I. PARTICLE AND FIELDS

CONTENTS

INTRODUCTION

The activity in the field of particle theory at BLTP in the years 2010-2011 followed the current trends and covered a wide spectrum of the development and applications of the Standard Model and its extensions. About 200 papers were published and about 50 conference talks were presented.

Special attention was paid to the description of topical experimental problems as well as various issues in the elaboration of the theory methods.

The theoretical developments are the subject of the first three contributions.

In the first contribution by D.I. Kazakov (in collaboration with L.V. Bork and G.I. Vartanov) the general problem of infrared and collinear singularities in quantum field theory was considered in the case of the $N = 4$ supersymmetric Yang-Mills (SYM)theory manifesting integrability and duality with a string theory. Particular emphasis was put on the construction of appropriate observables. It happens that remarkable symmetry properties of the Born amplitude are lost when corrections are considered and this may have a general reason, due to the appearance of dimensionful parameters breaking the conformal invariance.

The contribution by A.V. Kotikov (in collaboration with M. Beccaria, A.V. Belitsky and S. Zieme) deals with the solution of Baxter multiloop equations emerging in a similar context of the N=4 SYM theory. The current development includes a new basis of Wilson polynomials and a new representation for nonpolynomial terms.

The contribution by D.V. Fursaev deals with the enthropic approach to gravity which got a strong boost last year in the work of E. Verlinde. The previously derived (by the author) equation relating the increase of entanglement entropy to the particle dragging from the minimal surface is generalized providing a proof of one of the statements used by E. Verlinde as an axiom.

The following two contributions are dedicated to the development and applications of Analytic Perturbation Theory (APT). The first one, by A.P. Bakulev, S.V. Mikhailov and N.G. Stefanis, is devoted to the progress of fractional APT, generalizing the original analytization procedure of I.L. Solovtsov and D.V. Shirkov to the case of an arbitrary power dependence of coefficient functions on the kinematic observables. As an example, higher loop corrections to the Higgs decay are analyzed with extremely high precision.

The renormalization group equation for QCD coupling is combined with model assumption for the infrared region in the contribution of A.V. Nesterenko (in collaboration with a graduate student Yu. Belyakova) to get the static potential of a quark-antiquark interaction. This potential is fitted to lattice QCD data providing the expression for one-loop $\Lambda_{QCD} \sim 350 MeV$ compatible to other estimates within a potential model.

The QCD effective charge was also considered in the framework of relativistic potential model in the contribution of G. Ganbold. The fits based on the analysis of a meson spectrum using an analog of the ladder Bethe Salpeter equation allowed one to get the estimates of QCD coupling at low momentum transfers.

The following few contributions are devoted to various nonperturbative models being the necessary instrument of hadronic physics.

The recent development and various applications of the relativistic quark model are considered in the contribution by M.A. Ivanov and collaborators. Exotic tetraquark states and double heavy baryons are analyzed.

The model-independent analysis of meson spectrum is performed by Yu.S. Surovtsev and collaborators, concentrating on the higher mass scalar and vector states and their

experimental status.
The contribution by A. Dorokhov is dedicated to the so-called BABAR puzzle for pion transition form factor, corresponding to the difficulties in QCD description of the simplest exclusive hard process. It provides the description of these problems using the nonlocal quark models.
The popular nonperturbative model suggested by Nambu and Iona-Lasinio is discussed in the contributiion of A.B. Arbuzov, E.A. Kuraev, M.K. Volkov and Yu.M. Bystritsky. The important theoretical developments were combined with the new applications to hadronic physics.
The contribution by S.V. Molodtsov describes the applications of the instanton liquid model, in particular, to heavy-ion collisions.
The contributions by A.V. Efremov and collaborators deal with the model description of the Transverse Momentum Dependent (TMD) Parton Distributions and address the relations between various TMDs, which is quite important because of a large number of these new nonperturbative ingredients of QCD factorization.
More formal aspects of TMDs including the gauge invariance, relations to higher twist collinear parton correlators, and causality are discussed in the contribution of I.V. Anikin, I.O. Cherednikov, N.G. Stefanis and O.V. Teryaev. It deals also with the duality between different mechanisms of factorization for hard exclusive processes involving various kinds of Generalized Parton Distributions (GPDs).
The practical applications of GPDs to the processes of hard exclusive electroproduction of various mesons are studied in the contribution of S.V. Goloskokov and P. Kroll.
Finally, the contribution of E.A. Kuraev and collaborators is dedicated to the perturbative calculations of various radiative corrections and cross-sections of hadrons and jets production at the hadron and electron-positron colliders in different energy ranges.

FROM AMPLITUDES TO FORM FACTORS IN THE $\mathcal{N}=4$ SYM THEORY

L.V. Bork[a], D. I. Kazakov[b], G. S. Vartanov[c]

[a]Institute of Theoretical and Experimental Physics, Moscow 117218, Russia
[b]Bogoliubov Laboratory of Theoretical Physics, JINR, Dubna 141980, Russia
[c]Max Planck Inst. fur Gravitationsphysik, Golm, Potsdam D-14476, Germany

1. Much attention in the past few years has been dedicated to the study of the planar limit of the $\mathcal{N}=4$ SYM theory. It is believed that the hidden symmetries responsible for integrability properties of $\mathcal{N}=4$ SYM completely fix the structure of the amplitudes [1].
It was found [2] that these amplitudes revealed »the iterative structure (which was confirmed at the three-loop level)

$$\mathcal{M}_n \equiv \frac{A_n}{A_n^{tree}} = \exp\left[-\frac{1}{8}\sum_{l=1}^{\infty}\lambda^l\left(\frac{\gamma_{cusp}^{(l)}}{(l\epsilon)^2}+\frac{2G_0^{(l)}}{l\epsilon}\right)\sum_{i=1}^{n}\left(\frac{\mu^2}{-s_{i,i+1}}\right)^{l\epsilon}+\frac{1}{4}\sum_{l=1}^{\infty}\lambda^l\gamma_{cusp}^{(l)}F_n^{(1)}(0)+C(g)\right], \tag{1}$$

where $\lambda = g^2N_c/16\pi^2$ is the 't Hooft coupling which stays finite when $N_c\to\infty$, $\gamma_{cusp}(g) = \sum_l \lambda^l\gamma_{cusp}^{(l)}$ is the so-called cusp anomalous dimension and $G_0(g)=\sum_l\lambda^l G_0^{(l)}$ is the second function dependent on the IR regularization.
It is not surprising that the IR divergent parts of the amplitudes factorize and exponentiate. What is less obvious is that it is also true for the finite part. Note, however, that this ansatz is valid for $n=4,5$ but fails starting from $n=6$.
While all the UV divergences in $\mathcal{N}=4$ SYM are absent in scattering amplitudes, the IR ones remain and manifest themselves as poles in ϵ in (1). They are supposed to be canceled in properly defined quantities [3], according to the Kinoshita-Lee-Nauenberg theorem [4]. We concentrated on inclusive cross-sections in the hope that they reveal some factorization properties discovered in the regularized amplitudes and demonstrated the explicit cancelation of the infrared divergencies in properly defined inclusive cross-sections [5]. The same procedure was also applied to $\mathcal{N}=8$ SUGRA [6].
The natural generalization of the on-shell amplitudes are the form factors, i.e. the matrix elements of the form

$$\langle 0|\mathcal{O}|p_1^{\lambda_1}\dots p_n^{\lambda_n}\rangle. \tag{2}$$

where $\mathcal{O}$ is some gauge invariant operator which acts on vacuum and produces some state $|p_1^{\lambda_1}\dots p_n^{\lambda_n}\rangle$ with momenta $p_1\dots p_n$ and helicities $\lambda_1\dots\lambda_n$. One can wonder whether these objects at weak coupling possess similar features as the amplitudes. We study systematically the simplest types of form factors in planar $\mathcal{N}=4$ SYM at weak coupling for half-BPS operators $\mathcal{O}_I^{(n)}$ and the Konishi operator $\mathcal{K}$ [7, 8].
2. Our first aim is to evaluate the NLO correction to the inclusive differential polarized cross section and to trace the cancellation of the IR divergences [5].
We start with the $2\to 2$ MHV scattering amplitude with two incoming positively polarized gluons and two outgoing positively polarized gluons and consider the differential cross-section as a function of the scattering solid angle. At the tree level the cross-section is given by

$$\frac{d\sigma_{(g^+g^+\to g^+g^+)}}{d\Omega_{13}} = \frac{1}{J}\int d\phi_2|\mathcal{M}_4^{(tree)}|^2\mathcal{S}_2, \tag{3}$$

where $\mathcal{S}_2$ is the measurement function and the matrix element is $(s_{ij} = (p_i + p_j)^2)$

$$|\mathcal{M}_4^{(tree)(--++)}|^2 \;=\; g^4 N_c^2(N_c^2-1)\sum_{\sigma\in P_3}\frac{s_{12}^4}{s_{1\sigma(1)}s_{\sigma(1)\sigma(2)}s_{\sigma(2)\sigma(3)}s_{\sigma(3)1}}. \tag{4}$$

Within dimensional regularization(reduction) the cross-section looks like

$$\left(\frac{d\sigma_{2\to 2}}{d\Omega_{13}}\right)_0^{(--++)} = \frac{\lambda^2 N_c^2}{2E^2}\left(\frac{s^4}{t^2u^2}+\frac{s^2}{t^2}+\frac{s^2}{u^2}\right)\left(\frac{\mu^2}{s}\right)^\epsilon = \frac{\lambda^2 N_c^2}{E^2}\left(\frac{\mu^2}{s}\right)^\epsilon \frac{4(3+c^2)}{(1-c^2)^2}, \tag{5}$$

where s,t,u are the Mandelstam variables, E is the total energy in the center of mass frame, $c=\cos\theta_{13}$, μ and ϵ are the parameters of the dimensional regularization(reduction).
The NLO corrections include the virtual and real parts together with the splitting conterterms which appear because of the indistinguishability of the collinear particles in the initial and final states.
Virtual part. The one-loop contribution to the matrix element is already known

$$|\mathcal{M}_4^{(1l)(--++)}|^2 = -g^4N_c^2(N_c^2-1)\lambda\left[\frac{s^4}{s^2t^2}I_4^{(1l)}(s,t)+\frac{s^4}{s^2u^2}I_4^{(1l)}(s,u)-\frac{s^4}{t^2u^2}I_4^{(1-loop)}(-t,u)\right]. \tag{6}$$

where $I_4^{(1l)}(s,t)$ is the scalar box diagram

$$I_4^{(1l)}(s,t) \;=\; -2\frac{\Gamma(1+\epsilon)\Gamma(1-\epsilon)^2}{\Gamma(1-2\epsilon)}[\frac{1}{\epsilon^2}\left((\frac{\mu^2}{s})^\epsilon+(\frac{\mu^2}{-t})^\epsilon\right)+\frac{1}{2}\log^2\left(\frac{s}{-t}\right)+\frac{\pi^2}{2}]+\mathcal{O}(\epsilon).$$

