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CONFORMAL INVARIANCE IN DEFORMED N=4 SUPER YANG-MILLS THEORY

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1. During the last decade much attention has been paid to the $\mathcal{N}=4$ supersymmetric Yang-Mills theory (SYM) and its deformations obtained by the orbifold or orientifold projection, or by adding the marginal deformations to the Lagrangian. Such deformations lead to theories with less supersymmetry but possibly inheriting some attractive features of the original $\mathcal{N}=4$ SYM theory, namely, the conformal invariance, integrability in the planar limit and, especially, its connection with dual string theory via the AdS/CFT correspondence.

In the case of deformed $\mathcal{N}=4$ SYM theory the initial symmetry is broken down to $\mathcal{N}=1$ supersymmetry and $SU(4)_R$ global group down to $U(1)_R$. One of such examples is the so-called β-deformation of the original $\mathcal{N}=4$ SYM theory. Its gravity dual was constructed by Lunin and Maldacena and a significant role in this duality is played by the $U(1)\times U(1)$ global symmetry of the β-deformed theory which was associated with isometries of the deformed $AdS_5\times\tilde{S}_5$ background.

From the point of view of field theory the investigation was made of the so-called β-deformed case. We discuss here the conformal invariance of the full Leigh-Strassler deformation. Using the dimensional regularization (reduction) we found conditions of conformal invariance up to four loops in the planar limit and up to three loops in the nonplanar one. We also found solutions when the conformal conditions were exhausted in the one-loop order.

2. The so-called Leigh-Strassler deformation can be obtained by modification of the superpotential in the original $\mathcal{N}=4$ SYM theory written in terms of $\mathcal{N}=1$ superfields

$$S \;=\; \int d^8z Tr\left(e^{-gV}\bar{\Phi}_i e^{gV}\Phi^i\right)+\left(\frac{1}{2g^2}\int d^6z Tr(W^\alpha W_\alpha)+\int d^6z\ \mathcal{W}+h.c.\right) \qquad (1)$$

in such a way that

$$\begin{aligned}
\mathcal{W}_{N=4\ SYM} &= ig(Tr(\Phi_1\Phi_2\Phi_3)-Tr(\Phi_1\Phi_3\Phi_2)) \Rightarrow \qquad (2)\\
\mathcal{W}_{LS\ SYM} &= i[h_1Tr(\Phi_1\Phi_2\Phi_3)-h_2Tr(\Phi_1\Phi_3\Phi_2)+\frac{h_3}{3}\sum_{i=1}^{3}Tr(\Phi_i^3)],
\end{aligned}$$

where Φ_i with $i=1,2,3$ are the three chiral superfields of the original $\mathcal{N}=4$ SYM theory in the adjoint representation of the gauge group $SU(N)$ and the couplings h_1,h_2,h_3 are in the general complex. The beta-deformed case in the same notation corresponds to

$$h_1=hq,\ h_2=h/q,\ q=e^{i\pi\beta} \text{ and } h_3=0.$$

We are interested in the UV divergences. In our case, as in any $\mathcal{N}=1$ SYM theory formulated in terms of $\mathcal{N}=1$ superfields, one has two types of divergent diagrams, those of the chiral field propagator and of the gauge field one. Moreover, the gauge propagator is not independent: its divergences are related to the chiral propagators. Hence, the conformal invariance being understood as the vanishing of the beta function is valid on the submanifold in the coupling constant space which is defined by condition:

$$\gamma(g,\ \{h_i\}) = 0, \tag{3}$$

where $\{h_i\} = (h_1,\ h_2,\ h_3)$. One can solve this condition choosing the Yukawa couplings in the form of perturbation series over g:

$$h_i = \alpha_{0i}g + \alpha_{1i}g^3 + \alpha_{2i}g^5 + \ldots\ , i = 1...3. \tag{4}$$

If the anomalous dimensions of the chiral fields vanish, so do the gauge and Yukawa beta functions and the theory is conformally invariant.

Conformal invariance also means that the theory is finite, i.e. all UV divergences cancel (or in some gauges the sum of divergences) and the renormalization factors Z (or their products) are equal to 1 or finite. In the context of dimensional regularization this can be achieved by adding to expansion over g (4) a similar expansion over the parameter of dimensional regularization $\varepsilon = 4 - D$, i.e., $\alpha_{ij}(\varepsilon)$ and one has the twofold expansion instead of the onefold one. In a given order of PT equal to n one needs all terms of the double expansion with a total power of $g^2 \cdot \varepsilon$ equal to n.

Our goal now is to calculate several terms of the double expansion and to look for particular solutions when expansion breaks down at the first terms. In the 1-loop order the chiral field anomalous dimension is

$$\begin{aligned} \gamma^{(1)}(\{h_i\}, g) &= \frac{N}{(4\pi)^2}\left(f(\{h_i\}, N) - 2g^2\right), \\ f(\{h_i\}, N) &= (1 - \frac{2}{N^2})(|h_1|^2 + |h_2|^2) + \frac{2}{N^2}(h_1\bar{h}_2 + h_2\bar{h}_1) + (1 - \frac{4}{N^2})|h_3|^2. \end{aligned} \tag{5}$$

Thus, the one-loop conformal condition takes the form

$$f(\{h_i\}, N) - 2g^2 = 0. \tag{6}$$

To fulfil it, one has to choose the Yukawa couplings in a proper way. This condition is also valid up to 3 loops in the planar case and up to two loops in the nonplanar case. In higher orders new contributions appear and eq.(6) is modified.

3. Starting from three loops in the nonplanar case one has a new contribution that does not vanish when the one-loop conformal condition is applied

$$\gamma^{(3)}(\{h_i\}, g) = (f(\{h_i\}, N) - 2g^2)P_{31}(\{h_i\}, g^2, N) + G_{31}(\{h_i\}, N), \tag{7}$$

The explicit form of G_{31} is

$$\begin{aligned} G_{31}(\{h_i\}, N) &= -\frac{1}{128}\frac{6\zeta(3)}{(4\pi)^6}\frac{N^2-4}{N^3} \times \\ &\Big\{|h_1 - h_2|^2\left(N^2|h_1^2 + h_2^2 + h_1h_2|^2 - 9N^2|h_1|^2|h_2|^2 + 5|h_1 - h_2|^4\right) \\ &-18|h_3|^2\left((N^2-5)|h_1^2 + h_2^2|^2 - (N^2-10)\left(\bar{h}_1\bar{h}_2(h_2^2 + h_1^2) + c.c.\right) - 20|h_1|^2|h_2|^2\right) \\ &+\left(\bar{h}_3^3(h_1 - h_2)((N^2+20)(h_1^2 + h_2^2) + 10(N^2-4)h_1h_2) + c.c.\right) - 8(N^2-10)(|h_3|^2)^3\Big\} \end{aligned} \tag{8}$$

This means that in 3 loops the couplings $\{h_i\}$ must satisfy the following modified condition:

$$f(\{h_i\},N) \; = \; (1-\frac{2}{N^2})(|h_1|^2+|h_2|^2)+\frac{2}{N^2}(h_1\bar{h}_2+h_2\bar{h}_1)+(1-\frac{4}{N^2})|h_3|^2$$
$$= g^2\left\{2-\frac{\zeta_3}{128}G^{\Sigma}_{31}\varepsilon^2-\frac{2\zeta_3}{128}G^{\Sigma}_{31}\left(\frac{g^2N}{16\pi^2}\right)\varepsilon+\frac{18\zeta_3}{128}G^{\Sigma}_{31}\left(\frac{g^2N}{16\pi^2}\right)^2\right\}. \quad (9)$$

The situation is simplified in the planar (large N of the $SU(N)$ gauge group) limit. In this case, the first nonvanishing contribution comes in 4 loops

$$\begin{aligned} & c_{41}(\{h_i\},g^2,N) = \qquad (10)\\ & = \frac{5}{2}\zeta(5)\frac{N^4}{(4\pi)^8}\{(|h_1|^2+|h_2|^2+|h_3|^2)^4-(2g^2)^4+(|h_1|^2-|h_2|^2)^4+(|h_3|^2)^4 \\ & +6(|h_3|^2)^2(|h_1|^2+|h_2|^2)^2+24|h_3|^2|h_1|^2|h_2|^2(|h_1|^2+|h_2|^2)+ \\ & +8h_3^3(|h_2|^2\bar{h}_1^3-|h_1|^2\bar{h}_2^3)+8\bar{h}_3^3(|h_2|^2h_1^3-|h_1|^2h_2^3) \\ & -8|h_3|^2(h_2^3\bar{h}_1^3+h_1^3\bar{h}_2^3)-4|h_3|^2(|h_1|^2+|h_2|^2)^3-4(|h_3|^2)^3(|h_1|^2+|h_2|^2)\}. \end{aligned}$$

Again, as in the previous case, one gets a finite and conformal theory up to four loops if the renormalized Yukawa couplings are chosen to satisfy the condition

$$f(\{h_i\},N)=|h_1|^2+|h_2|^2+|h_3|^2=g^2\{2+\frac{5}{18}\zeta_5G^{\Sigma}_{41}\varepsilon^3+\frac{5}{3}\zeta_5G^{\Sigma}_{41}(\frac{g^2N}{16\pi^2})\varepsilon^2$$
$$+\;\; 5\zeta_5G^{\Sigma}_{41}(\frac{g^2N}{16\pi^2})^2\varepsilon+10\zeta_5G^{\Sigma}_{41}(\frac{g^2N}{16\pi^2})^3+...\}.$$

4. Consider now if one can find such values of (h_1,h_2,h_3) that G_{31} in the non-planar case and G_{41} in the planar case vanish meaning that the one-loop conformal condition is valid up to three or four loops, respectively.

In the nonplanar case, similar to the beta-deformed theory, we have not found any solution for vanishing of G_{31} which has a simple form and might be valid in any order of PT. In the planar case, on the contrary, we found two families of simple solutions of the equation $G_{41} = 0$. However, not all of these solutions are genuine. Some of them happen to be unitary equivalent to the β-deformed case.

The only nontrivial solution that exists in the planar limit and leads to conformal theory (up to 4 loops at least) corresponds to the superpotential which can be written in the form

$$\mathcal{W}=ih\int d^6z(qTr\Phi_1\Phi_2\Phi_3-\frac{1}{q}\sum_{i=1}^{3}\frac{Tr(\Phi_i^3)}{3}). \quad (11)$$

where $|h|^2=g^2$ and $|q|=1$, but $q\neq e^{i\frac{\pi n}{3}}$.

One may wonder if the theory defined by the superpotential (11) is conformal in the planar limit when $|q|=1$ in all loops precisely like the beta-deformed one. We have no rigorous proof of this statement, but the analysis of the diagrams suggests the conjecture that the theory defined by the superpotential (11) with $|q|=1$ is exactly conformal in the planar limit. We show in Fig.1 the solutions in the coupling constant space. The dashed areas correspond to conformal theories obtained via perturbative expansion. The exact solutions are denoted by dots.

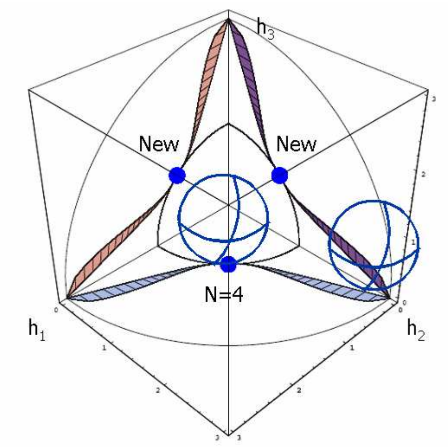

Рис. 1: The coupling constant space of a general Leight-Strassler theory. The $N = 4$ solution as well as the new one are shown by blue dots. At each point one has an arbitrary phase shown by a sphere.

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ORTHOPOSITRONIUM LIFETIME: ANALYTIC RESULTS IN $\mathcal{O}(\alpha)$ AND $\mathcal{O}(\alpha^3 \ln \alpha)$

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Quantum electrodynamics (QED), the gauged quantum field theory of the electromagnetic interaction, celebrated ground-breaking successes in the twentieth century. In fact, its multiloop predictions for the anomalous magnetic moments of the electron and the muon were found to agree with the highest-precision measurements within a few parts of 10^{-12} and 10^{-10}, respectively.

Another ultrapure laboratory for high-precision tests of QED is provided by positronium (Ps), the lightest known atom, being the electromagnetic bound state of the electron e^- and the positron e^+, which was discovered in the year 1951 [1]. In fact, thanks to the smallness of the electron mass m relative to typical hadronic mass scales, its theoretical description is not plagued by strong-interaction uncertainties and its properties such as decay widths and energy levels can be calculated perturbatively in nonrelativistic QED (NRQED) [2], as expansions in Sommerfeld's fine-structure constant α, with very high precision.

Ps comes in two ground states, $^1S_0$ parapositronium (p-Ps) and $^3S_1$ orthopositronium (o-Ps), which decay into two and three photons, respectively. Particulary, the lifetime of o-Ps has been the subject of a vast number of theoretical and experimental investigations. Numerically the o-Ps lifetime Γ can be evaluated from

$$\Gamma \;=\; \frac{2}{9}(\pi^2-9)\frac{m\alpha^6}{\pi}\left[1 + A\frac{\alpha}{\pi} + \frac{\alpha^2}{3}\ln\alpha + B\left(\frac{\alpha}{\pi}\right)^2 - \frac{3\alpha^3}{2\pi}\ln^2\alpha + C\frac{\alpha^3}{\pi}\ln\alpha\right], \qquad (1)$$

where the constants A, B and C have been known only numerically.

