{T}(p, k, q) &= \Delta v \Delta \tau (2\pi)^4 \delta(p + k - q) T(p, k), \quad (6) \\ T(p, k) &= \frac{2\alpha\alpha_s}{\pi} \int ds_1 ds_2 ds_3 dr_1 dr_2 dr_3 F(s_1, s_2, s_3, r_1, r_2, r_3 | p, k) \\ &\times \exp \left\{ p_{\perp}^2 \Phi_1(s_1, s_2, s_3, r_1, r_2, r_3) + p_{\perp}k_{\perp} \Phi_2(s_1, s_2, s_3, r_1, r_2, r_3) + \right. \\ &\quad \left. k_{\perp}^2 \Phi_3(s_1, s_2, s_3, r_1, r_2, r_3) - m_f^2 (s_1 + s_2 + s_3 + r_1 + r_2 + r_3) \right\}, \end{aligned}$$

The deconfinement phase

The functions

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

are positive in the whole region of integration and grow linearly for $s_j \rightarrow \infty$, 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

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

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 splitting 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 field, 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).

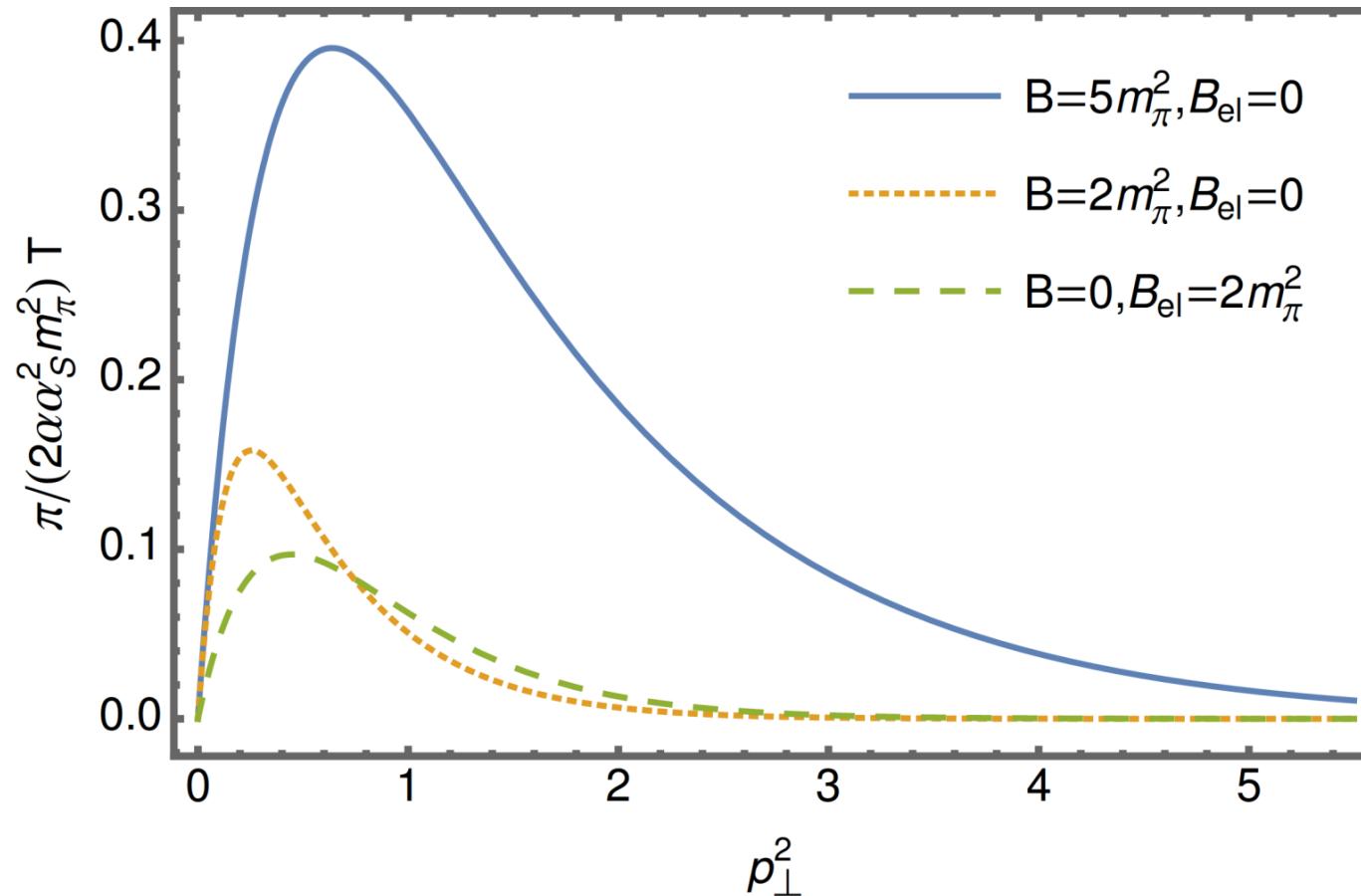
$$T(p, k) = \frac{2\alpha\alpha_s^2}{\pi} q_f^2 (2p_\perp^2 + k_\perp^2 + p_\perp k_\perp) \exp \left\{ -\frac{1}{|q_f B_{el}|} (p_\perp^2 + k_\perp^2 + p_\perp k_\perp) \right\}. \quad (8)$$