Real emission. One has to consider the amplitude with three outgoing particles. Here we have to define which is the process that we are interested in. There are several possibilities.
1. Three gluons with positive helicities: $g^+g^+\to g^+g^+g^+$. This is the MHV amplitude;
2. Two gluons with positive helicities and one with negative: $g^+g^+\to g^+g^+g^-$. This is the anti-MHV amplitude;
3. One of three final particles is the gluon with positive helicity and the rest is the quark-antiquark pair: $g^+g^+\to\ g^+q^-\overline{q}^+$ or $g^+g^+\to g^+q^+\overline{q}^-$. This is an anti-MHV amplitude;
4. One of three final particles is the gluon with positive helicity and the rest are two scalars: $g^+g^+\to g^+\Lambda\Lambda$. This is the anti-MHV amplitude.
The cross-section of these processes can be written as

$$\frac{d\sigma_{2\to 3}}{d\Omega_{13}} = \frac{1}{J}\int d\phi_3|\mathcal{M}_5^{(tree)}|^2\mathcal{S}_3, \tag{7}$$

where $d\phi_3$ is the 3 particle phase volume and $\mathcal{S}_3$ is the measurement function which constraints the phase space and defines a particular observable. To simplify the integration, in what follows we choose the universal measurement function $\mathcal{S}_3(p_3,p_4,p_5) = \Theta(p_3^0 - \frac{1-\delta}{2}E)\delta^{D-2}(\Omega_{Det}-\Omega_3)$, where we take $\delta=1/3$ in the case of identical particles and $\delta=1$ in other cases. We checked that the IR and collinear divergences cancel in observables for any value of δ.
Splitting. Taking into account the emission of additional soft quanta allows one to cancel the IR divergences (double poles in ϵ) but leaves the single poles originating from collinear ones. Indeed, in the case of massless particles the asymptotic states (both the initial and final ones) are not well defined since massless quantum can split into two parallel ones indistinguishable

from the original one. The emission of a massless gluon leads to a splitting described by *splitting functions* $P_{gg}(z)$. In the case of a gluon in the final state this corresponds to the fragmentation of the gluon into a pair of gluons or a pair of quarks or scalars.
Additional contributions from collinear particles in the initial or final states to inclusive cross-sections have the form, respectively

$$\begin{aligned} d\sigma_{2\to 2}^{InSplit} &= \frac{\alpha}{2\pi}\frac{1}{\epsilon}\left(\frac{\mu^2}{Q_f^2}\right)^{\epsilon}\int_0^1 dz P_{gg}(z)\sum_{i,j=1,2;\ i\neq j} d\sigma_{2\to 2}(zp_i,p_j,p_3,p_4)\mathcal{S}_2^{InSplit}(z), &(8)\\ d\sigma_{2\to 2}^{FnSplit} &= \frac{\alpha}{2\pi}\frac{1}{\epsilon}\left(\frac{\mu^2}{Q_f^2}\right)^{\epsilon} d\sigma_{2\to 2}(p_1,p_2,p_3,p_4)\int_0^1 dz\sum_{l=g,q,\Lambda} P_{gl}(z)\mathcal{S}_2^{FnSplit}(z), &(9)\end{aligned}$$

where the scale Q_f^2 belongs to the definition of the coherent asymptotic state and restricts the value of transverse momenta. The dependence of parton distribution on Q_f^2 is governed by the DGLAP equation.
Final result. In the NLO there are two sets of amplitudes, namely, the MHV and anti-MHV amplitudes which contribute to the observables. The leading order 4-gluon amplitude is both MHV and anti-MHV and we split it into two parts. Then one can construct three types of infrared-safe quantities in the NLO of perturbation theory, namely,

$$Finite = \frac{1}{2}\left(\frac{d\sigma_{2\to 2}}{d\Omega_{13}}\right)_{Virt}^{(--++)} + \left(\frac{d\sigma_{2\to 3}}{d\Omega_{13}}\right)_{Real}^{(--++j)} + \left(\frac{d\sigma_{2\to 3}}{d\Omega_{13}}\right)_{InSplit}^{(--++j)} + \left(\frac{d\sigma_{2\to 3}}{d\Omega_{13}}\right)_{FnSplit}^{(--++j)}, \quad (10)$$

where $j=+,-$ and $q\bar{q}+\Lambda\Lambda$ corresponds to the MHV, anti-MHV and matter amplitudes, respectively. In each case *all IR divergencies cancel* for arbitrary δ and only the finite part remains.
Defining now the physical condition for the observation we get several infrared-safe inclusive cross-sections. Relative simplicity of the virtual contribution does not hold for the real part. While the singular terms are simple enough and cancel completely, the finite parts are usually cumbersome and contain polylogarithms. The only expression where they cancel corresponds to the $\delta=1$ case for the anti-MHV cross-section (10). Choosing the factorization scale to be $Q_f=E$ we get:

$$\begin{aligned}\left(\frac{d\sigma}{d\Omega_{13}}\right)_{anti-MHV} &= \frac{4\lambda^2 N_c^2}{E^2}\left\{\frac{3+c^2}{(1-c^2)^2}\right.\\ &-\frac{\lambda}{4\pi}\left[2\frac{(c^4+2c^3+4c^2+6c+19)\log^2(\frac{1-c}{2})}{(1-c)^2(1+c)^4} - 2\frac{(11c^3-31c^2-47c-133)\log(\frac{1-c}{2})}{3(1+c)^3(1-c)^2} + (c\leftrightarrow -c)\right.\\ &\left.\left.-8\frac{(c^2+1)\log(\frac{1+c}{2})\log(\frac{1-c}{2})}{(1-c^2)^2} + \frac{6\pi^2(3c^2+13)-5(61c^2+99)}{9(1-c^2)^2}\right]\right\}. \quad (11)\end{aligned}$$

One can see that even this expression does not repeat the form of the Born amplitude and does not have any simple structure. The reason for this might be that constructing the infrared finite observable we mix the MHV and non-MHV amplitudes, thus loosing the fine properties of the former. The other reason might be that the MHV amplitudes themselves for the number of legs exceeding 5 do not follow the exponentiation pattern for the finite parts.

3. Now we turn to the calculation of the form factors [8]. It is convenient to use the $\mathcal{N}=1$ formulation of $\mathcal{N}=4$ SYM and perform an explicit computation in terms of the $\mathcal{N}=1$ superfields in momentum space. We choose the following set of $\mathcal{N}=1$ local operators:

$$\begin{aligned}
\mathcal{C}_{IJ} &= \mathrm{Tr}(\Phi_I\Phi_J), I\neq J\\
\mathcal{V}_I^J &= \mathrm{Tr}(e^{-gV}\bar{\Phi}^J e^{gV}\Phi_I), I\neq J\\
\mathcal{O}_I^{(n)} &= \mathrm{Tr}\left(\Phi_I^n\right),\\
\mathcal{K} &= \sum_I \mathrm{Tr}(e^{-gV}\bar{\Phi}^I e^{gV}\Phi_I),
\end{aligned} \tag{12}$$

where Φ_I are the chiral $\mathcal{N}=1$ superfields, and V is the $\mathcal{N}=1$ real vector superfield. The operators $\mathcal{C}_{IJ}$, $\mathcal{O}_I^{(n)}$ are chiral and $\mathcal{V}_I^J$, $\mathcal{K}$ are non-chiral from the $\mathcal{N}=1$ supersymmetric point of view.
We use the following notation for the form factor of the corresponding operator:

$$\mathcal{F}(p_1\dots p_n) = \langle p_1\dots p_n|\mathcal{O}(q)|0\rangle. \tag{13}$$

and consider the ratio $\mathcal{M}=\mathcal{F}/\mathcal{F}_{tree}$.
Form factors with $\Delta_0=2$. For the operators $\mathcal{C}_{IJ}$, $\mathcal{V}_I^J$ and $\mathcal{K}$ the form factors have a similar nature and the following form:

$$\log(\mathcal{M}) = \frac{1}{2}\sum_{i=1}^{2}\left(\hat{M}(s_{i,i+1}/\mu^2)\right) + O(\epsilon), \tag{14}$$

where

$$\hat{M}(s_{i,i+1}/\mu^2) = -\frac{1}{2}\sum_l\left(\frac{\lambda}{16\pi^2}\right)^l\left(\frac{\gamma_{cusp}^{(l)}}{(l\epsilon)^2}+\frac{G^{(l)}}{l\epsilon}+C^{(l)}\right)\left(\frac{s_{i,i+1}}{\mu^2}\right)^{l\epsilon}. \tag{15}$$

Here $\gamma_{cusp}^{(l)}$ are the coefficients of perturbative expansion of the cusp anomalous dimension, which is a universal quantity that governs the IR behavior of gauge theory amplitudes and the UV behavior of the Wilson loops, and some local gauge invariant operators. The quantities $G^{(l)}$ and $C^{(l)}$ are regularization and scheme dependent. We got

$$\log(\mathcal{M}) = \lambda\left(\frac{s_{12}}{\mu^2}\right)^{-\epsilon}\left(\frac{-2}{\epsilon^2}+\zeta_2\right) + \lambda^2\left(\frac{s_{12}}{\mu^2}\right)^{-2\epsilon}\left(\frac{\zeta_2}{\epsilon^2}+\frac{\zeta_3}{\epsilon}\right) + O(a^3) \tag{16}$$

where ζ_n are the Riemannian zeta functions. This gives the first two terms of perturbative expansion for the cusp and the collinear anomalous dimensions and the finite terms

$$\gamma_{cusp}^{(1)}=4,\ \gamma_{cusp}^{(2)}=-8\zeta_2,\quad G^{(1)}=0,\ G^{(2)}=-\zeta_3,\quad C^{(1)}=-\zeta_2,\ C^{(2)}=0. \tag{17}$$