After a laborious calculation, authors of [3] obtained the constant A in the form

$$\begin{aligned}
\frac{2}{9}(\pi^2-9)A \quad &= \frac{56}{27} - \frac{901}{216}\zeta_2 - \frac{11303}{192}\zeta_4 + \frac{19}{6}l_2 - \frac{2701}{108}\zeta_2 l_2 + \frac{253}{24}\zeta_2 l_2^2 \\
&+ \frac{251}{144}l_2^4 + \frac{913}{64}\zeta_2 l_3^2 + \frac{83}{256}l_3^4 - \frac{21}{4}\zeta_2 l_2 l_r - \frac{49}{16}\zeta_2 l_r^2 + \frac{7}{16}l_2 l_r^3 + \frac{35}{384}l_r^4 \\
&+ \frac{581}{16}\zeta_2 \,\mathrm{Li}_2\left(\frac{1}{3}\right) - \frac{21}{2}l_2\,\mathrm{Li}_3(-r) - \frac{7}{2}l_r\,\mathrm{Li}_3(-r) + \frac{63}{4}l_2\,\mathrm{Li}_3(r) \\
&+ \frac{63}{8}l_r\,\mathrm{Li}_3(r) - \frac{249}{32}\,\mathrm{Li}_4\left(-\frac{1}{3}\right) + \frac{249}{16}\,\mathrm{Li}_4\left(\frac{1}{3}\right) + \frac{251}{6}\,\mathrm{Li}_4\left(\frac{1}{2}\right) \\
&+ 7\,\mathrm{Li}_4(-r) - 7\mathrm{S}_{2,2}(-r) - \frac{63}{4}\,\mathrm{Li}_4(r) + \frac{63}{4}\,\mathrm{S}_{2,2}(r) + \frac{11449}{432}\zeta_3 \\
&- \frac{91}{6}\zeta_3 l_2 - \frac{35}{8}\zeta_3 l_r + \frac{1}{\sqrt{2}}\left[\frac{49}{2}\zeta_2 l_r - \frac{7}{72}l_r^3 - \frac{35}{6}l_r\,\mathrm{Li}_2(r)\right. \\
&\left. + \frac{35}{6}\,\mathrm{Li}_3(r) - \frac{175}{3}\,\mathrm{S}_{1,2}(r) + \frac{14}{3}\,\mathrm{S}_{1,2}(r^2) + \frac{119}{3}\zeta_3\right], \qquad (2)
\end{aligned}$$

where $r=(\sqrt{2}-1)/(\sqrt{2}+1)$, $l_x=\ln x$,

$$\mathrm{S}_{n,p}(x)=\frac{(-1)^{n+p-1}}{(n-1)!\,p!}\int_0^1\frac{\mathrm{d}t}{t}\ln^{n-1}t\ln^p(1-tx) \tag{3}$$

is the generalized polylogarithm, $\mathrm{Li}_n(x)=\mathrm{S}_{n-1,1}(x)$ is the polylogarithm of order n, and $\zeta_n=\zeta(n)=\mathrm{Li}_n(1)$, with $\zeta(x)$ being Riemann's zeta function.

An analytic expression for C is then simply obtained from (2) through the relationship [4]

$$C=\frac{A}{3}-\frac{229}{30}+8\ln 2, \tag{4}$$

but the constant B in Eq. (1) still remains analytically unknown.

From Eqs. (2) and (4), A and C can be numerically evaluated with arbitrary precision,

$$\begin{aligned} A &= -10.28661\,48086\,28262\,24015\,01692\,10991\,\ldots\,, \\ C &= -5.51702\,74917\,29858\,27137\,88660\,98665\,\ldots\,. \end{aligned} \tag{5}$$

These numbers agree with the best existing numerical evaluations [5].

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NLO QED Radiative Corrections to Exclusive Observables

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The continuously increasing experimental precision requires more and more accurate theoretical predictions. Among other effects higher order QED radiative corrections become of crucial importance. For a large class of processes the renormalization group (RG) can help to obtain the numerically most important part of those corrections. The RG approach is a very powerful method based on the principle of scale invariance. In high energy physics it helps to study the dependence of observable results on the energy scale. Evolution equations for large logarithms of the scale were first derived in the early 1970s for scalar QED and immediately extended to the QCD case known now as the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) evolution equations. The first application of the method for spinor QED in the leading logarithmic approximation (LLA) was made by Kuraev and Fadin only in 1985 in Ref. [1]. This gap in time is due to the following reasons: 1) there were other approaches in QED such as the method of *equivalent electrons* which were sufficient for the time being; 2) the smallness of the QED coupling constant, $\alpha_{QED} \ll \alpha_{QCD}$, makes QED corrections to be not so important numerically as the QCD ones; 3) typically there is a difference in the degree of inclusiveness of QED and QCD observables. For many applications the latter makes it impossible to apply the QCD RG method to QED without crucial modifications.

The QCD factorization theorem can be directly adopted for an inclusive QED NLO case *e.g.* for Bhabha scattering without imposing cuts on events with photon radiation in the following form:

$$\begin{aligned}
\mathrm{d}\sigma &= \int_{\bar{z}_1}^{1} \mathrm{d}z_1 \int_{\bar{z}_2}^{1} \mathrm{d}z_2 \sum_{a,b,c,d=e^{\pm},\gamma} \mathcal{D}^{\mathrm{str}}_{ae}(z_1)\mathcal{D}^{\mathrm{str}}_{be}(z_2) \left(\mathrm{d}\sigma^{(0)}(z_1,z_2)\right. \\
&+ \left.\mathrm{d}\bar{\sigma}^{(1)}(z_1,z_2) + \mathcal{O}\left(\alpha^2 L^0\right)\right) \int_{\bar{y}_1}^{1} \frac{\mathrm{d}y_1}{Y_1} \int_{\bar{y}_2}^{1} \frac{\mathrm{d}y_2}{Y_2} \mathcal{D}^{\mathrm{frg}}_{ec}(\frac{y_1}{Y_1}) \mathcal{D}^{\mathrm{frg}}_{ed}(\frac{y_2}{Y_2}),
\end{aligned} \tag{1}$$

where $\mathcal{D}^{\mathrm{str(frg)}}_{ba}(x)$ is a structure (fragmentation) function giving the probability density to find parton b in parton a with the relative energy fraction x. These functions contain the full dependence of the final result on the so-called large logarithm $L = \frac{M^2}{m_e^2}$, where M is the factorization scale and m_e is the electron (charged fermion) mass, $m_e \ll M$. In the NLO approximation the photonic part of the QED electron structure function takes the form

$$\begin{aligned}
\mathcal{D}^{\mathrm{str,frg}}_{ee}(z) &= \delta(1-z) + \frac{\alpha}{2\pi} d^{(1)}(z,\mu_0,m_e) + \frac{\alpha}{2\pi} L P^{(0)}(z) \\
&+ \left(\frac{\alpha}{2\pi}\right)^2 \left(\frac{1}{2} L^2 P^{(0)} \otimes P^{(0)}(z) + L P^{(0)} \otimes d^{(1)}(z,\mu_0,m_e)\right. \\
&\qquad \left. + L P^{(1,\gamma)\mathrm{str,frg}}_{ee}(z)\right) + \mathcal{O}\left(\alpha^2 L^0, \alpha^3\right),
\end{aligned} \tag{2}$$

where μ_0 is the renormalization scale, $P^{(0,1)}$ are QED splitting functions and $d^{(1)}$ is the relevant NLO initial condition.

The kernel cross sections $\mathrm{d}\sigma^{(0)}(z_1, z_2)$ and $\mathrm{d}\bar{\sigma}^{(1)}(z_1, z_2)$ describe the differential distributions of the process $a + b \to c + d$ in the Born and one-loop approximation, respectively. Partons a, b, c, d are considered as massless particles, the collinear singularities are removed by a subtraction, e.g., in the $\overline{\mathrm{MS}}$ scheme. See Refs. [2, 3] for further details on the notation.

Application of this master formula within the LLA to numerous processes became a standard approach mainly due to requests of very precise experiments at LEP [1, 4, 5]. This scheme can also be used also to get higher order radiative QED corrections to decay processes, e.g., to the muon decay spectrum [6].

The first application of the approach within the next-to-leading order (NLO) was done in Ref. [7] to the process of electron-positron annihilation with an inclusive treatment of photonic radiation. Only 15 year later the NLO case was considered in other cases [8, 9] again for inclusive observables.

In the master equation (1) integration over the real photon phase space is performed without any restriction. That cannot be realized in most of the modern high energy physics experiments. In [2, 3], a new scheme to evaluate NLO QED radiative corrections was proposed. It is based on the slicing of the photon phase space both in energy and angle. As usual we separate the virtual and soft radiation by a small parameter $\Delta \ll 1$, saying that photons with energy above (below) ΔE_{beam} are called *hard* (*soft*). In addition, radiation of photons with momenta collinear to one of the final state charged particles is separated by introduction of one more auxiliary parameter θ_0. So the cross section can be presented as a sum of several contributions with particular kinematics:

$$\mathrm{d}\sigma = \mathrm{d}\sigma^{(0)} + \mathrm{d}\sigma^{(1)}_{\mathrm{S+V}} + \mathrm{d}\sigma^{(1)}_{\mathrm{H}} + \mathrm{d}\sigma^{(2)NLO}_{\mathrm{S+V}} + \mathrm{d}\sigma^{(2)NLO}_{\mathrm{H}} + \mathrm{d}\sigma^{(3)LO} + \ldots \tag{3}$$

The hard photon contribution is also split:

$$\mathrm{d}\sigma^{(2)NLO}_{\mathrm{H}} = \mathrm{d}\sigma^{(2)}_{\mathrm{HH(coll)}} + \mathrm{d}\sigma^{(2)}_{\mathrm{HH(s-coll)}} + \mathrm{d}\sigma^{(2)}_{\mathrm{(S+V)H(n-coll)}} + \mathrm{d}\sigma^{(2)}_{\mathrm{(S+V)H(coll)}}, \tag{4}$$

where “coll” means collinear photon(s) with $\theta_\gamma < \theta_0$, “n-coll” means non-collinear photon with $\theta_\gamma > \theta_0$, and “HH(s-coll)” means semi-collinear kinematics of double photon emission, *i.e.* one collinear photon and one non-collinear photon.

Emission of a hard photon going at a small angle with respect to a charged particle a can be factorized out of the non-radiative differential cross section:

$$\begin{aligned} &\mathrm{d}\sigma[a(p_1) + b \to c + d + \gamma((1-z)p_1)] = \mathrm{d}\sigma[a(zp_1) + b \to c + d] \otimes R^{\mathrm{ISR}}_{\mathrm{H}}(z), \\ &R^{\mathrm{ISR}}_{\mathrm{H}}(z) = \frac{\alpha}{2\pi}\left[\frac{1+z^2}{1-z}\left(\ln\frac{E^2}{m_e^2} - 1 + l_0\right) + 1 - z + \mathcal{O}\left(\frac{m_e^2}{E^2}\right) + \mathcal{O}\left(\theta_0^2\right)\right], \qquad l_0 = \ln\frac{\theta_0^2}{4}, \end{aligned}$$

where z is the energy fraction of the charged particle after radiation of the photon(s), and E is the initial energy of the charged particle. The ISR and FSR radiation factors in $\mathcal{O}(\alpha)$ are well known and used for many applications. The second order contributions in the NLO approximation were derived in Ref. [3] with the help of the general RG approach. We got universal factors which describe emission of photons inside narrow cones along incoming or outgoing charged particles. In this way, for the initial state radiation of two photons in the same narrow cone we got

$$
\begin{aligned}
R^{\rm ISR}_{\rm HH}(z) &= \left(\frac{\alpha}{2\pi}\right)^2 L\biggl\{(L+2l_0)\biggl(\frac{1+z^2}{1-z}(2\ln(1-z)-2\ln\Delta-\ln z)+\frac{1+z}{2}\ln z \\
&- 1+z\biggr)+\frac{1+z^2}{1-z}\biggl(\ln^2 z+2\ln z-4\ln(1-z)+4\ln\Delta\biggr)+(1-z)\biggl(2\ln(1-z) \\
&- 2\ln\Delta-\ln z+3\biggr)+\frac{1+z}{2}\ln^2 z\biggr\}, \qquad L=\ln\frac{E^2}{m_e^2}\,. \qquad (5)
\end{aligned}
$$

The corresponding radiation factor for the final state radiation can be obtained from the ISR one by means of the modified Gribov–Lipatov relation:

$$
R^{\rm FSR}_{\rm HH}(z) \ = \ -zR^{\rm ISR}_{\rm HH}\left(\frac{1}{z}\right)\bigg|_{\ln\Delta\to\ln\Delta-\ln z;\ l_0\to l_0+2\ln z}. \qquad (6)
$$

For the single collinear hard photon radiation accompanied by one-loop virtual or soft photonic correction we got

$$
\begin{aligned}
&R^{\rm ISR}_{\rm H(S+V)}(z)\otimes {\rm d}\sigma(z) = \delta^{(1)}_{\rm (S+V)}R^{\rm ISR}_{\rm H}(z)\otimes {\rm d}\sigma^{(0)}(z) + \left(\frac{\alpha}{2\pi}\right)^2 L\biggl[2\frac{1+z^2}{1-z} \\
&\times\biggl({\rm L}i_2\,(1-z)-\ln(1-z)\ln z\biggr)-(1+z)\ln^2 z+(1-z)\ln z+z\biggr]\otimes {\rm d}\sigma^{(0)}(z). \quad (7)
\end{aligned}
$$

where $\delta^{(1)}_{\rm (S+V)}$ is the relative $\mathcal{O}\,(\alpha)$ contribution due to single soft photon emission and one-loop virtual correction. Note that the factorization of the hard and soft plus virtual corrections in Eq. (7) is not complete. The corresponding final state radiation factor can be received in the same way as in Eq. (6).