- 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.

$$|q_f B_{el}| \xrightarrow{\text{red}} |q_f B_{el} + \hat{n}B|$$

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

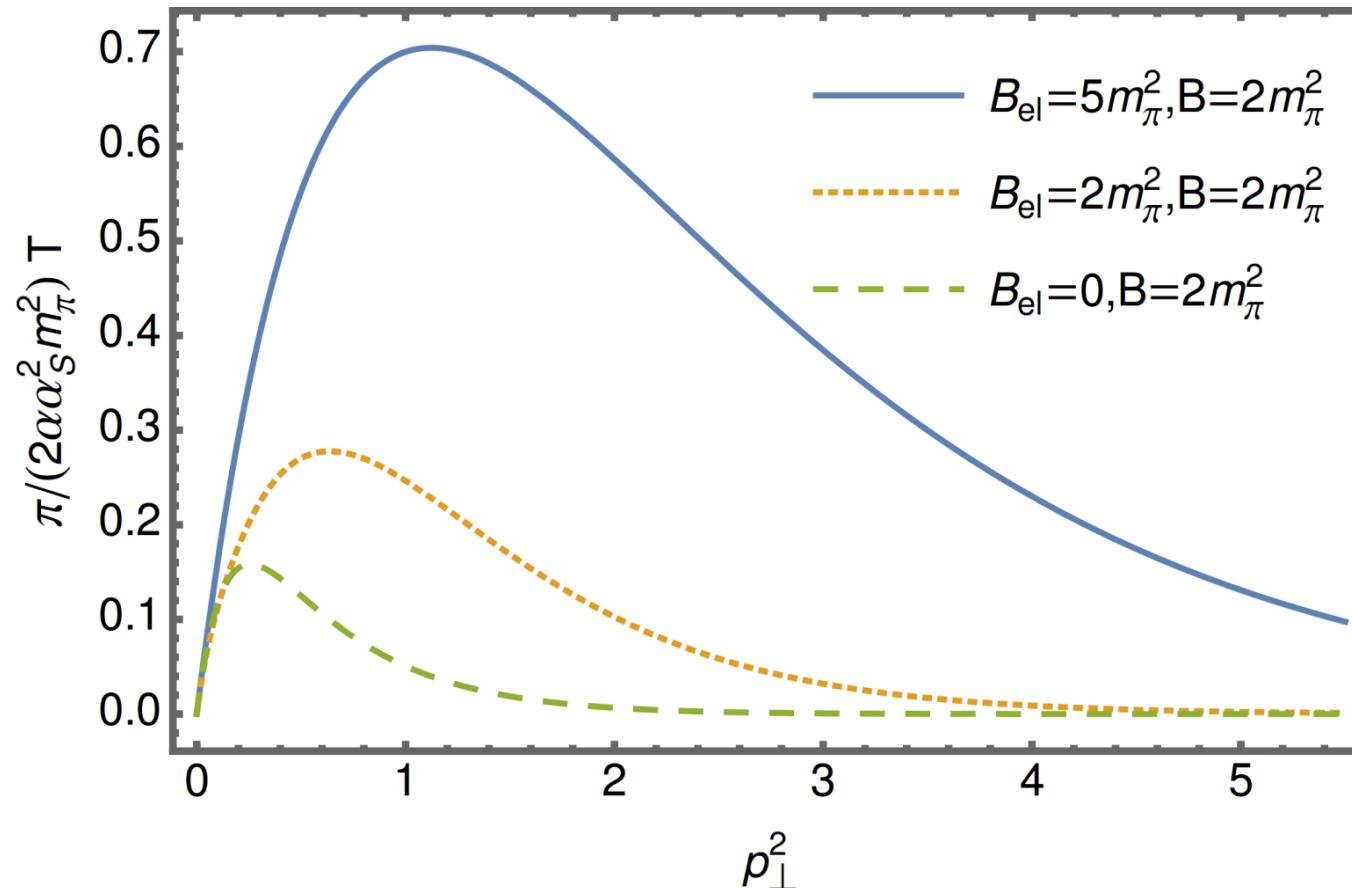
The squared amplitude / massless quarks



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.