Form factors with $\Delta_0=n$ for $n\geq 3$. Consider the chiral half-BPS operators $\mathcal{O}_I^{(n)}$ introduced earlier. At the second order of perturbation theory we obtain for $\log(\mathcal{M})$:

$$\log(\mathcal{M}) = \sum_{i=1}^{3} a\left(\frac{s_{ii+1}}{\mu^2}\right)^{-\epsilon}\left(-\frac{1}{\epsilon^2}+\frac{\zeta_2}{2}\right) + \sum_{i=1}^{3} a^2\left(\frac{s_{ii+1}}{\mu^2}\right)^{-2\epsilon}\left(\frac{\zeta_2}{2\epsilon^2}+\frac{7\zeta_3}{2\epsilon}\right) + \text{fin.part.} \tag{18}$$

This gives

$$\gamma_{cusp}^{(1)}=4, \gamma_{cusp}^{(2)}=-8\zeta_2,\quad G^{(1)}=0, G^{(2)}=-7\zeta_3. \tag{19}$$

Notice that the values of the cusp anomalous dimension $\gamma^{(l)}$ are universal and coincide with (17), while those of the collinear anomalous dimension depend on the form factor at hand. We would like to emphasize the highly nontrivial cancelations between the polylogarithms that occurred for $\log(\mathcal{M})$ for the whole set of scalar integrals. We see that the IR factorization property holds for the form factors as for the amplitudes.
As for the finite part at one loop it is trivial $F^{(1)} = 0$, and the two loop expression $F^{(2)}$ is a complicated function containing logarithms, polylogarithms and generalized Goncharov polylogarithms of several variables [8]. However, the result is still much simpler than in the non-supersymmetric case and the maximal transcendentality principle still holds.
4. We present here our recent achievements in understanding the IR nature of $\mathcal{N} = 4$ SYM. The MHV amplitudes having a number of remarkable mathematical properties do not have much sense being infrared divergent. And though the understanding of their behaviour is still an interesting mathematical issue, this cannot be considered as a final goal. At the same time, the physical quantities seem to be more complicated and we are far from "exact" solution. The form factors represent the next set of objects of interest. They all share the property of factorization of the IR divergences being governed by the universal anomalous dimension, but the finite parts are far from being simple.
The $\mathcal{N} = 4$ SYM theory is the first example of conformal quantum field theory in 4 dimensions. There are many indications that it can be integrable in some sense. However, to get integrability, one should have an infinite number of relations, i.e. an infinite number of conservation laws. It is quite possible that the Yangian symmetry [1] will provide us with them.

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ANALYTIC SOLUTION OF THE MULTILOOP BAXTER EQUATION

M. Beccaria[a], A.V. Belitsky[b], A.V. Kotikov[c] and S. Zieme[d]

[a]Physics Department, Salento University and INFN, 73100 Lecce, Italy
[b]Department of Physics, Arizona State University Tempe, AZ 85287-1504, USA
[c]Bogoliubov Laboratory of Theoretical Physics, JINR, Dubna 141980, Russia
[d]Max-Planck-Institut fur Gravitationsphysik Albert-Einstein-Institut Am Muhlenberg 1, D-14476 Potsdam, Germany and Faculty of Science, University of Herat, Afghanistan

The success of gauge theories in describing accurately the laws of nature is based on the availability of computational techniques, see e.g., Ref. [1], which allow for a systematic improvement of approximations involved. Perturbative expansions in the gauge coupling constant $g_{\rm YM}$ are conventionally deduced from Feynman diagrams. However, due to uncontrollable proliferation of the latter at higher orders in $g_{\rm YM}$, the rules quickly become unmanageable, making direct computations already at four-loop order highly nontrivial and require massive computer manipulations. On top of this, individual Feynman diagrams obscure underlying properties of the theory and reveal simple results enjoying sometimes enhanced symmetries only in their sum. One was therefore compelled to search for an alternative approach which presented itself recently.

On the one hand, some time ago it was established that at weak coupling one-loop spectra of anomalous dimensions of maximal-helicity gauge-invariant operators in QCD coincide with energy spectra of a one-dimensional non-compact Heisenberg magnet [2]. The latter can be diagonalized by means of the traditional Bethe ansatz formalism of integrable systems and yields anomalous dimensions of the corresponding four-dimensional gauge theory. These simplifications are echoed by higher loop contributions, especially in supersymmetric gauge theories. It was found in Refs. [3, 4] that all single-trace operators in planar, maximally supersymmetric gauge theory

$$\mathcal{O} = \mathrm{tr}\,\left(X(D_+^2 X)YZX\lambda XF_{+\perp}(D_+Y)\bar{\lambda}\ldots\right) \tag{1}$$

can be described by a long-range integrable spin-chain model with elementary excitations identified with the particle fields Y, Z, λ etc. of the gauge theory and/or covariant derivatives D_μ acting on them propagating on the vacuum state $|0\rangle = \mathrm{tr}\,\left(X^L\right)$.

On the other hand, the AdS/CFT correspondence [5] conjectures that the strongly coupled $\mathcal{N} = 4$ Yang-Mills theory is dual to a free type IIB super-string theory on an $\mathrm{AdS}_5\times\mathrm{S}^5$ background. The latter was found to be classically integrable as well [6]. Using this conjecture as a virtue led to a suggestion of an integrable structure which interpolates between weak and strong coupling regimes. Though the underlying spin chain model is not known, a set of Bethe ansatz equations is nevertheless available [4, 7], which has passed a number of non-trivial tests at weak coupling, see e.g. [7] and [8], as well as at strong coupling by positive comparison with perturbative string theory, see e.g. [9].

These findings suggest using the putative integrable structure as an alternative to the conventional Feynman diagrams technique for multiloop calculations of anomalous dimensions. In paper [10], we developed a practitioner's formalism building up on earlier considerations based on the all-order Baxter equation [11] for finding the spectrum of twist-two Wilson operators

$$\mathcal{O} = \mathrm{tr}\,\left(XD_+^M X\right)\,. \tag{2}$$

These arise in all gauge theories albeit with a different field content, the scalar X being specific to supersymmetric cousins of QCD. Their anomalous dimensions have been obtained diagrammatically to a considerably high-order (see, for example, the recent review [12]).
The Baxter equation is advantageous over the Bethe ansatz formalism if one is interested in a systematic analytical framework. However, they both enter on equal footings for numerical studies, and Bethe equations were used in the past together with the principle of maximal transcendentality [13] to perform phenomenal computations [14].
Our consideration [10] is a generalization of the study in Ref. [15] which was based on a deformation of the solution to the one-loop Baxter equation. What will differ in the current work is that we will introduce a new basis of functions used in the construction of next-to-leading order solutions, the so-called Wilson polynomials. For comparison, we also present the basis of continuous Hahn polynomials used in [15]. Furthermore, we obtain a new form for non-polynomial contributions which is free from multiple sums involving Stirling numbers. The latter property is essential for obtaining analytical results for anomalous dimensions in terms of nested harmonic sums.

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‘THERMODYNAMICS’ OF MINIMAL SURFACES AND ENTROPIC ORIGIN OF GRAVITY

D.V. Fursaev

International University “Dubna” and Bogoliubov Laboratory of Theoretical Physics, JINR, Dubna 141980, Russia

The fact that gravity is an emergent phenomenon dates back to ideas of the last century [1]. A renewal of the interest in this point of view in the last years has been motivated by attempts to find statistical explanation of the Bekenstein-Hawking entropy, see e.g. [2]. A possible source of the entropy is quantum correlations of underlying microscopical degrees of freedom across the black hole horizon.
By taking the black hole case as a guide a number of arguments have been presented in [3] that the entanglement entropy of fundamental degrees of freedom lying in a constant time slice and spatially separated by a surface $\mathcal{B}$, is

$$S(\mathcal{B}) = \frac{\mathcal{A}(\mathcal{B})}{4G} \quad . \tag{1}$$

Here G is the Newton coupling and $\mathcal{A}$ is the area of $\mathcal{B}$. Thus, (1) has the Bekenstein-Hawking form. Equation (1) holds in the semiclassical approximation if the low-energy limit of the fundamental theory is the Einstein gravity.
For realistic condensed matter systems the entanglement entropy associated with spatial separation of the system is a nontrivial function of microscopic parameters. Its calculation is technically involved and model dependent. The remarkable consequence of (1) is that the entanglement entropy in quantum gravity may not depend on a microscopic content of the theory, it is determined solely in terms of geometrical characteristics of the surface and low-energy gravity couplings.
Another feature established in [3] is related to the shape of the separating surface. Since $S(\mathcal{B})$ includes contributions of all fundamental degrees of freedom, quantum fluctuations of the geometry should be taken into account in a consistent way. For static space-times this requires that $\mathcal{B}$ is a minimal surface, i.e. a surface with a least area. A relevant physical example of a minimal surface is the intersection of a constant time slice and the horizon of a stationary black hole. Thus, the Bekenstein-Hawking entropy can be considered as a particular case of the entanglement entropy (1).
The fact that $S(\mathcal{B})$ is a macroscopic quantity that obeys certain dynamical laws points to similarity with a thermodynamical entropy. A natural question is whether the entanglement on the fundamental level admits a form of thermodynamical laws.
The first step in this direction has been made in [3]. A calculation made there in the weak field approximation shows that a shift by a distance l of a test particle with a mass m out of the minimal surface results in the entropy change

$$\delta S(\mathcal{B}) = -\pi m \ l \quad . \tag{2}$$

A work needed to drag the particle by the background gravitational field is also proportional to l. One can relate the entropy change (2) and the work, the relation being an analog of the first law of thermodynamics. This also yields a local temperature on the surface (proportional to the product of the acceleration of a static observer near the surface and the normal vector to $\mathcal{B}$).