Besides the photonic corrections, one has to take into account the effects related to creation of virtual and/or real electron-positron pairs. Pair emission also gives rise to contributions enhanced by the large logarithms in higher orders of perturbative QED. Singlet and non-singlet NLO pair contributions in $\mathcal{O}\,(\alpha^2 L)$ to inclusive observables can be described within the QCD-like factorization approach (1). However exclusive treatment here is of ultimate importance because practically all experiments impose cuts or some other special conditions on the final state charged particles. So we need a Monte Carlo generator for hard pairs, while the corresponding soft plus virtual part can be obtained directly from the master equation (1), so that

$$
{\rm d}\sigma^{(2)}_{\rm pair} = {\rm d}\sigma^{(2)MC}_{\rm H\ pair} + {\rm d}\sigma^{(0)}\times\delta^{(2)}_{\rm S+V\ pair}. \qquad (8)
$$

In this way we constructed a systematic scheme to evaluate the leading and next-to-leading QED logarithmic corrections to realistic high energy physics observables. The method can be applied to such processes as Bhabha scattering [10], muon decay spectrum, deep inelastic scattering, different electron-positron annihilation channels *etc.*.

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INFRARED BEHAVIOR OF THE ADLER D–FUNCTION AND THE INCLUSIVE τ LEPTON DECAY

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The Adler D–function, being defined as the logarithmic derivative of the hadronic vacuum polarization function, plays a crucial role in various issues of elementary particle physics. Specifically, it is essential for the theoretical analysis of such strong interaction processes as electron–positron annihilation into hadrons and inclusive τ lepton hadronic decay. Furthermore, the Adler D–function plays a key role for confronting precise experimental measurements of some electroweak observables (e.g., muon anomalous magnetic moment and shift of the electromagnetic fine structure constant) with their theoretical predictions, giving rise to decisive tests of the Standard Model and furnishing stringent constraints on possible new physics beyond it.

Recently, new integral representations for the Adler D–function and R–ratio of the electron-positron annihilation into hadrons have been derived [1, 2]:

$$D(Q^2) = \frac{Q^2}{Q^2+4m_\pi^2}\left[1+\int_{4m_\pi^2}^{\infty}\rho(\sigma)\frac{\sigma-4m_\pi^2}{\sigma+Q^2}\frac{d\sigma}{\sigma}\right], \tag{1}$$

$$R(s) = \theta(s-4m_\pi^2)\left[1+\int_{s}^{\infty}\rho(\sigma)\,\frac{d\sigma}{\sigma}\right], \tag{2}$$

where the spectral function $\rho(\sigma)$ can be determined either as the discontinuity of an explicit theoretical expression for $D(Q^2)$ across the physical cut, or, equivalently, as the logarithmic derivative of the experimental data on R–ratio:

$$\rho(\sigma) = \frac{1}{2\pi i}\lim_{\varepsilon\to 0_+}\Big[D_{\text{theor}}(-\sigma-i\varepsilon) - D_{\text{theor}}(-\sigma+i\varepsilon)\Big] = -\frac{d\,R_{\text{exp}}(\sigma)}{d\,\ln\sigma}. \tag{3}$$

The representations (1) and (2) capture both the effects due to the analytic continuation of spacelike theoretical results into timelike domain and the effects due to the nonvanishing mass of the lightest hadron state m_π. The latter plays a crucial role in this analysis, forcing the Adler D–function to vanish in the infrared limit. Within the approach in hand the Adler D–function was calculated [1,2] by employing its perturbative approximation as the only additional input, namely

$$\rho(\sigma) = \frac{1}{2\pi i}\lim_{\varepsilon\to 0_+}\Big[D_{\text{pert}}(-\sigma-i\varepsilon) - D_{\text{pert}}(-\sigma+i\varepsilon)\Big]. \tag{4}$$

The obtained result is in a reasonable agreement with the experimental prediction for the Adler D–function in the entire energy range, see Fig. 1.

The inclusive τ lepton decay is studied [3] in the framework of the developed approach. In particular, the nonstrange part of the inclusive semileptonic branching ratio associated with the vector quark currents can be represented in the following form:

$$R_{\tau,\text{V}} = \frac{N_c}{2}\,|V_{\text{ud}}|^2 S_{\text{EW}}\left(\Delta_{\text{QCD}} + \delta'_{\text{EW}}\right), \tag{5}$$

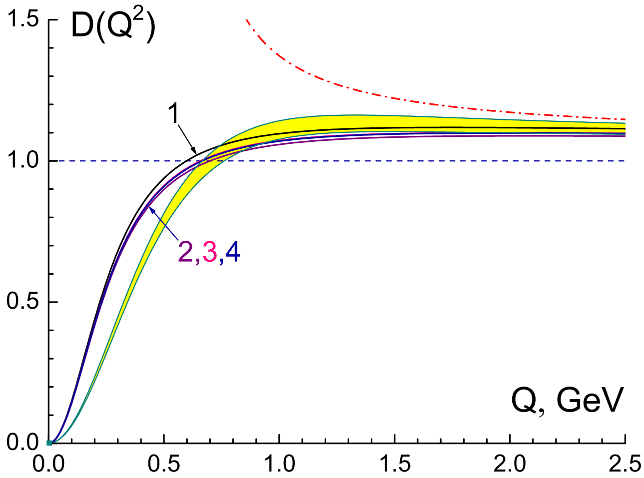

Рис. 1: The comparison of the Adler D–function (1) corresponding to perturbative spectral density (4) (solid curves, labels indicate the loop level) with its experimental prediction (shaded band). The one–loop perturbative approximation of the Adler D–function is denoted by dot–dashed curve. The values of parameters are: $\Lambda = 608\,\mathrm{MeV}$, $n_f = 3$ active flavors.

where $N_c = 3$ is the number of colors, $|V_{\mathrm{ud}}| = 0.97418 \pm 0.00027$ denotes the Cabibbo–Kobayashi–Maskawa matrix element, $S_{\mathrm{EW}} = 1.0194 \pm 0.0050$ and $\delta'_{\mathrm{EW}} = 0.0010$ are the electroweak corrections, and Δ_{QCD} can be expressed in terms of a weighted integral of the aforementioned R–ratio:

$$\Delta_{\mathrm{QCD}} = 2 \int_0^{M_\tau^2} \left(1 - \frac{s}{M_\tau^2}\right)^2 \left(1 + 2\,\frac{s}{M_\tau^2}\right) R(s) \frac{ds}{M_\tau^2}, \tag{6}$$

with $M_\tau \simeq 1.777\,\mathrm{GeV}$ being the τ lepton mass. The experimental measurement of the ratio (5) yields $R_{\tau,\mathrm{V}} = 1.764 \pm 0.016$.

In the framework of perturbative approach one usually reduces Eq. (6) to the contour integral in the complex s–plane along the circle of the radius of the squared mass of the τ lepton. At the one-loop level this eventually leads to

$$\Delta_{\mathrm{QCD}} = 1 + d_1\, \alpha_{\mathrm{s}}^{(1)}(M_\tau^2), \qquad d_1 = \frac{1}{\pi}. \tag{7}$$

At the same time, for the evaluation of Δ_{QCD} in the framework of the approach in hand, the integration in Eq. (6) can be performed in a straightforward way. Ultimately this leads to the following result at the one-loop level [3]:

$$\Delta_{\mathrm{QCD}} = 1 - \delta_\Gamma + d_1 \alpha_{\mathrm{TL}}^{(1)}(M_\tau^2) - d_1 \delta_\Gamma \alpha_{\mathrm{TL}}^{(1)}(m_\Gamma^2) + d_1\, \frac{4\pi}{\beta_0} \int_\chi^1 f(\xi)\, \rho^{(1)}(\xi M_\tau^2)\, d\xi, \tag{8}$$

where $f(\xi) = \xi^3 - 2\xi^2 + 2$, $\chi = m_\Gamma^2/M_\tau^2$, $\delta_\Gamma = \chi\, f(\chi)$,

$$\alpha_{\mathrm{TL}}^{(1)}(s) = \frac{4\pi}{\beta_0}\, \theta(s - m_\Gamma^2) \int_s^\infty \rho^{(1)}(\sigma)\, \frac{d\sigma}{\sigma} \tag{9}$$

is the one-loop timelike effective coupling, and $\rho^{(1)}(\sigma)$ denotes the relevant one-loop spectral function. Here m_Γ stands for the total mass of the lightest allowed hadronic decay mode of the τ lepton, e.g., for the vector channel $m_\Gamma = m_{\pi^0} + m_{\pi^-} \simeq 274.5\,\mathrm{MeV}$. In this case $\delta_\Gamma \simeq 0.048$ considerably exceeds the electroweak correction δ'_{EW}. It turns out that the effects due to the nonvanishing mass of the lightest hadron state play a substantial role in processing the experimental data on the inclusive τ lepton decay, see Ref. [3] for the details.

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AXIAL MASS IN REACTIONS OF QUASIELASTIC ANTINEUTRINO–NUCLEON SCATTERING WITH STRANGE HYPERON PRODUCTION

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In this study we investigate the reactions of quasielastic production of Λ, Σ^-, and Σ^0 hyperons in antineutrino interactions with nucleons, allowed by the selection rules $\Delta Y = \pm 1$, $\Delta I = 1/2$ and $\Delta Y = \Delta Q$ (where Y, I, and Q are, respectively, the hypercharge, isospin, and electric charge). In comparison with the usual $\Delta Y = 0$ quasielastic (anti)neutrino–nucleon scattering, the cross sections for the processes with $\Delta Y = 1$ are suppressed by the factor $\sin^2\theta_C \approx 0.05$ (where θ_C is the Cabibbo angle) and thus they are badly studied experimentally. That is why currently it is impossible to test the validity of the models for the transition form factors entering the weak hadronic current with $\Delta Y = 1$ by using the $\overline{\nu}N$ cross section data directly. The most uncertainty comes from are the axial-vector and pseudoscalar form factors. The pseudoscalar contribution into the $\nu_{e,\mu}n$ and $\overline{\nu}_{e,\mu}p$ cross sections is very small and cannot be measured in the present-day experiments. However it is important for predicting the quasielastic $\nu_\tau n$ and $\overline{\nu}_\tau p$ cross sections for which the experimental data are merely absent but the knowledge of which is necessary for studying the $\nu_\mu \leftrightarrow \nu_\tau$ oscillations in such experiments as ICARUS/ICANOE and OPERA, which will be able to identify the events with τ lepton in the final state.

We derive the most general formulae connecting the structure functions for the quasielastic νN and $\overline{\nu}N$ scattering off nucleons with the (complex) hadronic-current form factors, which take into account the final lepton mass and nonstandard (G parity violating) second-class currents [1, 2]. The transition form factors are related to the $\Delta Y = 0$ once

Таблица 1: Isovector and axial form factors for the $\Delta Y = 1$ reactions.

Reactions	F_V	F_M	F_A
$\overline{\nu}_\ell p \to \ell^+\Lambda$	$-\sqrt{\frac{3}{2}}F_1^p$	$-\sqrt{\frac{3}{2}}F_2^p$	$-\sqrt{\frac{3}{2}}\frac{1+2\xi}{3}G_A$
$\overline{\nu}_\ell n \to \ell^+\Sigma^-$	$-\left(F_1^p + 2F_1^n\right)$	$-\left(F_2^p + 2F_2^n\right)$	$(1-2\xi)\,G_A$
$\overline{\nu}_\ell p \to \ell^+\Sigma^0$	$-\frac{1}{\sqrt{2}}\left(F_1^p + 2F_1^n\right)$	$-\frac{1}{\sqrt{2}}\left(F_2^p + 2F_2^n\right)$	$\frac{1}{\sqrt{2}}(1-2\xi)\,G_A$

trough the approximate SU(3) symmetry [3]. These relations are collected in Table 1 (the notation is standard). For the pseudoscalar form factor we apply the hypothesis by Nambu [4] according to which

$$F_P(Q^2) = M^2\left(m_K^2 + Q^2\right)^{-1} F_A(Q^2), \tag{1}$$

where m_K is the kaon mass.