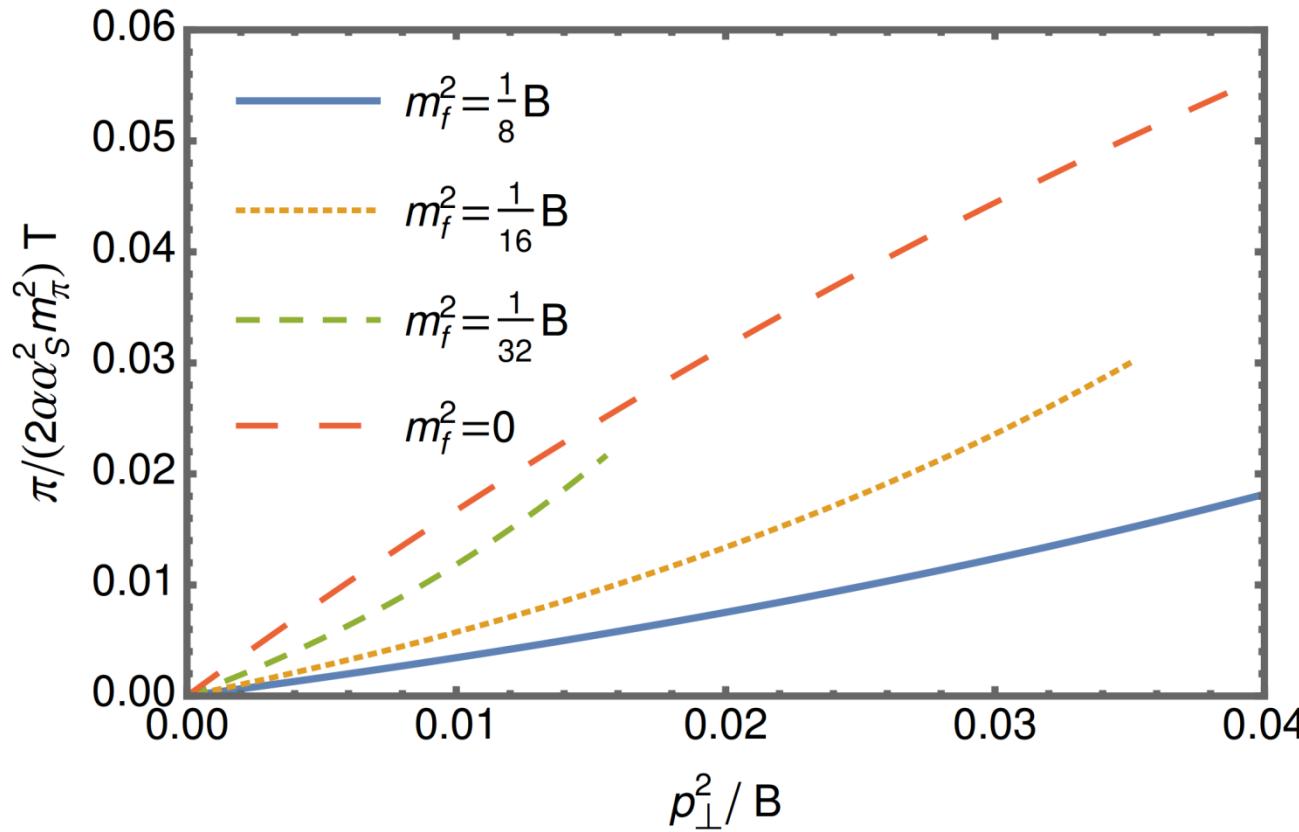
The amplitude squared / massless quarks



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 $B_{\text{el}}=2m_\pi^2$, $B_{\text{el}}=5m_\pi^2$ 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