An intriguing hypothesis was suggested by E. Verlinde [4] in 2010 that gradients of the entropy of microscopic degrees of freedom in an underlying quantum gravity theory determine gradients of the gravitational field. The classical force of gravity can be interpreted as an entropic force, thus losing its fundamental nature. The hypothesis is based on a number of assumptions for so called 'holographic screens' which store information about fundamental microstates ('bits') in such a way that a related entropy is proportional to the area of the screen. Interestingly, formula (2) for the entropy change is used by Verlinde as one of the postulates.
That is why deformations of minimal surfaces lying in constant time slices were studied in [5] for general static solutions to the Einstein equations. It was shown that formula (2) was remarkably universal and held for such a general set up. This result provides: 1) a strong support to the hypothesis [4] that gravity has an entropic origin, the minimal surfaces being a sort of holographic screens; 2) reduces by one the number of postulates of [4]; 3) suggests definite physical interpretation for the entropy on the screen as the entropy of entanglement across the screen of quantum gravity states.

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FRACTIONAL ANALYTIC PERTURBATION THEORY AND ITS APPLICATIONS

A.P. Bakulev, S.V. Mikhailov, N.G. Stefanis

International University "Dubna" and Bogoliubov Laboratory of Theoretical Physics, JINR, Dubna 141980, Russia

In a series of papers [1, 2], prepared partially in collaboration with Prof. A. I. Karanikas from the *Department of Physics, University of Athens (Greece)*, we developed in fully worked out detail a new generalization of the QCD Analytic Perturbation Theory (APT) of Shirkov and Solovtsov [3] – a version of the QCD perturbation theory without the Landau pole singularity. This extension aims to improve the computation of higher-order corrections in inclusive and exclusive processes by defining analytic images of the running coupling powers ($\mathfrak{A}_\nu^{\text{glob}}$ and $\mathcal{A}_\nu^{\text{glob}}$ in Minkowskian and Euclidean regions, respectively) for any *fractional* (real) power ν of the running coupling α_s^ν — by this reason it is called the Fractional Analytic Perturbation Theory (FAPT). As a result, in full analogy with the ATP case, a power series of the standard QCD PT $\sum_n d_n\,\alpha_s^{n+\nu}$ transforms with FAPT into non-power series $\sum_n d_n\,\mathfrak{A}_{n+\nu}^{\text{glob}}$ and $\sum_n d_n\,\mathcal{A}_{n+\nu}^{\text{glob}}$, respectively.
In our recent papers [4, 5], we used the technique of non-power series resummation in FAPT to analyze uncertainties of QCD predictions for the Higgs boson decay width, $\Gamma_{H^0\to\bar{b}b}(M_{\text{H}})$. To this end, we constructed a factorially growing model

$$d_n^{\text{H}} = d_1\,c^{n-1}\,\frac{\Gamma(n+1)+\beta\,\Gamma(n)}{1+\beta},$$

of perturbative coefficients d_n (known for $n=1,2,3,4$) with fixed parameters $c=2.4$ and $\beta=-0.52$, based on the generating function

$$P_{\text{H}}(t) = \frac{\beta+t/c}{c\,(\beta+1)}\,e^{-t/c}, \quad d_n^{\text{H}} = d_1\int_0^1 P_{\text{H}}(t)\,t^{n-1}\,dt. \tag{1}$$

Predictions of this model appear to be very close to the predictions obtained by using the Principe of Minimal Sensitivity [6], both for the known coefficients d_n with $n=1,2,3,4$ and for the unknown d_5. We analyzed the truncation errors

$$\Delta_N[L] = 1-\frac{\Gamma^{\text{FAPT}}_{H\to b\bar{b}}[L;N]}{\Gamma^{\text{FAPT}}_{H\to b\bar{b}}[L]}\,, \tag{2}$$

due to the truncation of the FAPT series for the decay width $\Gamma_{H^0\to\bar{b}b}[L]$ at the order N,

$$\Gamma^{\text{FAPT}}_{H\to b\bar{b}}[L;N] = \Gamma_0^b(\hat{m}^2)\,\left\{\mathfrak{A}_{\nu_0}^{\text{glob}}[L]+d_1\,\sum_{n=1}^{N}\frac{\tilde{d}_n}{\pi^n}\,\mathfrak{A}_{n+\nu_0}^{\text{glob}}[L]\right\}\,, \tag{3}$$

in the region $L=\ln(M_{\text{H}}^2/\Lambda_{\text{QCD}}^2)\in[11.7,13.6]$, corresponding to the Higgs boson mass region $80-180$ GeV at the QCD scale value $\Lambda_{\text{QCD}}^{N_f=3}=201$ MeV, that means $\mathfrak{A}_1^{\text{glob}}(m_Z^2)=0.122$. Our results show that already $\Gamma^{\text{FAPT}}_{H\to b\bar{b}}[L;2]$ gives an accuracy better than 2.5%, while $\Gamma^{\text{FAPT}}_{H\to b\bar{b}}[L;3]$ provides an accuracy of the order of 1%. That means that in order to predict the Higgs boson decay width with an accuracy of 1% in the region $m_{\text{H}}=60-180$ GeV, it is sufficient

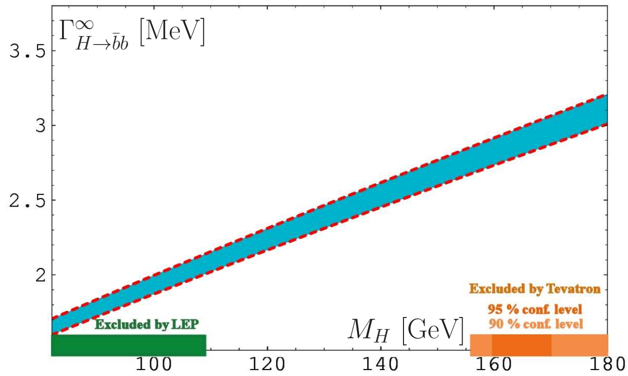

Fig. 1: Higgs-boson decay width $\Gamma^{\infty}_{H\to b\bar{b}}$ as a function of Higgs-boson mass M_H in resummed FAPT with variations $\hat{m}_b = 8.22 \pm 0.13$ GeV (in accord with Penin–Steinhauser estimates $\overline{m}_b(\overline{m}_b^2) = 4.35 \pm 0.07$ GeV [7]) at the two-loop running of the effective coupling. The window of the Higgs-boson-mass values which is still accessible to experiment is explicitly indicated.

to take into account terms with coefficients d_0, d_1, d_2 and d_3, whereas addition of the term with d_4 results in an accuracy benefit of 0.5% — that is not so important if one compares it with 2% uncertainty of the renormgroup-invariant mass.

In addition, we analyzed the sensitivity of our FAPT resummation result to the details of the model for the generating function $P(t)$: we used two deformed models $P_{\pm}(t)$ for which the coefficients $\tilde{d}_n$ are enhanced ($P_+(t)$) or reduced ($P_-(t)$) from 5 to 15%. We showed that this type of uncertainties provides not more than 0.6%. At the same time, the RG-invariant mass uncertainty appears to be $\simeq 2\%$, so that the resulting uncertainty of our predictions is of the order of 3%.

To conclude: FAPT allows one to estimate the relative importance of the higher-order perturbative corrections. Having the RG-invariant mass uncertainty of the order of 1% it is too early to take into account the four-loop correction.

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A NONPERTURBATIVE MODEL FOR THE STRONG RUNNING COUPLING WITHIN THE POTENTIAL APPROACH

Yu.O. Belyakova[a], A.V. Nesterenko[b]

[a]Department of Theoretical Physics, Moscow State University, Moscow 119991, Russia
[b]Bogoliubov Laboratory of Theoretical Physics, JINR, Dubna 141980, Russia

Theoretical description of hadron dynamics at large distances remains a crucial challenge of elementary particle physics for a long time. The asymptotic freedom of Quantum Chromodynamics (QCD) allows one to apply perturbation theory to study some "short–range" processes, for example, the high–energy hadronic reactions. However, the study of many phenomena related to the "long–range" dynamics (such as confinement of quarks, structure of the QCD vacuum, etc.) can be performed only within nonperturbative methods. In what follows, we shall employ the so–called potential approach [1] that involves the construction of the QCD invariant charge which satisfies certain nonperturbative requirements.

In accordance with the basic idea of the potential approach, we shall construct the strong running coupling $\alpha(Q^2)$ that coincides with perturbative QCD invariant charge in the ultraviolet domain ($Q^2 \to \infty$) and meets the requirement of the infrared enhancement at low energies ($Q^2 \to 0_+$). In terms of the renormalization group (RG) β–function

$$\frac{d\,\ln a(\mu^2)}{d\,\ln \mu^2} = \beta(a), \qquad \beta(a) \simeq \beta^{(\ell)}_{\text{pert}}(a) = -\sum_{n=0}^{\ell-1} B_n \left[a^{(\ell)}(\mu^2)\right]^{n+1}, \quad a \to 0_+ \tag{1}$$

the afore-mentioned conditions can be equivalently rewritten as

$$\beta(a) \simeq \beta_{\text{pert}}(a), \quad a \to 0_+, \qquad \beta(a) \simeq -1, \quad a \to \infty. \tag{2}$$

Here $a(Q^2) \equiv \alpha(Q^2)\beta_0/(4\pi)$ denotes the so–called "couplant", $\ell = 1, 2, 3, ...$ stands for the loop level, $B_n = \beta_n/\beta_0^{n+1}$ is the ratio of the QCD β–function perturbative expansion coefficients, $\beta_0 = 11 - 2\,n_f/3$, and n_f denotes the number of active quarks. One of the possible expressions for the β–function that satisfies conditions (2) reads [2]

$$\beta^{(\ell)}(a) = \beta^{(\ell)}_{\text{pert}}(a)\,\frac{1 - \exp(-2/a)(1 - \ell^2/B_{\ell-1})}{1 + \ell^2 a^\ell}, \tag{3}$$

where $\beta^{(\ell)}_{\text{pert}}(a)$ is the ℓ–loop perturbative QCD β–function (1).