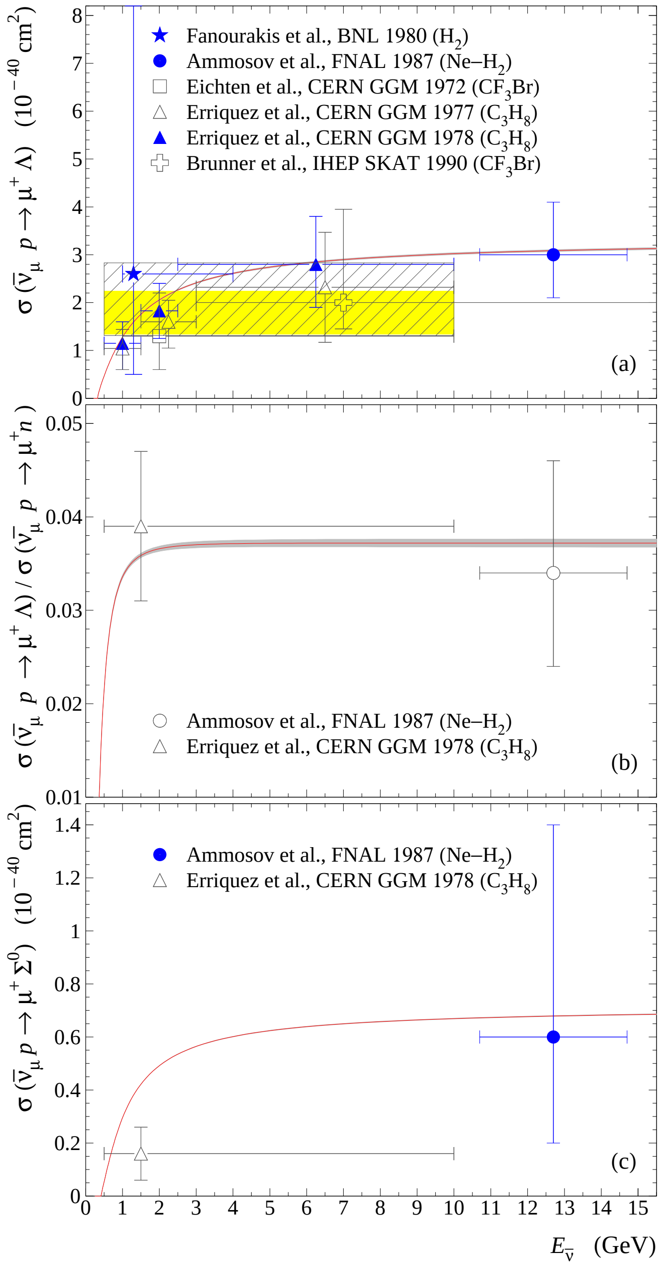

Рис. 1: Total cross sections for the reactions $\overline{\nu}_\mu p \to \mu^+\Lambda$ (a) and $\overline{\nu}_\mu p \to \mu^+\Sigma^0$ (c) and the cross section ratio for the reactions $\overline{\nu}_\mu p \to \mu^+\Lambda$ and $\overline{\nu}_\mu p \to \mu^+ n$ (b), measured in the experiments BNL 1980 [7], FNAL 1987 [8], CERN GGM 1972 [9], CERN GGM 1977 [10], CERN GGM 1978 [11] and IHEP SKAT 1990 [12]. The curves show the calculated cross sections obtained with the BBBA(07) model for the electromagnetic form factors and with $M_A = 0.999$ GeV. The widths of the narrow bands correspond to the statistical error $\Delta M_A = \pm 0.011$ GeV.

A statistical analysis of all available accelerator data on the total and differential cross sections as well as the Q^2 distributions of events for quasielastic scattering of ν_μ and $\overline{\nu}_\mu$ with and without change of hypercharge ($\Delta Y = 0, 1$) on hydrogen, deuterium, carbon, argon, iron, propane, freon, propane–freon mixture, and neon–hydrogen mixture has been performed [1, 2, 5]. In this analysis, we apply the same selection criteria for the experimental data as in [5]. For the electromagnetic form factors G_E^p, G_M^p, G_E^n, and G_M^n, the most accurate phenomenological models were applied, which describe well the available experimental results extracted from the ep scattering data. As an illustration, we show the comparison of our calculations with the available data [7–12] on the cross sections of $\overline{\nu}_\mu p \to \mu^+\Lambda$ and $\overline{\nu}_\mu p \to \mu^+\Sigma^0$ reactions and the cross section ratio for the reactions $\overline{\nu}_\mu p \to \mu^+\Lambda$ and $\overline{\nu}_\mu p \to \mu^+ n$. The curves in the Figure show the calculated cross sections obtained with the BBBA(07) model for the electromagnetic form factors [13]. This model is an accurate Kelly type parametrization of the current experimental data on the form factors G_E^p, G_M^p, G_E^n, G_M^n, and ratio G_E^p/G_M^p, which uses the Nachtmann scaling variable to relate elastic and inelastic form factors, and imposes quark-hadron duality asymptotic constraints at high momentum transfers where the quark structure dominates. Quark-hadron duality implies that the squared ratio of neutron and proton magnetic form factors should be the same as the ratio of the corresponding inelastic structure functions F_2^n and F_2^p in the limit $Q^2 \to \infty$:

$$\left(\frac{G_M^n}{G_M^p}\right)^2 = \frac{F_2^n}{F_2^p} = \frac{1+4(d/u)}{4+(d/u)}, \quad Q^2 \to \infty.$$

Here d and u are the partonic density functions. The data are fitted under the two assumptions: $d/u = 0$ and $d/u = 0.2$. One more duality-motivated constraint is the equality

$$\left(\frac{G_E^n}{G_M^n}\right)^2 = \left(\frac{G_E^p}{G_M^p}\right)^2$$

applied for the highest Q^2 data points for the neutron electric form factor included into the BBBA(07) fit.

Within the Standard Model assumptions, it has been obtained the world-averaged value of the axial mass of the nucleon $M_A = 0.999 \pm 0.011$ GeV/c^2 which is concordant with the results obtained by fitting all the data on exclusive and inclusive νN and $\overline{\nu} N$ reactions. [1, 6].

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NUCLEON IS MORE COMPLICATED THAN IT WAS THOUGHT

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The quark models of the sixties gave a comparatively simple spin structure of the nucleons: three quarks of three different colors, the spins of two u in one direction,and spin of d in the opposite one. "Spin Crisis" of the eighties shows that the matter is much more complicated and that more than half of the nucleon complete spin is carried by either gluons or/and by quark orbital angular momentum. Last measurements of the gluon spin contribution by COMPASS collaboration on the longitudinally polarized target, however, do not support a large gluon spin contribution.

Meanwhile, measurements of spin asymmetry in the semi-inclusive deeply inelastic scattering on the transversely polarized target by HERMES and COMPASS revealed other interesting details of the nucleon spin structure. It occurs that quarks are polarized even in the nonpolarized nucleon; moreover, the direction of their polarization is connected with the direction of their internal momentum (so-called "Sivers asymmetry"and that the transversely polarized quark scattered from the nucleon differently fragments into a hadron depending on that where this hadron flies, to the right or to the left (so-called "Collins asymmetry").

The asymmetry of Sivers and Collins are the most noticeable among the single spin asymmetries in the semi-inclusive deeply inelastic scattering on the transversely polarized target. In the review [1], our today's understanding of these phenomena is presented. In particular, the possibility was considered [2] of the extraction of these functions from the data of COMPASS and HERMES experiments. The HERMES data on Sivers's effect for the pions and kaons indicate a significant contribution of the sea quarks in the region $x \simeq 0,15$. We proposed a new fit of these data which includes all appropriate sea quark distributions; however, in spite of a statistically satisfactory description of all data the fit does not give an ideal description of the HERMES data on K^+ asymmetry. We believe that the measurements of the Sivers effect for the pions and kaons in CLAS12 and COMPASS will explain this situation.

In the general case, there are eight leading (twist-2) quark distribution functions dependent on spin, the fraction of longitudinal momentum x and the quark transversal momentum $\mathbf{p}_T^2$ which describe the spin structure of the nucleon

$$f_1^a,\ f_{1T}^{\perp a},\ g_{1L}^a,\ g_{1T}^a,\ h_{1T}^a,\ h_{1L}^{\perp a},\ h_{1T}^{\perp a},\ h_1^{\perp a}, \tag{1}$$

where the indices L(T) indicate the longitudinal (transverse) polarization of nucleon. Integration over $\mathbf{p}_T^2$ leaves only three independent "collinear"distribution functions: unpolarized $f_1^a(x)$, helicity $g_1^a(x)$ and transversity

$$h_1^a(x) = \int \mathrm{d}^2\mathbf{p}_T\{h_{1T}^a(x,\mathbf{p}_T^2) + \mathbf{p}_T^2/(2M_N^2)h_{1T}^{\perp a}(x,\mathbf{p}_T^2)\}.$$

The cross section differential in the azimuthal angle ϕ of the produced hadron has schematically the following general decomposition:

$$\begin{aligned}\frac{\mathrm{d}\sigma}{\mathrm{d}\phi} &= F_{UU}+\cos(2\phi)F_{UU}^{\cos(2\phi)}+S_L\sin(2\phi)F_{UL}^{\sin(2\phi)}+\lambda\left[S_LF_{LL}+S_T\cos(\phi-\phi_S)F_{LT}^{\cos(\phi-\phi_S)}\right]\\ &+S_T\left[\sin(\phi-\phi_S)F_{UT}^{\sin(\phi-\phi_S)}+\sin(\phi+\phi_S)F_{UT}^{\sin(\phi+\phi_S)}+\sin(3\phi-\phi_S)F_{UT}^{\sin(3\phi-\phi_S)}\right]\end{aligned} \tag{2}$$

In F_{XY}^{weight} the index $X = U(L)$ denotes the unpolarized (longitudinally polarized, helicity λ) beam; $Y = U(L,T)$ denotes the unpolarized target (longitudinally, transversely, with respect to the virtual photon, polarized target), ϕ and ϕ_S are azimuthal angles around virtual photon direction counted from the lepton scattering plane.

Each structure function arises from a different TMD. The chirally even f's and g's enter into the observables in connection with the unpolarized fragmentation function D_1^a, the chirally odd h's in connection with the chirally odd Collins fragmentation function $H_1^{\perp a}$

$$F_{UU} \propto \sum_a e_a^2\, f_1^a \otimes D_1^a, \qquad F_{LT}^{\cos(\phi-\phi_S)} \propto \sum_a e_a^2\, g_{1T}^{\perp a} \otimes D_1^a\,, \tag{3}$$

$$F_{LL} \propto \sum_a e_a^2\, g_1^a \otimes D_1^a, \qquad F_{UT}^{\sin(\phi-\phi_S)} \propto \sum_a e_a^2\, f_{1T}^{\perp a} \otimes D_1^a\,, \tag{4}$$

$$F_{UU}^{\cos(2\phi)} \propto \sum_a e_a^2\, h_1^{\perp a} \otimes H_1^{\perp a}, \qquad F_{UT}^{\sin(\phi+\phi_S)} \propto \sum_a e_a^2\, h_1^a \otimes H_1^{\perp a}\,, \tag{5}$$

$$F_{UL}^{\sin(2\phi)} \propto \sum_a e_a^2\, h_{1L}^{\perp a} \otimes H_1^{\perp a}, \qquad F_{UT}^{\sin(3\phi-\phi_S)} \propto \sum_a e_a^2\, h_{1T}^{\perp a} \otimes H_1^{\perp a}\,. \tag{6}$$

More precisely, TMDs and fragmentation functions enter into the respective *tree-level* expressions in certain convolution integrals, indicated by $\otimes$ in (3-6), which entangle transverse parton momenta from TMDs and fragmentation functions.

Early studies of the spin asymmetries on longitudinally polarized target were based on assumptions about $H_1^{\perp}(x)$ which are not supported by results of extraction of Collins function from deeply inelastic processes on transversally polarization target (COMPASS and HERMES) and e^+e^--annihilation (BELLE). These new data give the beginning of a new consistent picture of $H_1^{\perp}$ [3], which makes it necessary to reexamine our description of processes on longitudinal polarized target. In the work [4], we made this for the asymmetry of $A_{UL}^{\sin 2\phi} \propto \sum_a e_a^2 h_{1L}^{\perp(1)} H_1^{\perp}$ from the point of view the specific problem: how accurate are the approximate relations between different spin distribution functions?

QCD equations of motion connect different distribution functions responsible for different asymmetries in (2) with additional terms of pure twist-3. Neglecting of these terms leads to approximate relations of Wandzura-Wilczek (WW) type, in particular, between $h_{1L}^{\perp(1)}$ and transversity h_1

$$h_{1L}^{\perp(1)a}(x) \overset{!?}{\approx} -x^2 \int_x^1 \frac{\mathrm{d}y}{y^2}\, h_1^a(y)\,, \tag{7}$$

Our study shows that the data do not exclude a possibility that such approximations of WW-type do actually work (Fig. 1). As a by-product it is obtained that the data on two asymmetries due to the Collins effect, $A_{UL}^{\sin 2\phi}$ and $A_{UT}^{\sin(\phi+\phi_S)}$, are noncontradictory. For a more definite conclusion, more precise data on these asymmetries, preferably in the region $x \sim 0.3$, are necessary. Increase in statistics by an order of magnitude, expected from experiments CLAS and COMPASS, will improve our understanding of these and other spin asymmetries and shed light on the spin-orbital correlations of quarks in the nucleon.