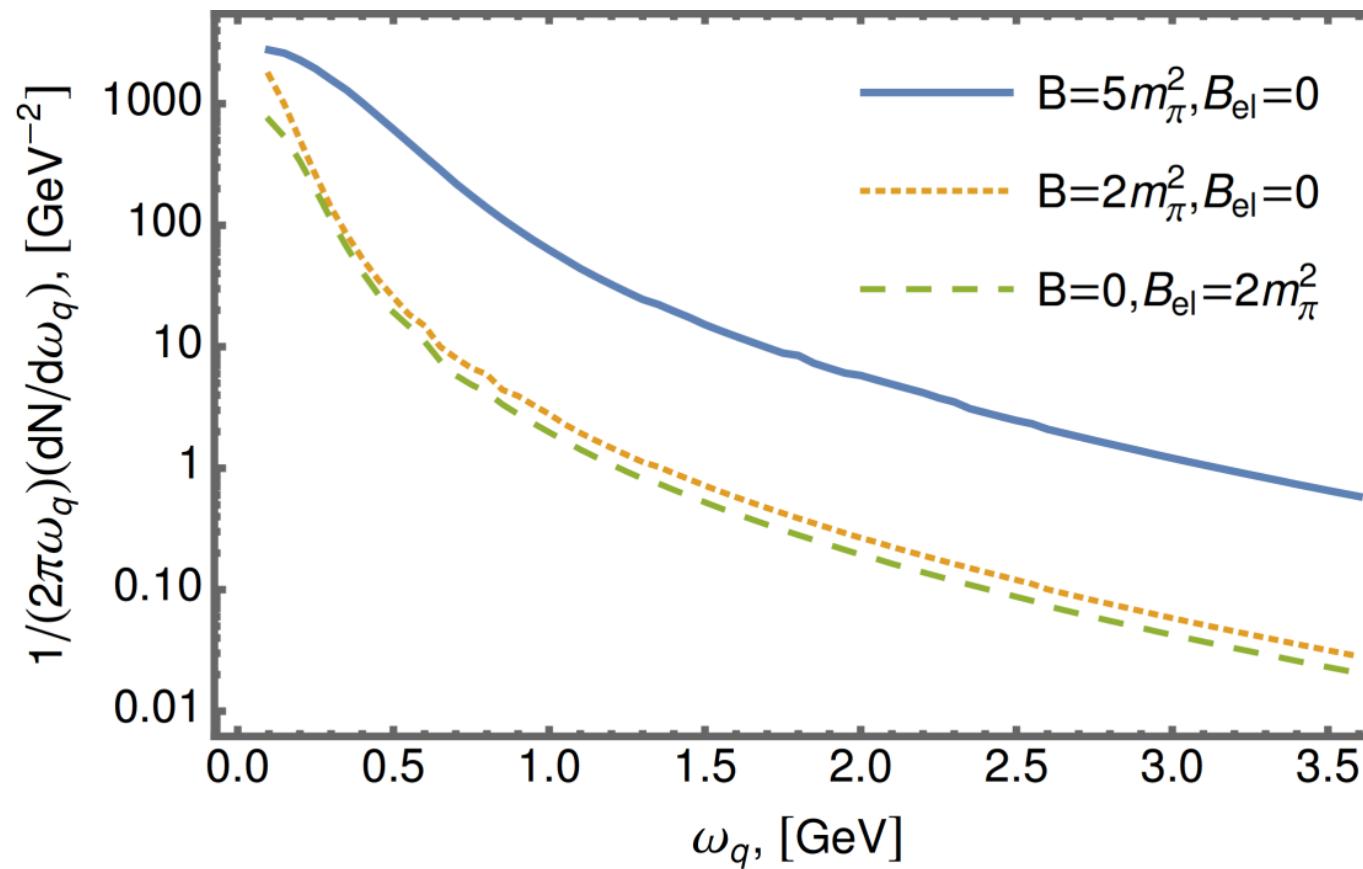
The invariant photon distribution in the presence of chromomagnetic and magnetic fields can be represented in the form

$$\frac{1}{2\pi\omega_q} \frac{dN}{d\omega_q} = v\Delta\tau \frac{\alpha\alpha_s^2\pi}{2N_c(2\pi)^6\omega_q} q_f^2 \text{Tr}_{\hat{n}} \int_0^{\omega_q} d\omega_p (2\omega_p^2 + \omega_q^2 - \omega_p\omega_q) e^{-g_f^B(\omega_p, \omega_q)} \\ \times \left[I_0(g_f^B(\omega_p, \omega_q)) - I_1(g_f^B(\omega_p, \omega_q)) \right] \left(n(\omega_p) n(|\omega_q - \omega_p|) \right), \quad (10)$$

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

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

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 $m_\pi = 0.135 \text{ GeV}$ 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 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!