At the one–loop level ($\ell = 1$) RG equation for the QCD invariant charge $\alpha(Q^2)$ corresponding to the β–function (3) can be solved explicitly:

$$\alpha^{(1)}(Q^2) = \frac{4\pi}{\beta_0}\,\frac{1}{W_0(Q^2/\Lambda^2)}, \qquad \Lambda^2 = Q_0^2\,\frac{\beta_0}{4\pi}\,\alpha^{(1)}(Q_0^2)\exp\left[-\frac{4\pi}{\beta_0}\,\frac{1}{\alpha^{(1)}(Q_0^2)}\right], \tag{4}$$

where $W_0(x)$ denotes the principal branch of the Lambert W–function. At the higher loop levels ($\ell > 1$) the RG equation with β–function (3) can only be integrated numerically. The plots of the couplants $a^{(\ell)}(Q^2)$ and $a^{(\ell)}_{\text{pert}}(Q^2)$ at the ℓ–loop level ($\ell = 1, ..., 4$) are presented in Figure 1. As one can infer from this figure, the QCD invariant charge $\alpha(Q^2)$ corresponding to

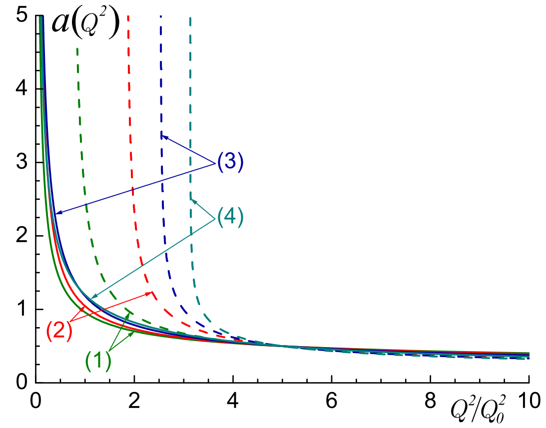

Fig. 1: The ℓ–loop couplant $a^{(\ell)}(Q^2)$ corresponding to β–function (3) (solid curves) and the perturbative couplant $a^{(\ell)}_{\text{pert}}(Q^2)$ (dashed curves). The functions are computed for $n_f = 3$ active quarks and normalized to the value $a(Q^2) = 1/2$ at $Q^2 = 5\,Q_0^2$. The numerical labels indicate the loop level.

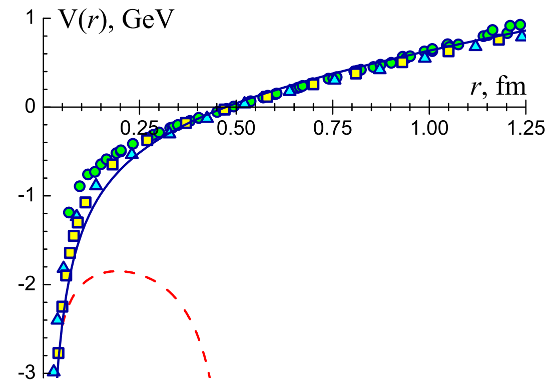

Fig. 2: The quark–antiquark potential $V(r)$ (Eq. (8), solid curve), lattice data (Ref. [3], "○"), Cornell potential (Ref. [4], "□"), Richardson's potential (Ref. [5], "△"), and the perturbative result (Eq. (9), dashed curve). Values of parameters: $\Lambda = 375\,\text{MeV}$, $V_0 = 315\,\text{MeV}$, $n_f = 3$.

β–function (3) possesses elevated (with respect to perturbative results) higher loop correction stability in the intermediate energy range. It is worthwhile to note also that the proposed strong running coupling $\alpha(Q^2)$ is free of low–energy unphysical singularities.

In the framework of the potential approach the static potential of quark–antiquark interaction $V(r)$ is related to the strong running coupling $\alpha(Q^2)$ which contains no unphysical singularities and satisfies the afore-mentioned requirements by

$$V(r) = -\frac{16\pi}{3}\int_0^\infty \frac{\alpha(\boldsymbol{Q}^2)}{\boldsymbol{Q}^2}\,\frac{\exp(i\boldsymbol{Q}\boldsymbol{r})}{(2\pi)^3}\,d\boldsymbol{Q}, \tag{5}$$

see, e.g., reviews [1] and references therein for details. In what follows, for the construction of the static quark–antiquark potential we use the invariant charge (4). After integration over angular variables, Eq. (5) in this case takes the form

$$V(r) = \frac{8\pi}{3\beta_0}\,\Lambda\Big[U_1(R) + U_2(R)\Big], \qquad U_i(R) = -\frac{1}{R}\,\frac{4}{\pi}\int_0^\infty a_i\!\left(\frac{x^2}{R^2}\right)\frac{\sin x}{x}\,dx, \tag{6}$$

where $a_1(z) = 1/z$, $a_2(z) = 1/W_0(z) - 1/z$, $z = Q^2/\Lambda^2$, $R = \Lambda r$, and $x = Qr$. The function $U_1(R)$ diverges and requires regularization, whereas $U_2(R)$ is regular and can be computed numerically.

To regularize the function $U_1(R)$, we shall employ the method similar to that used in Ref. [6]. Specifically, in terms of an auxiliary function

$$I(t) = \int_0^\infty x^t \sin x\,dx = \sqrt{\pi}\,2^t\,\frac{\Gamma(1+t/2)}{\Gamma\big((1-t)/2\big)} \tag{7}$$

the singular part of the potential (6) reads $U_1(R) = -4\,R\,I(-3)/\pi = R$. The right–hand side of Eq. (7) represents the analytic continuation of the function $I(t)$ to the entire complex

t–plane (except for the points $t = -2N$, $N = 1, 2, 3, \dots$) that plays the role of regularization of the function $U_1(R)$, see Refs. [6, 7, 2] for details.
Thus, the static quark–antiquark potential (6) takes the following form:

$$V(r) = V_0 + \frac{8\pi}{3\beta_0}\,\Lambda\left[R - \frac{1}{R}\,\frac{4}{\pi}\int_0^\infty a_2\!\left(\frac{x^2}{R^2}\right)\frac{\sin x}{x}\,dx\right], \qquad R = \Lambda r, \tag{8}$$

where V_0 is an additive self–energy constant. At small distances this potential possesses the standard behavior determined by the asymptotic freedom, whereas at large distances potential (8) proves to be linearly rising implying the confinement of quarks:

$$V(r) \simeq V_{\text{pert}}(r) = \frac{8\pi}{3\beta_0}\,\frac{\Lambda}{R\,\ln R}, \quad r \to 0, \qquad V(r) \simeq V_{\text{conf}}(r) = \frac{8\pi}{3\beta_0}\,\Lambda\,R, \quad r \to \infty. \tag{9}$$

The potential (8) satisfies also the concavity condition $dV(r)/dr > 0$, $d^2V(r)/dr^2 \le 0$, which is a general property of the gauge theories (see Ref. [8] for the details).
Figure 2 presents the quark–antiquark potential (8), lattice simulation data [3], Cornell's potential [4], and Richardson's potential [5]. Equation (8) has been fitted to the lattice data [3] by making use of the least square method, Λ and V_0 being the varied parameters. The estimation of the scale parameter in the course of this comparison gives $\Lambda = (375 \pm 40)$ MeV (this value corresponds to the one–loop level with $n_f = 3$ active quarks), which agrees with previous estimations of this parameter within the potential approach. As one can infer from Figure 2, in the region $r \lesssim 0.05$ fm the derived potential (8) coincides with the perturbative result (9). At the same time, in the nonperturbative physically–relevant range $0.3\,\text{fm} \lesssim r \lesssim 1.2\,\text{fm}$ the obtained potential (8) reproduces the lattice data [3] fairly well. Additionally, $V(r)$ (8) is in good agreement with both the Cornell [4] and Richardson [5] potentials.

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QCD EFFECTIVE CHARGE AND MESON SPECTRUM

G. Ganbold

The behaviour of the QCD effective charge α_s has been investigated by exploiting the conventional meson spectrum within a relativistic quantum-field model based on analytic (or, infrared) confinement. The spectra of two-quark bound states were defined by using a master equation similar to the ladder Bethe-Salpeter equation. We derive meson mass formula and adjust the model parameters (quark constituent masses and the confinement scale Λ) by fitting heavy meson masses [1]. Having adjusted model parameters, we estimate $\alpha_s(M)$ in the low-energy domain by exploiting light meson masses. We found new, independent and specific infrared-finite behavior of QCD coupling below energy scale 1 GeV [2]. Particularly, an infrared-fixed point is extracted at $\alpha_s(0) \simeq 0.757$ for $\Lambda = 345$ MeV. As an application, we estimate masses of intermediate and heavy mesons and obtain results in reasonable agreement with recent experimental data. We demonstrate that global properties of the low-energy phenomena such as QCD running coupling and conventional meson spectrum may be explained reasonably in the framework of a simple relativistic quantum-field model if one guesses a correct symmetry structure of the quark-gluon interaction in the confinement region and uses simple forms of propagators in the hadronization regime.

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DEVELOPMENT AND APPLICATIONS OF RELATIVISTIC CONSTITUENT QUARK MODEL

M.A. Ivanov

The relativistic constituent quark model developed in Dubna was refined [1] to include the confinement of quarks. It is done, first, by introducing the scale integration in the space of alpha-parameters, and, second, by cutting this scale integration on the upper limit which corresponds to an infrared cutoff. In this manner, one removes all possible thresholds presented in the initial quark diagram. The cutoff parameter is taken to be the same for all physical processes.

The consequences of treating the X(3872) meson as a tetraquark bound state was explored [2, 3] within a relativistic constituent quark model with infrared confinement. The decay widths of the observed channels $X \to J/\psi + 2\pi(3\pi)$ and $X \to \bar{D}^0 + D^0 + \pi^0$ were calculated. For reasonable values of the size parameter of the X(3872) the consistency of the theoretical results with the available experimental data were found.

The flavor-conserving radiative decays of double heavy baryons were studied by using a manifestly Lorentz covariant constituent three-quark model [4, 5]. Decay rates were calculated and compared to each other in the full theory, keeping masses finite, and also in the heavy quark limit.

The strong couplings $g_{D^*D\pi}$ and $g_{B^*B\pi}$ were computed [6] by using a framework in which all elements are constrained by Dyson-Schwinger equation studies of QCD and, therefore, incorporate a consistent, direct and simultaneous description of light- and heavy-quarks and the states they may constitute. These couplings were found to be $g_{D^*D\pi} = 15.9^{+2.1}_{-1.0}$ and $g_{B^*B\pi} = 30.0^{+3.2}_{-1.4}$. A comparison between them indicates that when the c-quark is a system's heaviest constituent, $\Lambda_{\rm QCD}/m_c$-corrections are not under good control.