In [5], we study the transverse moment dependent parton distribution function of the leading twist $h_{1T}^{\perp}$ which is sometimes called "pretzelosity". So strange name is connected with the fact that this function determines the form of distribution of transversely the

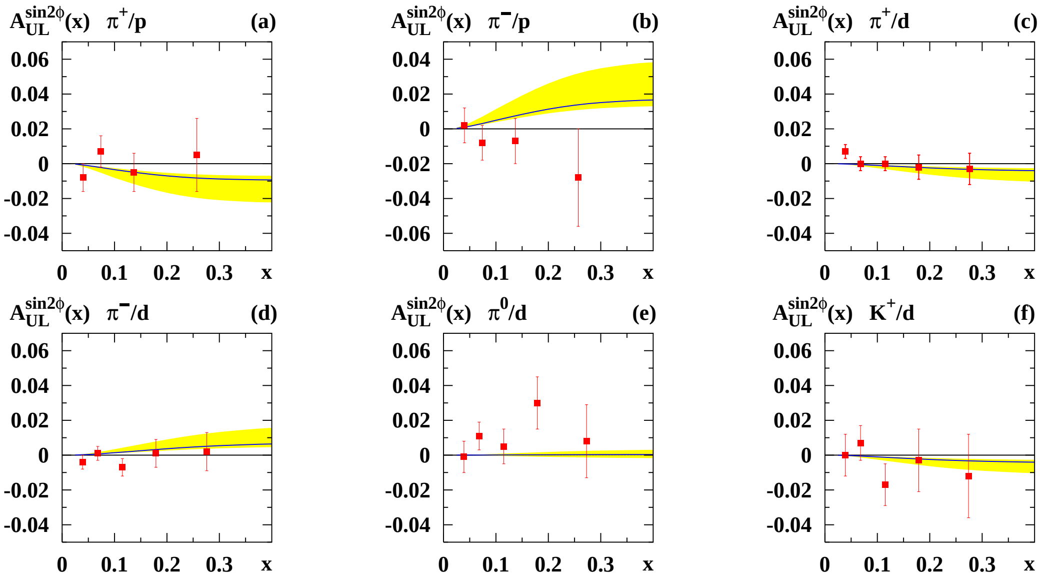

Рис. 1: Single spin asymmetries $A_{UL}^{\sin 2\phi}$ as a function of x for proton (a, b) and deuterium (c-f). Data by HERMES collaboration. Theoretical curves are obtained with the use of information about the Collins fragmentation function from [3] and validity of Wandzura-Wilczek type approximation Eq. (7). The shaded regions are due to uncertainty of Collins fragmentation function.

polarized partons in transversely polarized nucleon which in certain cases can take very specific forms.

The theoretical properties of this function and its prediction in the bag model, supplemented by a detailed comparison with the spectator model results are considered. For the first time, we notice an interesting relation valid in the majority of the relativistic models – the difference between the helicity distribution function and transversity, which frequently serves as a measure of relativistic effects in the nucleon, is nothing more than the pretzelosity distribution:

$$\boxed{g_1^q(x) - h_1^q(x) = h_{1T}^{\perp(1)q}(x)\,.} \tag{8}$$

The pretzelosity is chiral-odd and can be measured in combination with the Collins fragmentation in semi-inclusive deep inelastic scattering on transversely polarized target, where it leads to the azimuthal asymmetry proportional to $\sin(3\phi - \phi_S)$. In this connection, preliminary data of COMPASS collaboration on this asymmetry on deuteron target are discussed (Fig. 2a,b), and on this base the predictions are given for future experiments on different targets in JLab, COMPASS and HERMES. This work was presented and discussed at the 2nd International Conference on Spin Phenomena on transversely polarized target [6].

The pretzelosity distribution functions $h_{1T}^{\perp}$ are found also [7] in the covariant quark-parton model, proposed by us earlier, which describes the structure of nucleon from the point of view of three-dimensional internal motion of quarks in nucleon. It is shown that this model confirms relation (8) but without assumption about SU(6) spin-flavor symmetry. On this basis the prediction for CLAS12 experiments on the semi-inclusive deeply inelastic scattering on transversely polarized proton target are obtained (Fig. 2c).

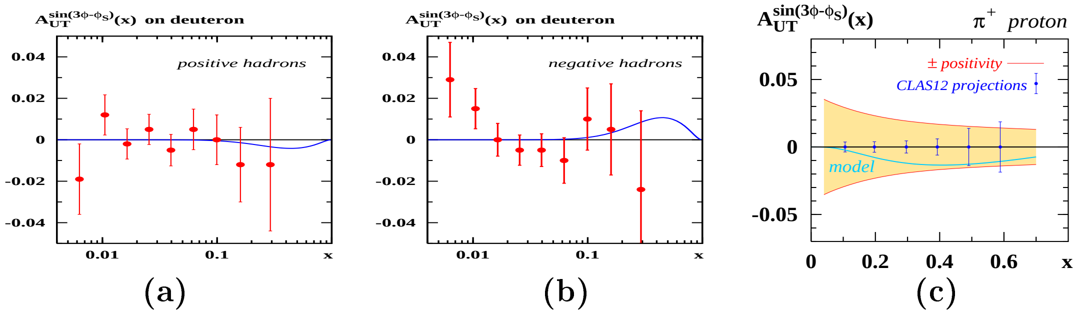

Рис. 2: Single spin asymmetries $A_{UT,\pi}^{\sin(3\phi-\phi_S)}(x)$ for the semi-inclusive process $l+d\uparrow\to l'+h^{\pm}+X$ on transversely polarized deuterium target in comparison with preliminary data of COMPASS for positive **(a)** and negative **(b)** hadrons.
(c). The same for π^+ production on the protons in kinematics of CLAS12 (projected accuracy is indicated by error bars). The shaded region indicate the range, allowed by the positivity condition.

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DISPERSION RELATIONS AND QCD FACTORIZATION IN HARD REACTIONS

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Introduction. We study analytical properties of the hard exclusive process amplitudes are studied. It is found that QCD factorization for deeply virtual Compton scattering and hard exclusive vector meson production results in the subtracted dispersion relation with the subtraction constant determined by the Polyakov-Weiss D-term. Hard exclusive reactions described by the Generalized Parton Distributions (GPDs) are the subject of extensive theoretical and experimental studies [1]- [6]. The analytical properties of deeply virtual Compton scattering (DVCS) and hard exclusive vector meson production (VMP) amplitudes constitute the important aspect of these studies [7]- [10]. The crucial point in application of the relevant dispersion relations is a possible ambiguity due to the subtraction constants. In these notes we discuss the problem of dispersion relations and subtractions in the framework of the leading order QCD factorization [7, 11].

Dispersion relations in skewness-plane. We restrict our study to the case: $\{s,\, Q^2\} \to \infty, t \ll \{s,\, Q^2\}$, where QCD factorization is applicable. At the leading order, DVCS and VMP amplitudes can be expressed via

$$\mathcal{A}_f(\xi,\, t) = \lim_{\epsilon\to 0} \mathcal{A}_f(\xi - i\epsilon,\, t) = \lim_{\epsilon\to 0} \int\limits_{-1}^{1} dx\, \frac{H_f^{(+)}(x,\, \xi,\, t)}{x - \xi + i\epsilon}. \tag{1}$$

Here $H^{(+)}(x,\, \xi,\, t)$ denotes the singlet ($C = +1$) combination of GPDs summing the contributions of quarks and anti-quarks and of s- and u-channels.

To prove the analyticity of the amplitude for $|\xi| > 1$, one represents the denominator as the geometric series [7]:

$$\mathcal{A}(\xi) = -\sum_{n=0}^{\infty} \xi^{-n-1} \int\limits_{-1}^{1} dx\, H^{(+)}(x,\, \xi) x^n. \tag{2}$$

This series is convergent thanks to the polynomiality condition which reflects the Lorentz invariance [4]. One can write down the fixed-t dispersion relations for the LO amplitude:

$$\mathrm{Re}\, \mathcal{A}(\xi) = \frac{v.p.}{\pi} \int\limits_{-1}^{1} dx\, \frac{\mathrm{Im}\mathcal{A}(x + i\epsilon)}{x - \xi} + \Delta(\xi), \tag{3}$$

where the limit $\epsilon \to 0$ is implied. In other words, we write

$$v.p. \int\limits_{-1}^{1} dx\, \frac{H^{(+)}(x,\, \xi)}{x - \xi} = v.p. \int\limits_{-1}^{1} dx\, \frac{H^{(+)}(x,\, x)}{x - \xi} + \Delta(\xi). \tag{4}$$

Here $\Delta(\xi)$ is a possible subtraction. This expression represents the holographic property of GPD: the relevant information about hard exclusive amplitudes in the considered leading approximation is contained in the one-dimensional sections $x = \pm\xi$ of the two dimensional space of x and ξ. These holographic as well as tomographic properties in momentum space are complementary to the often discussed holography and tomography in coordinate space.

Prove now that $\Delta(\xi)$ is finite and independent of ξ. To this end, one considers the following representation [7]:

$$\Delta(\xi) = v.p. \int_{-1}^{1} dx\, \frac{H^{(+)}(x,\xi) - H^{(+)}(x,x)}{x-\xi} = -v.p. \int_{-1}^{1} dx \sum_{n=1}^{\infty} \frac{1}{n!} \frac{\partial^n}{\partial \xi^n} H^{(+)}(x,\xi)\Big|_{\xi=x} (\xi - x)^{n-1}.$$

Due to the polynomiality condition the only survived highest power term in this series is equal to a finite subtraction constant: $\Delta(\xi) \equiv \Delta$. This can also be derived with the use of the Double Distributions (DDs) formalism [7, 11]. It should be emphasized that both integrals in (4) are divergent at $\xi = t = 0$, and these divergences do not cancel for $\xi \to 0$. Thus, Δ cannot be defined for $\xi = t = 0$. However, the point $\xi = 0$ (or $Q^2 = 0$) cannot be accessed in DVCS and VMP experiments. Note that for an arbitrarily small ξ the integrals are finite and, therefore, Δ is well-defined.

Taking into account the parameterization for the D-term as [4]

$$D(\beta) = (1-\beta^2) \sum_{n=0}^{\infty} d_n C_{2n+1}^{(3/2)}(\beta), \tag{5}$$

and keeping only the lowest term one gets $\Delta = -4d_0$. This lowest term d_0 was estimated within the framework of different models. We focus on the results of chiral quark-soliton model [12]: $d_0^{\mathrm{CQM}}(N_f) = d_0^u = d_0^d = -\frac{4.0}{N_f}$, where N_f is the number of active flavours and lattice simulations [13]: $d_0^{\mathrm{latt}} = d_0^u \approx d_0^d = -0.5$. The subtraction constant varies as $\Delta^p_{\mathrm{CQM}}(2) = \Delta^n_{\mathrm{CQM}}(2) \approx 4.4$, $\Delta^p_{\mathrm{latt}} \approx \Delta^n_{\mathrm{latt}} \approx 1.1$ for the DVCS on both the proton and neutron targets.

Dispersion relations in ν-plane. Let us now compare the dispersion relation with the dispersion relation written in ν-plane where $\nu = (s-u)/4m_N$. In terms of new variables ν', ν related to x, ξ as $x^{-1} = 4m_N\nu'/Q^2$, $\xi^{-1} = 4m_N\nu/Q^2$. The fixed-t dispersion relation becomes the subtracted one:

$$\mathrm{Re}\,\mathcal{A}(\nu, Q^2) = \frac{v.p.}{\pi} \int_{Q^2/4m_N}^{\infty} d\nu'^2 \mathrm{Im}\,\mathcal{A}(\nu', Q^2) \left[\frac{1}{\nu'^2 - \nu^2} - \frac{1}{\nu'^2}\right] + \Delta. \tag{6}$$

This subtracted (at the symmetric unphysical point $\nu = 0$) dispersion relation is our principal result. It can be considerably simplified provided $\mathrm{Im}\,\mathcal{A}(\nu)$ decreases fast enough so that both terms in the squared brackets can be integrated separately:

$$\mathrm{Re}\,\mathcal{A}(\nu) = \frac{v.p.}{\pi} \int_{\nu_0}^{\infty} d\nu'^2\, \frac{\mathrm{Im}\,\mathcal{A}(\nu')}{\nu'^2 - \nu^2} + \mathbf{C}_0\,, \tag{7}$$

where

$$\mathbf{C}_0 = \Delta - \frac{1}{\pi}\int\limits_{\nu_0}^{\infty} d\nu'^2\, \frac{\mathrm{Im}\,\mathcal{A}(\nu')}{\nu'^2} = \Delta + \int\limits_{-1}^{1} dx\, \frac{H^{(+)}(x,\, x)}{x}. \tag{8}$$

Using (4) with $\xi = 0$,

$$\Delta(t) \;\; = \;\; 2\int\limits_{-1}^{1} dx\, \frac{H(x,\, 0, t) - H(x,\, x, t)}{x}, \tag{9}$$

one can see that the D-term is cancelled from the expression for the subtraction constant

$$\mathbf{C}_0(t) = 2\int\limits_{-1}^{1} dx \frac{H(x,\, 0,\, t)}{x}. \tag{10}$$

This constant is similar to the result obtained in the studies of the fixed pole contribution to the forward Compton amplitude. For ξ, $t = 0$ GPDs are expressed in terms of standard parton distributions

$$\mathbf{C}_0(0) = 2\int\limits_{0}^{1} dx \frac{q(x) + \bar{q}(x)}{x} = 2\int\limits_{0}^{1} dx \frac{q_v(x) + 2\bar{q}(x)}{x}. \tag{11}$$

However, the integrals in (9) and (10) diverge at low x in both the valence and sea quark contributions [10, 11]. Therefore, for $t = 0$ we should consider (6) as a correct general form of the dispersion relation which includes an infinite subtraction at the point $\nu = 0$ and the subtraction constant associated with the D-term [11].

For $t \neq 0$, the integrals in (9) and (10) converge for sufficiently large t [11]. In the case of Regge inspired parameterization [4] $H(x,\, 0,\, -t) \sim x^{-\alpha(0)+\alpha' t}$, this condition reads as $t > \alpha(0)/\alpha'$ resulting in $t \gtrsim 1(10)\mathrm{GeV}^2$ for the valence (sea) quark distributions.

The subtracted dispersion relation for the forward Compton scattering amplitude is

$$\mathrm{Re}\,\mathcal{A}(\nu) = \frac{\nu^2}{\pi} v.p. \int\limits_{0}^{\infty} \frac{d\nu'^2}{\nu'^2}\, \frac{\mathrm{Im}\,\mathcal{A}(\nu')}{(\nu'^2 - \nu^2)} + \Delta. \tag{12}$$

One can see that, for the proton target, our subtraction combined with the lattice simulations is rather close to the low energy Thomson term (note that $\Delta_{\mathrm{Thomson}} = 1$ for our normalization of the Compton amplitude).

Note that approaching the value $Q^2 = 0$ for Compton amplitude involves an infinite tower of higher twists. The observed numerical coincidence may signal about the sort of duality similar to Bloom-Gilman type duality between the leading twist contribution and the full result [14].

Conclusions. We can conclude that the fixed-t DR for the DVCS and VMP amplitudes require the infinite subtractions at $\nu = 0$ with the subtraction constants associated with the D-terms. For the productions of the mesons defined by valence ($C = -1$) GPDs the finite subtraction is absent. The appearance of the subtraction expressed in terms of (forward) parton distributions may be investigated in the framework of the leading order QCD factorization. The possibility of continuation of our results to the real photons limit has been considered.

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LOW ENERGY TESTS OF THE STANDARD MODEL

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The theoretical study of π^0 and η decaying into lepton pairs and the comparison with the experimental rates offers an interesting possibility to test our understanding of the long-distance dynamics of the standard model. The situation with these decays became more pressing after recent KTeV E799-II experiment at Fermilab in which the pion decay into an electron-positron pair was measured using the $K_L \to 3\pi$ process as a source of tagged neutral pions [1]. The branching ratio was determined to be equal to

$$B^{\mathrm{KTeV}}_{\mathrm{no-rad}}\left(\pi^0 \to e^+e^-\right) = (7.49 \pm 0.29 \pm 0.25) \cdot 10^{-8}. \tag{1}$$

The standard model prediction based on the use of CLEO data on the transition form factor $\pi \to \gamma\gamma^*$ [2] gives [3]

$$B^{\mathrm{Theor}}\left(\pi^0 \to e^+e^-\right) = (6.2 \pm 0.1) \cdot 10^{-8}, \tag{2}$$

which is 3.3σ below the KTeV result (1). Therefore, it is extremely important to trace possible sources of the discrepancy between the experiment and theory. There are a number of possibilities: (1) problems with (statistic) experiment procession, (2) inclusion of QED radiation corrections by KTeV is wrong, (3) unaccounted mass corrections are important, and (4) effects of new physics. At the moment the last possibilities were reinvestigated. In [4], the contribution of QED radiative corrections to the $\pi^0 \to e^+e^-$ decay, which must be taken into account when comparing the theoretical prediction (2) with the experimental result, (1) was revised. Comparing with earlier calculations [5], the main progress is in the detailed consideration of the $\gamma^*\gamma^* \to e^+e^-$ subprocess and revealing of dynamics of large and small distances. Occasionally, this number agrees well with the earlier prediction based on calculations [5] and, thus, the KTeV analysis of radiative corrections is confirmed. In the present paper, we show that the mass corrections are under control and do not resolve the problem. So our main conclusion is that the inclusion of radiative and mass corrections is unable to reduce the discrepancy between the theoretical prediction for the decay rate (2) and experimental result (1). The effects of new physics were considered in [6] where the excess of experimental data over theory is explained by the contribution of low mass (~ 10 MeV) vector bosons appearing in some models of dark matter. Another possibility is the contribution of light CP-odd Higgs entering into the next-to-minimal supersymmetric standard model [7]. Further independent experiment at KLOE, NA48, WASAa@COSY, BES III and other facilities will be crucial for resolution of the problem.

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RELATIONS BETWEEN SU(2) AND SU(3) LECs IN ChPT AT TWO-LOOP LEVEL

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We consider Green functions of quark currents in the framework of QCD with three flavours. At low energies, the Green functions can be analysed in the framework of chiral perturbation theory (ChPT) [1–3]. It is customary to perform the pertinent quark mass expansion either around $m_u = m_d = 0$, with the strange quark mass held fixed at its physical value (ChPT$_2$), or to consider an expansion in all three quark masses, around $m_u = m_d = m_s = 0$ (ChPT$_3$). The relevant effective lagrangians contain low-energy constants (LECs) which are not determined by chiral symmetry alone. The two expansions are not independent: one can express the LECs in the two-flavour case through the ones in ChPT$_3$. These relations were given at one-loop order in [3] and were used to obtain information on the LECs in ChPT$_3$ from those known in the two-flavour case. Because there are many two-loop calculations available now, both in the two- and three-flavour case [4, 5], it is expedient to have the relevant relations between the LECs at two-loop accuracy as well, both, to obtain more additional information, and for internal consistency checks.

The relations that occur at order p^2 and p^4 in ChPT$_2$ were found in [6]. The result is

$$\begin{aligned}
Y &= Y_0\left[1 + a_Y\,x + b_Y\,x^2 + O(x^3)\right], \qquad Y = F\,, \Sigma\,,\\
l_i^r &= a_i + x\,b_i + O(x^2)\,,\ \ i \neq 7\,,\\
l_7 &= \frac{F_0^2}{8B_0 m_s} + a_7 + x\,b_7 + O(x^2)\,,\\
x &= \frac{B_0 m_s}{N F_0^2}\,, \qquad N = 16\pi^2\,, \qquad \Sigma = F^2 B\,, \qquad \Sigma_0 = F_0^2 B_0\,.
\end{aligned}$$

We denote the contributions proportional to a_i (b_i) as NLO (NNLO) terms. Their explicit expressions are given in [6]

We have also obtained [7] such relations at leading order for the NNLO LECs c_i in absence of external sources $s,\ p$. Due to this restricted framework, only relations for LECs not involving monomials dependent on the sources s or p are nontrivial. Moreover, in the restricted framework, an additional relation among the remaining SU(2)-monomials has been found:

$$\begin{aligned}
&\frac{4}{3}P_1 - \frac{1}{3}P_2 + P_3 - \frac{10}{3}P_{24} + \frac{4}{3}P_{25} + 2P_{26} - \frac{8}{3}P_{28} - \frac{1}{2}P_{29} + \frac{1}{2}P_{30} - P_{31} + 2P_{32} - \frac{1}{2}P_{33}\\
&+\frac{4}{3}P_{36} - \frac{4}{3}P_{37} - \frac{11}{6}P_{39} + \frac{5}{6}P_{40} + \frac{7}{3}P_{41} - \frac{4}{3}P_{42} - \frac{3}{2}P_{43} + \frac{1}{2}P_{44} - \frac{1}{2}P_{45} - P_{51} - P_{53} = 0\,.
\end{aligned}$$

The relation is no longer valid in the full framework, since the EOM is different there. We used it to exclude the monomial P_1 from the consideration. The monomial P_{27} can be discarded also, as described in [8]. As a result, we give the matching for the 27 combinations of c_i^r. In the full framework, an additional relation (apart from the ones for the monomials involving the sources s and p) for c_1^r would be present.

To render the formulae more compact, we found it convenient to express the bare kaon mass squared $B_0 m_s$ through its equivalent $\bar{M}_K^2$ in the chiral limit. Then, the final result may be written in the form

$$x_i = p_i^{(0)} + p_i^{(1)} \ln(\bar{M}_K^2/\mu^2) + p_i^{(2)} \ln^2(\bar{M}_K^2/\mu^2)\,. \tag{1}$$

Due to the relations among the SU(2) monomials, the x_i on the left hand side are a combination of c_1^r, c_{27}^r and one other c_i^r, given in Table . The explicit expressions for the polynomials $p_i^{(n)}$ are displayed in the paper [7].

i	x_i	i	x_i	i	x_i
1	$c_2^r + \frac{1}{4}c_1^r$	10	$c_{32}^r - \frac{3}{2}c_1^r - c_{27}^r$	19	$c_{43}^r + \frac{9}{8}c_1^r + \frac{1}{4}c_{27}^r$
2	$c_3^r - \frac{3}{4}c_1^r$	11	$c_{33}^r + \frac{3}{8}c_1^r + \frac{1}{4}c_{27}^r$	20	$c_{44}^r - \frac{3}{8}c_1^r - \frac{1}{4}c_{27}^r$
3	$c_{24}^r + \frac{5}{2}c_1^r$	12	$c_{36}^r - c_1^r$	21	$c_{45}^r + \frac{3}{8}c_1^r + \frac{1}{4}c_{27}^r$
4	$c_{25}^r - c_1^r$	13	$c_{37}^r + c_1^r$	22	c_{50}^r
5	$c_{26}^r - \frac{3}{2}c_1^r$	14	c_{38}^r	23	$c_{51}^r + \frac{3}{4}c_1^r + \frac{1}{2}c_{27}^r$
6	$c_{28}^r + 2c_1^r - c_{27}^r$	15	$c_{39}^r + \frac{11}{8}c_1^r + \frac{1}{4}c_{27}^r$	24	c_{52}^r
7	$c_{29}^r + \frac{3}{8}c_1^r + \frac{1}{4}c_{27}^r$	16	$c_{40}^r - \frac{5}{8}c_1^r - \frac{1}{4}c_{27}^r$	25	$c_{53}^r + \frac{3}{4}c_1^r + \frac{1}{2}c_{27}^r$
8	$c_{30}^r - \frac{3}{8}c_1^r - \frac{1}{4}c_{27}^r$	17	$c_{41}^r - \frac{7}{4}c_1^r - \frac{1}{2}c_{27}^r$	26	c_{55}^r
9	$c_{31}^r + \frac{3}{4}c_1^r + \frac{1}{2}c_{27}^r$	18	$c_{42}^r + c_1^r$	27	c_{56}^r

Таблица 1: The quantities x_i of Eq. (1)

Following the strategy outlined in [6,7], we investigated three-flavour chiral perturbation theory including virtual photons in a limit where the strange quark mass is much larger than the external momenta and the up and down quark masses, and where the external fields are those of two-flavour chiral perturbation theory [9]. As a result we worked out the strange quark mass dependence of the LECs C, k_i in the two-flavour case. We expect that the relations may be useful for a more precise determination of the pertinent LECs.

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B-MESON FORM FACTORS

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Transition form factors that characterise the decays of B-mesons into light pseudoscalar and vector mesons, the so-called heavy-to-light decays, are basic to an understanding of this heavy-meson exclusive semileptonic and rare radiative decays. These form factors also provide the factorisable amplitudes that appear in B-meson exclusive nonleptonic charmless decays. An understanding of all these processes is essential to the reliable determination of CKM matrix elements, and transitions mediated by electroweak and gluonic penguin operators. Moreover, they should provide a means of searching for non Standard Model effects and CP violation. Considering all these factors, it is not surprising that heavy-light form factors are the subject of much experimental and theoretical scrutiny.

The analysis of heavy-to-light processes has two facets. One is factorisation; viz., the feature that in exclusive decays of B-mesons there exist strong interaction effects that do not correspond to form factors. These may be radiative corrections to purely hadronic operators in the weak effective Lagrangian or final state interactions between daughter hadrons. The development of soft collinear effective-field theory (SCET) provides a means of simplifying that problem, yielding factorisation theorems which enable a systematic approximation to be developed for a given process in terms of products of soft and hard matrix elements (see, e.g. [1]).

The second facet, once factorisation for a given process is assumed or proved, is to evaluate the hadronic transition form factors. Naturally, they cannot be calculated in perturbation theory. The relevant matrix elements involve single hadrons in the initial and final states. Hence, their calculation requires information about the structure of both heavy- and light-mesons. A variety of theoretical approaches have been applied to this problem, recent amongst which are analyses using light-cone sum rules [2, 3], light-front quark models [4], a constituent-quark model in a dispersion relation formulation [5], and relativistic quark models – e.g., Ref. [6].

Models based on results obtained via QCD's Dyson-Schwinger equations (DSEs) have also been employed [7]. The calculated $B \to P(V)$ heavy-to-light transitions are shown in Fig. 1 for $q^2 \in [0, q^2_{\rm max}]$.

It is noteworthy and phenomenologically important that in our DSE-based approach all form factors can be calculated on the entire domain of physically accessible momenta. Moreover, the chiral limit is directly accessible and the consequences of Goldstone's theorem are manifest, so that both pseudoscalar and vector light-quark mesons are realistically described. No extrapolation in any quantity is required.

Our calculated results are satisfactorily *interpolated* by the simple function

$$F(q^2) = \frac{F(0)}{1 - as + bs^2}, \qquad s = q^2/m_B^2. \tag{1}$$

We list the values of the form factors at the maximum recoil point, $q^2 = 0$, and the parameters a and b in Tables 1 and 2. Analytically, $a_0(0) = a_+(0) = g(0)$, a result preserved by the interpolation function. We provide the interpolations so that our results may readily be adapted as input for other analyses.