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MODEL INDEPENDENT ANALYSIS OF MESON SPECTRUM

Yu.S. Surovtsev

In a model-independent approach, based on analyticity and unitarity, combined 3-channel analyses of data on isoscalar S-wave processes $\pi\pi \to \pi\pi, K\overline{K}, \eta\eta, \eta\eta'$ [1] and also with adding the data from DM2 and Mark3 on decays $J/\psi \to \phi\pi\pi, \phi K\overline{K}$ [2] were performed to study f_0-mesons lying below 1.9 GeV. First, it is shown that the data admit two possibilities for parameters of $f_0(600)$ with mass close to the ρ-meson mass, and with the total width about 600 and 1000 MeV. These two possibilities are related to two solutions admitted by the data, for the phase shift of the $\pi\pi$-scattering amplitude below 1 GeV: "up" and "down". As to a combined description of the considered processes, it is impossible to prefer any of these solutions. However, the "up" solution remarkably accords with the prediction on the basis of the mended symmetry by Weinberg (S.Weinberg, PRL **65** (1990) 1177) as to the mass and the width, the "down" one as to mass.
Furthermore, unlike our previous analyses, we considered all relevant possibilities of the representation of resonances by pole clusters (of poles and zeros on the Riemann surface). In our model-independent approach the 3-channel resonances (depending on their nature) are represented by seven types of the pole clusters. It is shown that for the "up" solution there are four scenarios of the representation of resonances $f_0(1370)$, $f_0(1500)$ (as a superposition of two states, broad and narrow) and $f_0(1710)$ ($f_0(600)$ and $f_0(980)$ are given by the pole clusters of the same types in all cases) giving about a similar description of the above processes and, however, quite different parameters of some resonances. For $f_0(600)$, $f_0(1370)$ and $f_0(1710)$ a spread of values is obtained for the masses and widths 605-735 and 567-686 MeV, 1326-1404 and 223-345 MeV, and 1751-1759 and 118-207 MeV, respectively. On the other hand, the results for $f_0(980)$ and $f_0(1500)$ are more stable and confirm conclusions of our previous analyses. Note the mass value about 1007 MeV for $f_0(980)$, whereas in analyses only of the J/ψ-decays one obtains this mass value below the $K\overline{K}$ threshold.
Finally, we performed the analysis of the above processes excluding the state $f_0(1370)$. This is needed by the uncertain situation when, e.g., D. Bugg (Eur. Phys.J. C**52** (2007) 55; arXiv: 0710.4452 [hep-ex]) has indicated a number of data requiring apparently the existence of $f_0(1370)$: the Crystal Barrel data on $\bar{p}p \to \eta\eta\pi^0$ and on $\bar{p}p \to 3\pi^0$ also the BES data on $J/\psi \to \phi\pi^+\pi^-$, and the GAMS data for $\pi^+\pi^- \to \pi^0\pi^0$ at large $|t|$. On the other hand, e.g., in works by W. Ochs and P. Minkowski (arXiv:1001.4486v1 [hep-ph]; EPJ C**9** (1999) 283; hep-ph/0209225) one did not find evidence for the existence of $f_0(1370)$. Our conclusion: the existence of $f_0(1370)$ does not contradict the considered data. The description of the $\pi\pi$ scattering is a bit improved whereas the one of the $\pi\pi \to K\overline{K}$ process is made worse, especially as to the phase shift; and the J/ψ-decays are also described a little worse. An interpretation of $f_0(1370)$ as dominated by the $s\bar{s}$ component explains why one did not find this state considering only the $\pi\pi$ scattering.
The model-independent analysis of the isovector P-wave of the 3-channel $\pi\pi$-scattering was performed [5, 6], i.e., in addition to the $\pi\pi$- and $\rho 2\pi$-threshold, we took into account the threshold of the third effective channel in the corresponding uniformizing variable. This threshold is found to be at about 1512.35 MeV and interpreted by us as related to the $\rho\sigma$ channel. Since the σ-meson still provokes many questions, it is interesting to observe the $\rho\sigma$ final state in the π-π collisions. There is confirmed a conclusion of our previous 2-channel analysis of the $\pi\pi$-scattering that the first ρ-like meson is $\rho(1250)$ ($m_{res} = 1274.7 \pm 32.4$ MeV,

$\Gamma_{tot} = 304.3 \pm 23.6$ MeV). A possible $q\bar{q}$ classification is proposed with taking $\rho(1250)$ into account. In this classification, $\rho(1450)$ (whose existence does not contradict the data) might be the ${}^3D_1$ $q\bar{q}$ state with the possible SU(3) partners: the isodoublet $K^*(1680)$ and the isoscalars $\omega(1650)$ and $\phi(1680)$.
Since the result on $\rho(1250)$ leads to an important conclusion as to mainstream quark models (see a discussion in [5]), it is worthwhile to check if this result is supported by investigation of other mesonic sectors. Considering the (J, M^2)-plot for the daughter ρ-trajectory related to $\rho(1250)$ one can conclude that there should exist a 1^+3^{--}-state at about 1950 MeV – "$\rho_3(1950)$". Our analysis [7] of the F-wave $\pi\pi$ scattering data (B. Hyams et al., NP B **64** (1973) 134) indicates that, except the known $\rho_3(1690)$ (in our analysis $m_{res} \approx 1703$ MeV and $\Gamma_{tot} \approx 175$ MeV), there might be one more state lying above 1830 MeV. Since the $\pi\pi$ scattering data above 1890 MeV are absent, it is impossible to say something serious about the parameters of this state. However, $\rho_3(1950)$ does not contradict the data and even improves a little the obtained parameters of $\rho_3(1690)$ and its branching ratios when comparing them with the PDG tables.
All calculations were performed in parallel in the program Mathematica and Fortran program which utilizes MINUIT.

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MODEL DESCRIPTION OF TRANSITION FORM FACTORS

A.E. Dorokhov

The present study is devoted to the so-called BABAR puzzle. New very precise data were obtained by the BABAR collaboration for the photon-pion transition form factor in a very wide kinematic region up to large photon virtualities $Q^2 \approx 40$ GeV^2. The data overshoot the asymptotic limit for $Q^2 F_{\pi\gamma\gamma^*}(Q^2)$, predicted by Brodsky and Lepage, and have a tendency to grow further. Both facts are in strong contradiction with the standard QCD factorization approach, which constitutes the BABAR puzzle.

The main problem is the unstopped growth of the new data points for $Q^2 F_{\pi\gamma\gamma^*}(Q^2)$ that is inconsistent with the predicted $Q^2 F_{\pi\gamma\gamma^*}(Q^2) \to$ constant, following from simple asymptotic properties of the massless quark propagator. The key point, to solve this problem, is to consider the properties of the pion vertex function $F(k_1^2, k_2^2)$ which is an analog of the light-cone pion wave function. There are two possibilities for the momentum dependence of the pion vertex function. In the limit, when one quark virtuality k_1^2 goes to infinity, and the other, k_2^2, remains finite, the vertex function may not necessarily tend to zero. When it goes to zero, the pion DA $\varphi_\pi(x)$, which is a functional of the pion vertex function, is zero at the endpoints, $\varphi_\pi(0) = \varphi_\pi(1) = 0$, with either strong or weak suppression in the neighborhood of the endpoints $x = 0$ and $x = 1$. For the situation of strong suppression, the asymptotic $1/Q^2$ behavior of the pion form factor in asymmetric kinematics $(Q_1^2 = Q^2, Q_2^2 = 0)$ is developed very early, in contradiction with the BABAR data. For weak suppression (resembling a flat distribution amplitude of the pion), the asymptotic $1/Q^2$ behavior is developed quite late, and can give a reasonable description of the data in the BABAR region with the $\ln Q^2/Q^2$ behavior in this region. For the other case of non-vanishing pion vertex function in the above limit, the pion DA $\varphi_\pi(x)$ is not zero at the endpoints and, therefore, the asymptotic $\ln Q^2/Q^2$ behavior persists over the whole range, in particular, in the BABAR region.

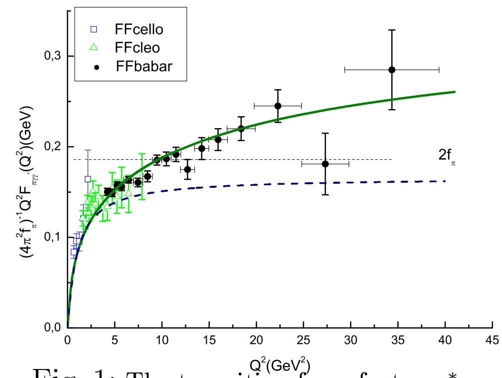

Fig. 1: The transition form factor $\gamma^*\gamma \to \pi^0$. The data are from the CELLO (empty squares), CLEO (empty triangles) and BABAR Collaborations (filled circles). The solid line is the model of this work, the dashed line is the Brodsky-Lepage prediction, the short-dashed line is massless QCD asymptotic limit .

In order to fit the available data on the photon-pion transition form factor from CELLO, CLEO and BABAR, we have analyzed the parameter space of two examples of nonperturbative models, motivated by the instanton and the chiral quark models, characterized by the two parameters, dynamical quark mass M_q and the parameter of nonlocality Λ. The main conclusion is that the fit to the data requires a quite small dynamical quark mass $M_q \approx 125$

MeV with rather small uncertainty (Fig. 1). As a consequence, the parameter of nonlocality, that fits the pion decay constant f_π, is very small, $\Lambda \sim 0.01$ GeV^{-2}. Thus, one has an almost local quark model with very flat regulators in momentum space which considerably diminishes the difference between the nonperturbative models considered in this work. We would like to point out that the problem of QCD radiative corrections and evolution has to be considered in addition. Concluding we may say that the BABAR data being unique in their accuracy and covering a very wide kinematic range are consistent with considerations based on nonperturbative QCD dynamics and may indicate specific properties of the pion wave function.

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DEVELOPMENT AND APPLICATIONS OF THE NAMBU-IONA-LASINIO MODEL

A.B. Arbuzov, M.K. Volkov, E.A. Kuraev, Yu.M. Bystritskiy

The development of the Nambu-Jona-Lasinio (NIL) model and its application to the description of strong interactions of mesons at low energies was continued. The model is based on the chiral symmetry principle and exploits the spontaneous symmetry breaking mechanism. It is well known for a very good description of the light meson spectra and interactions. On the other hand, the NJL model is an effective nonrenormalizable field theory limited for applications at energies below 1-2 GeV. An important step in understanding and justification of the model was performed in papers [1,2]: by means of the Bogoliubov compensation method it was shown for the first time that QCD at low energies can be reduced to the NJL model. The nonlocal version of the NJL model with Polyakov loops (PNJL) was applied [3] for the description of strong interactions in the conditions of dense matter at nonzero temperature which occur in heavy ion collisions. Within the framework of the NJL model, many radiative decays of pseudoscalar, scalar and vector mesons were described [4]. It was shown that for a correct theoretical description of scalar meson radiative decays, in particular, for the two-photon decays f0(600)?2?, a0(980)?2?, and f0(980)?2?, besides quark loops, it is necessary to take into account also the meson ones, i.e., to work in the 1/Nc approximation. That allows one to get good agreement between the theoretical predictions obtained in the NJL model with the experimental data [5-7]. The results of the calculations were further used for the description of several processes of meson production in colliding electron-positron beams at such machines as VEPP-2M, VEPP-2000 (Novosibirsk), BEPC (Beijing) and DAFNE (Frascati) [8]. A new theoretical description of the experimental data on the pion polarizability was also obtained [9]. During the 2010 year, within the framework of the extended NJL model several predictions were made for radiative decays with participation of pions, rho-, and omega-mesons in the ground and in the first radial excited states [10,11].

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INSTANTON LIQUID MODEL AND HEAVY-ION COLLISIONS

S.V. Molodtsov

The possibility of self-consistent determination of instanton liquid parameters was considered together with the defnition of optimal pseudo-particle configurations and comparing the various pseudo-particle ensembles. The weakening of repulsive interactions between pseudo-particles was argued and estimated [1].

The origin of the lightest scalar mesons was studied in the framework of the instanton liquid model of the QCD vacuum. The impact of phonon-like vacuum excitations on the σ-meson features was qualitatively analyzed. In particular, it was noticed that the changes produced in the scalar sector may unexpectedly become quite considerable in spite of insignificant values of corrections to the dynamical quark masses and then the medley of σ-meson and those excitations may reveal as broad resonance states of different masses [2].

The quark behavior while influenced by a strong stochastic gluon field was analyzed. An approximate procedure for calculating the effective Hamiltonian was developed and the corresponding ground state within the Hartree-Fock-Bogoliubov approach was found. The comparative analysis of various Hamiltonian models was given and a transition to the chiral limit in the Keldysh model was discussed in detail [3], [4].

The parameter responsible for the choice of quantum operator representation was discussed, and with the help of the variational principle its optimal value was established. Interpretation of deviations from equilibrium value as a dynamical variable leads to an idea of a scalar field of exceptional nature responsible for ordering of the operators [5].

We study meson correlation functions for several models with a four-fermion quark interaction. We show that although the system average energy and quark condensate, as previously noted, are singular, the meson observables are finite, completely recognizable, and comparable to the experimental data on the energy scale. This permits using a wide set of Hamiltonians to model nonequilibrium states of quark and hadronic systems, which is a relevant problem in studying the physics of ultrarelativistic heavy-ion collisions. We obtain analytic expressions for the meson correlation functions in the Keldysh model [6].

Considering quarks as quasiparticles of the model Hamiltonian with four-fermion interaction we study a response to the process of filling up the Fermi sphere with quarks, calculate the vacuum pressure and demonstrate the existence of filled-in state degenerate with the vacuum one [7], [8], [9].

The spectra and e+e- decay widths of the heavy quarkonia as a function of the temperature, generated in the ultrarelativistic heavy-ion collisions was discussed. The fluctuations of the vacuum gluon fields was estimated within the instanton liquid model approach. It is noticed that the mentioned parameters can be applied as an indicator of the temperature of gluons [10].

Some features of hot and dense gas of quarks which are considered as quasiparticles of the model Hamiltonian with four-fermion interaction were studied. Being adapted to the Nambu-Jona-Lasinio model this approach allows us to accommodate a phase transition similar to the nuclear liquid-gas one at the proper scale. It allows us to argue the existence of the mixed phase of vacuum and baryonic matter (even at zero temperature) as a plausible scenario of chiral symmetry (partial) restoration [11].

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THE TMD PDF RELATION IN THE BAG AND OTHER MODELS

H. Avakian[1], A.V. Efremov[2], P. Schweitzer[3], O.V. Teryaev[2] and P. Zavada[4]

[1] *Thomas Jefferson National Accelerator Facility, Newport News, VA 23606, U.S.A.*
[2] *Joint Institute for Nuclear Research, Dubna, 141980 Russia*
[3] *Department of Physics, University of Connecticut, Storrs, CT 06269, U.S.A.*
[4] *Institute of Physics AS CR, Na Slovance 2, CZ-182 21 Prague 8, Czech Rep.*

Introduction. The Transverse Momentum Dependent distributions functions (TMDs) are a generalization of parton distribution functions (PDFs) promising to extend our knowledge of the nucleon structure far beyond what we have learned from PDFs about the longitudinal momentum distributions of partons in the nucleon. In addition to the latter, TMDs carry also information on transverse parton momenta and spin-orbit correlations. TMDs enter the description of leading-twist observables in deeply inelastic reactions on which data are available like: semi-inclusive DIS (SIDIS) Drell-Yan process or hadron production in e^+e^- annihilations. The interpretation of these data is not straight-forward though. In SIDIS one deals with convolutions of a priori unknown transverse momentum distributions in nucleon and fragmentation process, and in practice is forced to *assume* models for transverse parton momenta such as the Gaussian Ansatz. In the case of subleading twist observables, one moreover faces the problem that several twist-3 TMDs and fragmentation functions enter the description of one observable (we recall that presently factorization is not proven for subleading-twist observables).

In this situation, information from models is valuable for several reasons. Models can be used for direct estimates of observables, though it is difficult to reliably apply the results, typically obtained at low hadronic scales, to experimentally relevant energies. Another aspect concerns relations among TMDs observed in models. Such relations, especially when supported by several models, could be helpful — at least for qualitative interpretations of the first data. Furthermore, model results allow one to test assumptions made in literature, such as the Gaussian Ansatz for transverse momentum distributions or certain approximations. In addition to such practical applications, model studies are also of interest because they provide important insights into nonperturbative properties of TMDs.

The purpose of this draft is to present main results for the TMD relation in the framework of the MIT bag and other models. More details and references can be found in [1].

TMDs in the bag model. We begin with the quark-quark correlation function in polarized nucleon from [2] with the MIT bag model quark field in the lowest mode

$$\varphi_m(\vec{k}) = i\sqrt{4\pi}NR_0^3 \begin{pmatrix} t_0(k)\chi_m \\ \vec{\sigma}\cdot\hat{k}\, t_1(k)\chi_m \end{pmatrix} , \tag{1}$$

where $\hat{k} = \vec{k}/k$ with $k = |\vec{k}|$ and N is a normalization factor, t_0 represents the S-wave component, whereas t_1 represents the P-wave component of the proton wave functions. Moreover, we assume $SU(6)$ spin-flavor symmetry of the proton wave function such that spin-independent TMDs of definite flavor are given in terms of respective 'flavor-less' expressions multiplied by a 'flavor factor' N_q, and spin-dependent TMDs of definite flavor follow from multiplying the respective 'flavor-less' expressions by a 'spin-flavor factor' P_q with $N_u = 2$, $N_d = 1$, $P_u = \frac{4}{3}$, $P_d = -\frac{1}{3}$.

Since there are no explicit gluon degrees of freedom, T-odd TMDs vanish in this model and there stay 14 (twist-2 and 3) T-even TMDs.
In the notation introduced above, the results for the T-even leading twist TMDs are given by (A is a common factor)

$$
\begin{aligned}
f_1^q(x,k_\perp) &= N_q A\big[t_0^2+2\widehat{k}_z\,t_0t_1+t_1^2\big] & g_1^q(x,k_\perp) &= P_q\,A\big[t_0^2+2\widehat{k}_z\,t_0t_1+(2\widehat{k}_z^2-1)\,t_1^2\big]\\
h_1^q(x,k_\perp) &= P_q\,A\big[t_0^2+2\widehat{k}_z\,t_0t_1+\widehat{k}_z^2\,t_1^2\big] & g_{1T}^{\perp q}(x,k_\perp) &= P_q\,A\big[\ \ 2\widehat{M}_N(t_0t_1+\widehat{k}_z\,t_1^2)\big]\\
h_{1L}^{\perp q}(x,k_\perp) &= P_q\,A\big[-2\widehat{M}_N(t_0t_1+\widehat{k}_z\,t_1^2)\big] & h_{1T}^{\perp q}(x,k_\perp) &= P_q\,A\big[-2\widehat{M}_N^2\,t_1^2\big]
\end{aligned}
$$

and for the subleading twist TMDs we obtain

$$
\begin{aligned}
e^q(x,k_\perp) &= N_q A\big[t_0^2-t_1^2\big] & f^{\perp q}(x,k_\perp) &= N_q A\big[2\widehat{M}_N\,t_0t_1\big]\\
g_T^q(x,k_\perp) &= P_q\,A\big[t_0^2-\widehat{k}_z^2\,t_1^2\big] & g_L^{\perp q}(x,k_\perp) &= P_q\,A\big[2\widehat{M}_N\,\widehat{k}_z\,t_1^2\big]\\
g_T^{\perp q}(x,k_\perp) &= P_q\,A\big[2\widehat{M}_N^2\,t_1^2\big] & h_L^q(x,k_\perp) &= P_q\,A\big[t_0^2+(1-2\widehat{k}_z^2)t_1^2\,\big]\\
h_T^{\perp q}(x,k_\perp) &= P_q\,A\big[2\widehat{M}_N\ t_0t_1\big] & h_T^q(x,k_\perp) &= P_q\,A\big[-2\widehat{M}_N\widehat{k}_z\,t_1^2\big]
\end{aligned}
$$