	F_+	F_-	F_T	A_0	A_+	A_-	V	a_0	a_+	g
$F(0)$	0.24	-0.24	0.24	0.32	0.25	-0.32	0.32	0.25	0.26	0.26
a	1.87	1.97	1.92	1.16	2.08	2.27	2.21	1.26	2.10	2.21
b	0.93	1.04	1.00	0.32	1.14	1.38	1.30	0.48	1.16	1.29

Таблица 1: $B \to \pi(\rho)$ form factors: values of the parameters in Eq. (1).

	F_+	F_-	F_T	A_0	A_+	A_-	V	a_0	a_+	g
$F(0)$	0.29	-0.28	0.32	0.40	0.30	-0.38	0.37	0.30	0.30	0.30
a	1.85	1.95	1.90	0.98	1.92	2.10	2.05	1.04	1.95	2.05
b	0.96	1.09	1.02	0.034	0.97	1.19	1.13	0.16	1.00	1.12

Таблица 2: $B \to K(K^*)$ form factors: values of the parameters in Eq. (1).

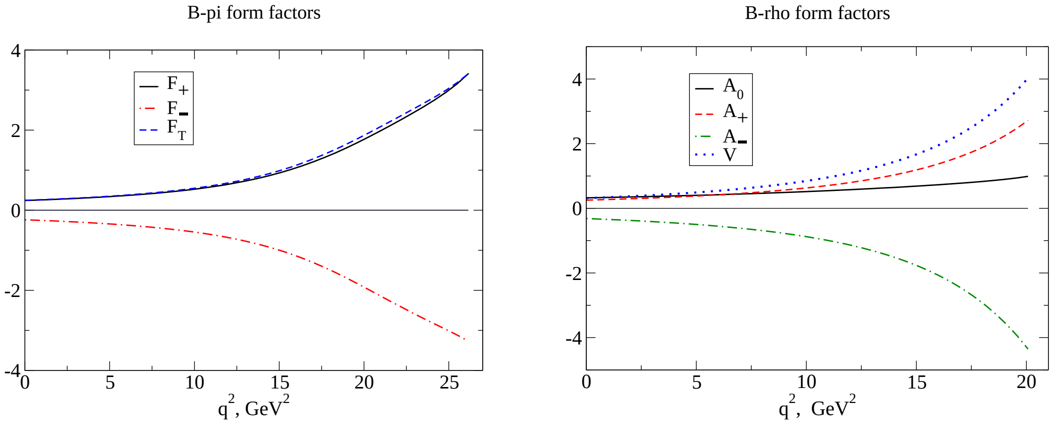

Рис. 1: Our results for the $B \to \pi, \rho$ form factors.

In Table 3 we collect our predictions for the form factors at the maximum recoil point and provide a comparison with extant results obtained within other frameworks. The figures and tables highlight the wide range of phenomena accessible within our approach.

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	DSE [7]	LCSR [2]	LCSR [3]	LCQM [4]	DQM [5]	RQM [6]
$f^{+}_{B\pi}(0)$	0.24	0.25±0.05	0.258±0.031	0.25	0.29	0.22
$f^{+}_{BK}(0)$	0.30	0.31±0.04	0.331±0.041	0.30	0.36	
$f^{T}_{B\pi}(0)$	0.25	0.21±0.04	0.253±0.028	0.25	0.28	
$f^{T}_{BK}(0)$	0.32	0.27±0.04	0.358±0.037	0.33	0.35	
$V^{B\rho}(0)$	0.31	0.32±0.10		0.30	0.31	0.30
$V^{BK^*}(0)$	0.37	0.39±0.11		0.34	0.44	
$A_1^{B\rho}(0)$	0.24	0.24±0.08		0.23	0.26	0.27
$A_1^{BK^*}(0)$	0.29	0.30±0.08		0.25	0.36	
$A_2^{B\rho}(0)$	0.25	0.21±0.09		0.22	0.24	0.28
$A_2^{BK^*}(0)$	0.30	0.26±0.08		0.23	0.32	
$T_1^{B\rho}(0)$	0.26	0.28±0.09		0.26	0.27	
$T_1^{BK^*}(0)$	0.30	0.33±0.10		0.29	0.39	

Таблица 3: Our calculated values of $B \to \pi, K$ and $B \to \rho, K^*$ form factors at the maximum recoil point compared with the results obtained by other authors. Based on the widespread application of our approach herein and elsewhere, we estimate that the relative systematic uncertainty in our calculated results is $\sim 15\%$.

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EXOTIC HADRON STATES: NEWS FROM THEORY AND EXPERIMENT

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The constituent quark models are widely used for the description of hadron properties. In the framework of this approach most of the observed meson states are quark-antiquark bound states and baryons are a three-quark system. However, there no evident reasons to forbid the existence of so-called exotic states. For example, in various versions of the constituent quark model the multiquark states with the number of quark and antiquark more than three should exist as well. Moreover, the quark-gluon hybrid states and glueballs which include valence gluons are under discussion now. There are two types of exotic states. Hidden exotic states can have the same quantum numbers as the ordinary hadrons. Open exotic states has quantum numbers which are impossible to obtain within the quark-antidiquark and three-quark model for hadrons.

Recently, the development of the exotic spectroscopy was related mainly with attempts to discribe properties of θ^+ pentaquark which was expected to have a small width (about 15 MeV) and a small mass (about 1540 MeV) which has been predicted within the coliton model for baryons [1]. Within the constituent quark model a state like this is a bound state of two ud diquarks and one strange antiquark. Unfortunately, experimental situation around this state is highly controversial. Some of the experimental groups [2], including JINR experimental groups [3], report observation of this state, but other high statistics experiments (see, for example [4]) have not seen this sort of resonance.

We should like to emphasize that within the quark model it is rather difficult to explain the modern experimental restriction on the width of this resonance, $\Gamma < 1 MeV$ [2]. Furthermore, precise calculation of the mass of θ^+ within QCD sum rules [5] gives a larger value of its mass with respect to the soliton model prediction and shows a very weak signal for the bound state. Within the soliton model θ^+ is a member of a flavor antidecuplet. Therefore, if this model is correct, other members of the antidecuplet should exist as well. At the present time, the candidate for a nonstrange pentaquark $N^*(1685)$ is under discussion [6].

Increasing attention is now given to the problem of four-quark states called tetraquarks. Interest in these states is related to the necessity to explain the scalar meson spectrum, which does not follow by prediction of the naive quark-antiquark model. The central problem here is the sigma($f_0(600)$)- meson which probably has a very complicated internal structure. For a long time even the existence of state like this was doubtful because the pion-pion scattering phase does not change by 90° at resonance. The problem has been solved in the recent papers by Achasov with collaborators [7]. They show that within the sigma-model the sigma-pole contribution is hidden in the large background amplitude of pion-pion scattering. At the present time, sigma-meson is considered as a good established resonance with the mass around 440 MeV and the width about 540 MeV [8]. From the theoretical point of view, the sigma-meson may include a large mixture of a four-quark state [9] or/and glueball [10], [11]. Furthermore, the properties of sigma-meson in quark-gluon plasma (QGP) and in vacuum might be different. This observation opens a new way to investigate the properties of QGP through changing of properties of sigma-meson which is produced in heavy ion collisions [12], [13].

Very interesting bound states predicted within different QCD based approaches are hybrids and quark-gluon bound states. The famous candidate for such a hybrid is the $\pi(1600)$ state with exotic quantum numbers, $J^{PC} = 1^{-+}$. The evidence for $\pi(1600)$ was obtained for first time by the VES Collaboration at Protvino [14] and recently the search for this state was continued by the E852 Collaboration at Brookhaven and by CLAS at CEBAF. The result of the analysis of data coming from these experiments is rather controversial [15], [16]. Therefore, the intensive search for hybrids is being continued in several current experiments, including COMPASS experiment at CERN.

Glueball states are one of the firm predictions of QCD and their properties are studied in different approaches based on QCD, for example, within lattice QCD and QCD sum rules (see review [17]). The main activity in this field is related to the investigation of low mass glueball states with zero spin and quantum numbers $J^{PC} = 0^{\pm 00}$ and to tensor glueball, $J^{PC} = 2^{++}$. Recent calculations show significant mixing of zero spin glueballs with ordinary quarkonium states and, therefore, the ambiguity problem of theoritical interpretation of the experimental data for such states arises. From our point of view, a cleaner glueball channel is a tensor channel, where the mixing with quark-antiquark states is expected to be very small.

In conclusion, we would like to mention the large numbers of exotic candidates, so-called XYZ mesons, with charm quark content, which were found recently in BES-II, BELLA and BaBar experiments. Most of such states have unexpected values of the masses and widths [18]. Investigation of the hadron exotic is included also in future experiments: PANDA (FAIR), GlueX (CEBAF) and BES-III.

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A RELATIVISTIC MEAN-FIELD MODEL WITH SCALED HADRON MASSES AND COUPLINGS

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The study of thermodynamic properties and phase structure of strongly interacting nuclear matter at high baryon densities and temperatures has recently become important in view of plans to construct new mahines: NICA (JINR, Dubna) [1], the Nuclotron-based collider for searching a quark-hadron mixed phase, and FAIR (GSI Darmstadt), a fixed-target accelerator covering the same heavy-ion energy range of (5-40) AGeV [2], for investigating bulk nuclear properties and rare decays, as well as a low-energy campaign at RHIC [3] aimed at identification of the critical end-point.

Following Ref. [4] we assume here a relevance of the (partial) chiral symmetry restoration at high baryon densities and/or temperatures [5] manifesting in form of the Brown-Rho scaling hypothesis [6]: masses and coupling constants of all hadrons decrease with a density increase in approximately the same way. Our main goal is to construct some effective thermodynamic model of Equation of State (EoS) that would incorporate the decrease of hadron masses and couplings with increase of the baryon density n_B and temperature T and, simultaneously, would fulfill various constraints known from analysis of atomic nuclei, neutron stars and HIC. Our consideration is based on the generalization of the so-called KVOR model [7] to finite temperatures [8,9] named here the Scaled Hadron Mass and Coupling (SHMC) model.

Within our relativistic mean-field SHMC model we present the Lagrangian density of the isospin symmetric hadronic matter as the sum of several terms:

$$\mathcal{L} = \mathcal{L}_{\rm bar} + \mathcal{L}_{\rm MF} + \mathcal{L}_{\rm ex} \ . \tag{1}$$

For baryons interacting via σ, ω mean fields the Lagrangian density is as follows, cf. [4],

$$\mathcal{L}_{\rm bar} \;=\; \sum_{b\in\{b\}} \left[i\bar{\Psi}_b \left(\partial_\mu - i\, g_{\omega b}\, \chi_\omega\ \omega_\mu\right)\gamma^\mu\, \Psi_b - m_b^*\, \bar{\Psi}_b\, \Psi_b \right]. \tag{2}$$

The considered baryon set is $\{b\} = N(938)$, $\Delta(1232)$, $\Lambda(1116)$, $\Sigma(1193)$, $\Xi(1318)$, $\Sigma^*(1385)$, $\Xi^*(1530)$, and $\Omega(1672)$. The used σ-field dependent effective masses of baryons are [4, 8, 9]

$$m_b^*/m_b = \Phi_b(\chi_\sigma \sigma) = 1 - g_{\sigma b}\ \chi_\sigma\ \sigma/m_b\,,\ b\in\{b\}\ . \tag{3}$$

In (2), (3) $g_{\sigma b}$ and $g_{\omega b}$ are coupling constants and $\chi_\sigma(\sigma)$, $\chi_\omega(\sigma)$ are coupling scaling functions.

The σ-, ω-meson mean-field contribution

$$\mathcal{L}_{\rm MF} \;=\; \frac{\partial^\mu\sigma\ \partial_\mu\sigma}{2} - \frac{m_\sigma^{*2}\,\sigma^2}{2} - U(\sigma) + \frac{m_\omega^{*2}\,\omega_0^2}{2} - \frac{1}{4}F^{\mu\nu}F_{\mu\nu}\ , \quad \omega^\mu = (\omega^0, \vec{0}). \tag{4}$$

The mass terms of the mean fields are

$$m_m^*/m_m \;=\; |\Phi_m(\chi_\sigma\sigma)|\,, \quad \{m\} = \sigma, \omega\,. \tag{5}$$

The dimensionless scaling functions Φ_b and Φ_m, as well as the coupling scaling functions χ_m depend on the scalar field in the combination $\chi_\sigma(\sigma)\ \sigma$. Following [4] we assume an approximate validity of the Brown-Rho scaling ansatz in the simplest form

$$\Phi = \Phi_N = \Phi_\sigma = \Phi_\omega = \Phi_\rho = 1 - f,\ \ f = g_{\sigma N}\ \chi_\sigma\ \sigma/m_N\,. \tag{6}$$

We keep the standard form for the non-linear self-interaction term (potential U) of relativistic mean-field models, but now in terms of the new variable f:

$$U \ \ = \ \ m_N^4(\frac{b}{3}\,f^3 + \frac{c}{4}\,f^4). \tag{7}$$

The third term in the Lagrangian density (1) includes meson excitations

$$\mathcal{L}_{\rm ex} = \sum_{bos\in\{ex\}} \mathcal{L}_{ex}, \quad \{bos\} \ \ = \ \ \pi; K, \bar{K}; \eta(547); \sigma', \omega', \rho'; K^{*\pm,0}(892), \eta'(958), \phi(1020).$$

It is convenient to introduce the coupling ratios

$$x_{mb} = g_{mb}/g_{mN},\ \ m \in \{m\} = \sigma, \omega, \tag{8}$$

the renormalized constants

$$C_m = \frac{m_N\ g_{mN}}{m_m} \tag{9}$$

and, instead of χ_m, other variables

$$\eta_m(f) = \Phi_m^2(f)/\chi_m^2(f)\,. \tag{10}$$

In terms of these new variables the contribution of mean fields to the pressure is :

$$P_{\rm MF}[f,\omega_0] \ \ = \ \ -\frac{m_N^4\,f^2}{2\,C_\sigma^2}\,\eta_\sigma(f) - U(f) + \frac{m_N^2\eta_\omega(f)}{2\,C_\omega^2}\,[g_{\omega N}\,\chi_\omega\,\omega_0]^2. \tag{11}$$

The pressure of every particle is the usual pressure of ideal Bose-gas with effective masses and chemical potentials for σ', ω', ρ', η and K mesons and ideal Fermi gas for baryons. For other particles we use bare masses.

In order to get $P_\sigma^{\rm part}$, we should expand total pressure $P[\sigma,\omega_0(\sigma)]$ in $\sigma^{'} = \sigma - \sigma^{\rm cl}$. The term linear in $\sigma^{'}$ does not give a contribution due to subsequent requirement of the pressure minimum in $\sigma^{\rm cl}$. The quadratic term produces effective $\sigma^{'}$ particle mass squared,

$$(m_\sigma^{\rm part*})^2 = -\frac{d^2 P_{\rm MF}[\sigma,\omega_0(\sigma)]}{d\sigma^2} = -\frac{d^2 P_{\rm MF}[f,\omega_0(f)]}{df^2}\left(\frac{df}{d\sigma}\right)^2. \tag{12}$$

The first-order derivative $dP/df = 0$, as it follows from the equations of motion. Keeping only quadratic terms in all thermodynamical quantities in fluctuating fields we disregard the $P_{\rm bos.ex.}$ term in (12) and suppress the baryon contribution in (12).

Within our approximation effective masses of ω' and ρ' excitations prove to be the same as those follow from the mean-field mass terms

$$m_\omega^{\rm part*} = m_\omega|\Phi_\omega(f)|, \quad m_\rho^{\rm part*} = m_\rho|\Phi_\omega(f)|. \tag{13}$$

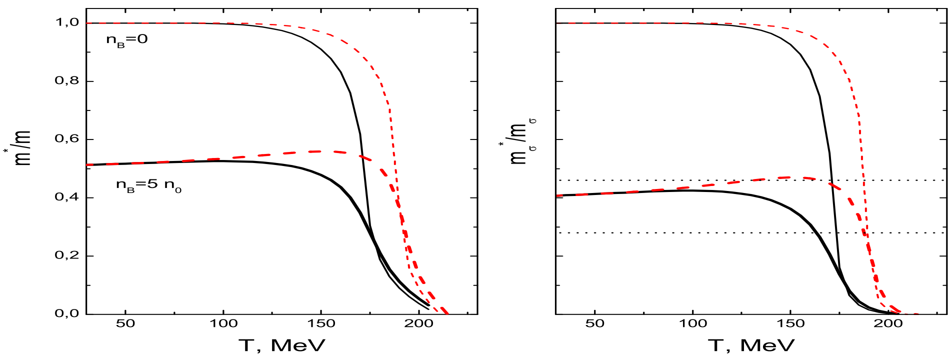

Рис. 1: Temperature dependence of the effective-to-bare mass ratio for nucleon-ω-ρ (left panel) and for σ (right panel) for $n_B = 0$ and $n_B = 5\ n_0$. To guide the eye the horizontal dots show $m^* = 2m_\pi$ and $m^* = m_\pi$ thresholds. Dashed/dotted curves show appropriate results without boson excitations [8].

Parameters of the relativistic mean field model, C_σ, C_ω and the self-interaction potential U, are to be adjusted to reproduce the nuclear matter properties at the saturation for $T = 0$. Values of parameters of the SHMC model [8, 9] are the same as in the KVOR model [4]: the binding energy $e_{\rm bind} = -16$ MeV and nuclear saturation density $n_0 = 0.16\ \ {\rm fm}^{-3}$. Effective nucleon mass and the compressibility coefficient are

$$m^*_N(n_0)/m_N = 0.805, \quad K = 275\ \text{MeV}. \tag{14}$$

One can demonstrate [8, 9] that the SHMC model describes the nucleon optical potential $U_{\rm opt}$ in an optimal way and the pressure at $T = 0$ calculated in the SHMC model well satisfies the experimental constraints coming from the analysis of elliptic flow.

In Fig. 1 we show excited particle mass ratios calculated by eq. (3) for the nucleon and (13) for ω', and ρ' excitations (left panel) and following eq. (12) for σ meson excitation (right) as a function of the temperature. Thin and bold lines are calculated at $n_B = 0$ and $n_B = 5\ n_0$, respectively. The temperature dependence of the effective masses of nucleon and sigma excitations is weak up to $T \sim 150$ MeV. For higher temperatures the effective masses begin to decrease abruptly. For $n_B = 0$ we have $T^{\sigma\pi}(m^*_\sigma = 2m_\pi) \approx T^{\rm chir} \approx 180$ MeV. If one proceeds to the dense matter ($n_B = 5n_0$) the difference between these critical temperatures is about few MeV [9]. Within this narrow temperature interval the second derivative of the effective mass with respect to temperature changes the sign. The effective nucleon-ω'-ρ' and σ' excitation masses reach zero at the same critical temperature $T_{c\sigma}$, about 205 MeV. Effective masses of all excitations exhibit a similar behavior as functions of the temperature and the density, in a line with the mass-scaling hypothesis that we have exploited for the mean fields. Since the coupling scaling functions χ_σ and χ_ω follow the same dropping trend as the mass scaling function Φ, in vicinity of $T_{c\sigma}$ we deal with a gas of almost massless excitations.

If one turns to the case $n_B \simeq 0$ which is close to conditions realized at RHIC, we see the decrease of the hadron masses with increase of the temperature. A similar effect for $n_B \simeq 0$ has been found in [10] as the consequence of the blurring of the baryon and meson vacua.

This sharp decrease of the hadron mass manifests in a strong enhancement of the particle and especially antiparticle yields near the transition temperature, see Fig.2. The

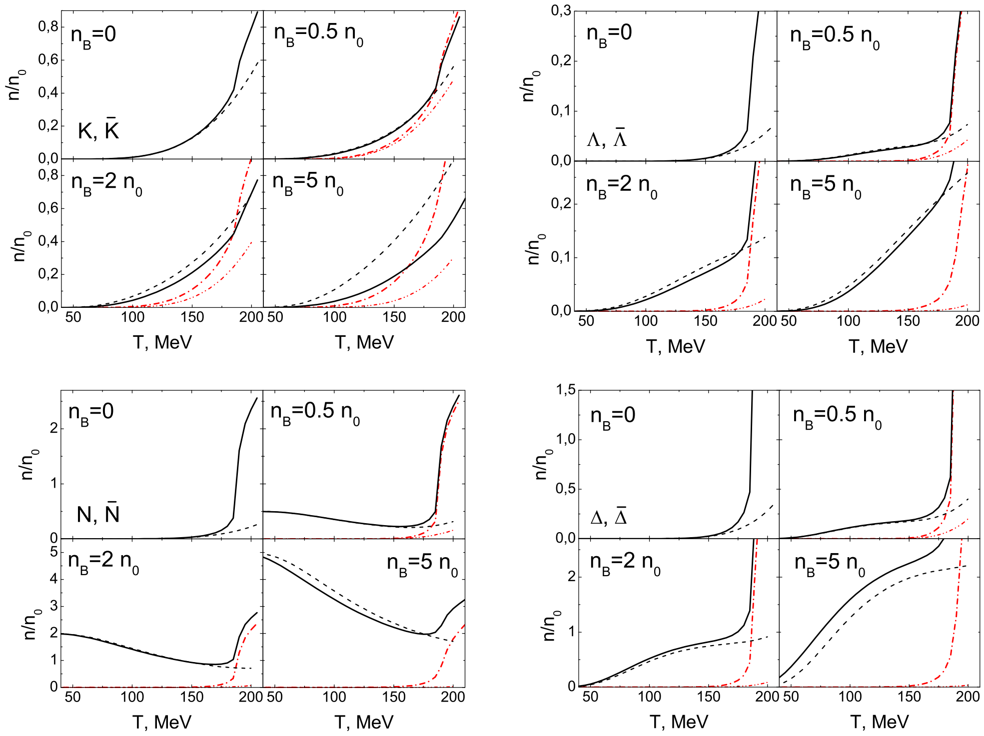

Рис. 2: Particle number density in n_0 units versus temperature at different baryon densities for K, Λ, N, Δ and their antiparticles, $N = Z$. Our model results are plotted by solid line for particles and by dashed-dotted line for antiparticles. Dashed and dashed-double-dotted curves correspond to particles and antiparticles in the ideal gas model.

temperatures of 150-160 MeV are quite reachable at the NICA collider. The appearance of this blurring phase was predicted in [10].

Finally, the modified relativistic mean-field σ-ω-ρ model with scaled hadron masses and couplings was generalized to finite temperatures. Besides nucleon and mean fields the model treats low-lying baryon resonances and their antiparticles, as well as excitations (following the $SU(3)$ concept) including σ'- ω'- ρ'-excitations on the ground of mean fields. The EoS for $T = 0$ satisfies general constraints known from atomic nuclei, neutron stars and those coming from the flow analysis of HIC data. The model supposes the simplest choice of the Brown-Rho scaling when the change of all masses mentioned above follows the same universal law. The developed SHMC model describes the EoS of hot and dense hadronic matter in a broad range of temperatures and baryon densities. The temperature dependence of the EoS at T near $T_{c\sigma}$ reminds the "phase boundary"behavior, being far away from the ideal gas EoS. Our model allows one to predict possible isentropic paths of excited fireball evolution in the (T, μ_B)-plane as discussed in a detail in our papers [8, 9].

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PRECISION SPECTROSCOPY OF LIGHT MOLECULAR ATOMS AND IONS

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In the past few years experimental [1] and theoretical [2] efforts brought the antiprotonic helium spectroscopy to the new level of precision, which allowed to use it in determination of the atomic mass of an electron [3]. The atomic mass deduced from comparison of the antiprotonic helium theory and experiment [3] is now

$$A_r(e) = 0.000\,548\,579\,908\,81(91)[1.7\times10^{-9}].$$

Thus, this exotic atomic system may serve as a new alternative and competitive tool in measuring the fundamental physical constants.

Ro-vibrational transitions for the hydrogen molecular ions H_2^+ and HD^+ were obtained with the relative accuracy of about 3×10^{-10} [4]. That is the ultimate theoretical precision to date. The main uncertainty which sets limits to the present level of precision comes from the yet uncalculated contribution of the one-loop self-energy correction at the $m\alpha^7$ order. This contribution is currently under consideration.

The other major achievements of this year are calculations of the $m\alpha^6(m/M)$ order corrections (including the finite nuclear size corrections) to the hyperfine structure of the ro-vibrational states of the hydrogen molecular ions [5]. That has allowed one to reduce the discrepancy in comparison of transition frequencies with the most precise experiment carried out by Jefferts in 1969 by a factor of six. The Zeeman effect and the HFS intensities of the spectral lines, which are of great importance for the analysis of experiment, were obtained in [6, 7].

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PHOTOIONIZATION AND RECOMBINATION OF A HYDROGEN-LIKE ATOM IN A HOMOGENEOUS MAGNETIC FIELD

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Due to the development of laser and accelerator engines and detection technologies, the investigation of dynamics of light atoms and molecules under the action of fast particles, laser pulses and magnetic fields is considerably interesting for planning and interpretetion of modern experiments. To enhance the laser-stimulated recombination of antihydrogen from cold positron-antiproton plasma in a trap we propose to use a new resonance mechanism involving the quasi-stationary states of the positron which arise from the joint action of the Coulomb field of the antiproton and the strong magnetic field of the trap [1]. The recombination rate is expressed via the cross-section of laser ionization of the atom that has strongly nonmonotonic frequency dependence due to the presence of quasi-stationary states merged into the continuum background [2]. The estimates obtained by using previously calculated ionization cross-section show the possibility to enhance the laser-stimulated recombination of antihydrogen in the cold positron-antiproton plasma conditions of the ATHENA experiment(CERN) by means of the optimal laser frequency and magnetic field choice [3]. To study this possibility, we developed methods of calculations: the photo-ionization cross-sections of a hydrogen-like atom in a homogeneous magnetic field by using the basis of oblate spheroidal functions in the formalism of stationary scattering theory with account taken of axial symmetry of the problem [4–7], the excitation–deexcitation probabilities of Zeeman wave packets of Hydrogen-like atom in the time-dependent electric field [8] and ionization cross-sections of a helium-like atom by electron impact in the framework the wave-packet-evolution method in the paraxial approximation [9]. In this way, resonance transmission and total reflection effects for elastic scattering processes of electrons (positrons) on protons (antiprotons) in a homogenous magnetic field are manifested [10] and possibilities of stabilization and control probability functions of a Zeeman wave packet of Hydrogen-like atom in high-frequency pulse electric field are revealed [11]. These pure quantum effects will be taken into account in calculations of plasma kinetics in magneto-optical traps. Charge-scaling law for angular correlation in double photoionization of the light ions and atoms with two active electrons is analyzed and violation of the Wannier threshold law for the photoionization of a hydrogen negative ion is observed [12].

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