In the bag model, there are 9 *linear* relations among the 14 TMDs, which can be written as follows:

$$\mathcal{D}^q\, f_1^q(x,k_\perp)+g_1^q(x,k_\perp)=2h_1^q(x,k_\perp) \qquad (2)$$

$$\mathcal{D}^q\, e^q(x,k_\perp)+h_L^q(x,k_\perp)=2g_T^q(x,k_\perp) \qquad (3)$$

$$\mathcal{D}^q\, f^{\perp q}(x,k_\perp)=h_T^{\perp q}(x,k_\perp) \qquad (4)$$

$$g_{1T}^{\perp q}(x,k_\perp)=-h_{1L}^{\perp q}(x,k_\perp) \qquad (5)$$

$$g_T^{\perp q}(x,k_\perp)=-h_{1T}^{\perp q}(x,k_\perp) \qquad (6)$$

$$g_L^{\perp q}(x,k_\perp)=-h_T^q(x,k_\perp) \qquad (7)$$

$$g_1^q(x,k_\perp)-h_1^q(x,k_\perp)=h_{1T}^{\perp(1)q}(x,k_\perp) \qquad (8)$$

$$g_T^q(x,k_\perp)-h_L^q(x,k_\perp)=h_{1T}^{\perp(1)q}(x,k_\perp) \qquad (9)$$

$$h_T^q(x,k_\perp)-h_T^{\perp q}(x,k_\perp)=h_{1L}^{\perp q}(x,k_\perp)\,, \qquad (10)$$

where $j^{(1)q}(x,k_\perp)=\frac{k_\perp^2}{2M_N^2}\,j^q(x,k_\perp)$ and $\mathcal{D}^q=\frac{P_q}{N_q}$.
Why are there 9 linear relations? In fact, naively, one could have expected even more relations, since all TMDs are expressed in terms of only two functions, t_0 and t_1. However, having *linear relations* in mind, the combinations t_0^2, t_0t_1, t_1^2 are to be considered as independent structures. Also, $\widehat{k}_z t_0 t_1$ and $\widehat{k}_z t_1$, are linearly independent as there is no way of relating one with the other in a model-independent way. However, there are *nonlinear relations*, for example,

$$h_1^q(x,k_\perp)\,h_{1T}^{\perp q}(x,k_\perp) \;=\; -\frac{1}{2}\left[h_{1L}^{\perp q}(x,k_\perp)\right]^2, \qquad (11)$$

$$g_T^q(x,k_\perp)\,g_T^{\perp q}(x,k_\perp) \;=\; \frac{1}{2}\left[g_{1T}^{\perp q}(x,k_\perp)\right]^2-g_{1T}^{\perp q}(x,k_\perp)\,g_L^{\perp q}(x,k_\perp)\,. \qquad (12)$$

Equations (11, 12) are independent in the sense that it is impossible to convert one into the other upon use of the linear relations (2–10).
With the 9 linear relations (2–10) and the 2 non-linear relations (11, 12) we find altogether 11 relations among 14 TMDs in the bag model.

How general are quark model relations among TMDs? The deeper reason, why in the bag model relations among TMDs appear, is ultimately related to Melosh rotations which connect longitudinal and transverse nucleon and quark polarization states in a Lorentz-invariant way [3]. An important issue, when observing relations among TMDs in a model, concerns their presumed validity beyond that particular model framework. For that it is instructive to compare first to other models.

- Eq. (2): its $k_\perp$-integrated version was discussed in the bag model in [4] and [5, 6] and in the light-cone constituent models in [7]. The unintegrated version was discussed in the bag and light-cone constituent models [8, 9].
- Eq. (3): its integrated version was observed in the bag model previously in [5].
- Eq. (5): was first observed in the spectator model of [10] and later also in the light-cone constituent models [8] and the covariant parton model of Ref. [11].
- Eq. (7): was found in the spectator model of Ref. [10].
- Eq. (8): was first observed in the bag [9]. It is valid also in the spectator [10], light-cone constituent [8], and the covariant parton [11] models.
- Eqs. (4, 6, 9, 10): are new in the sense of not having been mentioned previously in literature. But the latter 3 are satisfied by the spectator model results from [10].
- The nonlinear relation (11), which connects all T-even, chirally-odd leading-twist TMDs was observed in the covariant parton model approach [11]. Eq. (12) was not discussed so far in literature.

The detailed comparison of the models where these relations hold and where they are violated, gives some insight into the question to which extent these relations are model-dependent.

Let us discuss first Eqs. (2–4) which connect polarized and unpolarized TMDs. For these relations SU(6)-spin-flavor symmetry is necessary, but not sufficient. For example, the spectator model of [10] is SU(6) symmetric. However, it does not support (2–4) which are spoiled by the different masses of the (scalar and axial-vector) spectator diquark systems. Also, (2, 3) are not supported in the covariant parton model approach of [11]. However, also in that approach it is possible to 'restore' these relations by introducing additional, restrictive assumptions, see [11] for a detailed discussion. We conclude that the relations (2–4) require strong model assumptions. It is difficult to estimate to which extent such relations could be useful approximations in nature, though they could hold in the valence-x region with an accuracy of (20–30) % (see [12]).

From the point of view of model dependence, it is 'safer' [9] to compare relations which include only polarized or only unpolarized TMDs. We know no example for the latter, however, the relations (5–10) are of the former type. It is gratifying to observe that these relations are satisfied not only in the bag model, but also in the spectator model version of Ref. [10]. The relations among the leading twist TMDs, Eqs. (5, 8), hold also in the light-cone constituent [8] and the covariant parton [11] models.

Of course, quark model relations among TMDs have limitations, even in quark models. In [13], various versions of spectator models were used, and in some versions the relations were not supported (5, 8). Also, the quark-target model [14] did not support the relations (5, 8).

Finally, in QCD none of such relations is valid, and all TMDs are independent structures. It would be interesting to 'test' such quark model relations in other models, lattice QCD, and in experiment.
It is worthwhile to discuss in some more detail one particularly interesting relation which can be obtained in this way. By eliminating the transverse moment of the pretzelosity distribution from Eqs. (8, 9) and integrating over transverse momenta we obtain

$$g_1^q(x) - h_1^q(x) = g_T^q(x) - h_L^q(x) = h_{1T}^{\perp q}(x)\,. \tag{13}$$

This relation holds also in its unintegrated form. There are several reasons why this relation is interesting.
First, it involves only collinear parton distribution functions, which is the only relation of such type in the bag model. The QCD evolution equation for all these functions is different, which shows the limitation of this relation: even if for some reason (13) was valid in QCD at a certain renormalization scale μ_0, it would break down at any other scale $\mu \neq \mu_0$.
Second, for the first Mellin moment this relation is valid model-independently presuming the validity of the Burkardt-Cottingham sum rule and an analog sum rule for $h_L^q(x)$ and $h_1^q(x)$. In QCD there are doubts especially concerning the validity of the Burkardt-Cottingham sum rule. However, it is valid in many models such as the bag [4] or the chiral quark soliton model [15].
Third, it would be interesting to learn whether (13) is satisfied in nature approximately. Also, this relation can be tested on the lattice, especially for low Mellin moments and in the flavour non-singlet case. Lattice QCD calculations for Mellin moments of $g_T^q(x)$ were reported in [16].
Forth, relation (13) can be tested in models where collinear parton distribution functions were studied. Some results can be found in literature. For example, calculations of parton distribution functions in the bag models [4, 5] support this relation. Moreover, the spectator model [10] supports this relation. One counter-example is known though: the chiral quark-soliton model does not support this relation [15, 17]. The models where (13) holds include only the components in the nucleon wave-function with the quark orbital angular momenta up to $L = 0, 1, 2$ at most. The chiral quark soliton model, which does not support (13), contains *all* quark angular momenta $L = 0, 1, 2, 3, 4, \ldots$.
Fifth, an important aspect of model relations is that they inspire interpretations. The relation (13) means that the difference between g_T^q and h_L^q is a 'measure of relativistic effects in the nucleon' to the same extent as the difference between helicity and transversity. Both these differences are related to the transverse moment of pretzelosity, see Eqs. (8, 9) and [9].

Pretzelosity and quark orbital angular momentum. In quark models, in contrast to gauge theories, one may unambiguously define the quark orbital angular momentum operator as $\hat{L}_q^i = \bar{\psi}_q \varepsilon^{ikl} \hat{r}^k \hat{p}^l \psi_q$ where for clarity the 'hat' indicates a quantum operator. This definition follows (in the absence of gauge fields) uniquely, for instance, from identifying that part of the generator of rotations not associated with the intrinsic quark spin. It will be convenient to introduce a 'non-local version' of this operator, by defining $\hat{L}_q^i(0,z) = \bar{\psi}_q(0)\varepsilon^{ikl}\hat{r}^k\hat{p}^l\psi_q(z)$. In the bag model it is convenient to work in the momentum space where $\hat{r}^k = i\,\frac{\partial}{\partial p^k}$ and $\hat{p}^l = p^l$. Next let us define the quantity

$$L_q^j(x,p_T) = \int \frac{\mathrm{d}z^- \mathrm{d}^2\vec{z}_T}{(2\pi)^3}\, e^{ipz}\; N(P,S^3)|\bar{\psi}_q(0)\varepsilon^{ikl}\hat{r}^k\hat{p}^l\psi_q(z)|N(P,S^3)\Big|_{z^+=0,\,p^+=xP^+}\,. \tag{14}$$

In order to find a connection to TMDs we must consider a longitudinally polarized nucleon, choosing the polarization vector as $S = (0, 0, 1)$, for definiteness, and we must focus on the $j = 3$ component in (14), i.e. on the component of the angular momentum operator along the light-cone space-direction.
Evaluating expression (14) in the bag model we obtain just as in [18, 19]

$$L^3_q(x, p_T) = (-1)\, h_{1T}^{\perp(1)q}(x, p_T)\ . \tag{15}$$

In order to demonstrate the consistency of this result we compute the contribution to the total angular momentum of the nucleon J^3_q due to flavour q; J^3_q is composed of contributions from intrinsic quark spin, $S^3_q = \frac{1}{2}\int \mathrm{d}x g_1^q(x)$, and quark orbital angular momentum $L^3_q = \int \mathrm{d}x \int \mathrm{d}^2\vec{p}_T L^3_q(x, p_T) = (-1) \int \mathrm{d}x\ h_{1T}^{\perp(1)q}(x)$. We obtain $2J^3_q = 2S^3_q + 2L^3_q = P_q$ and $2J_u + 2J_d = 1$.
It is important to observe that the relation of pretzelosity and orbital angular momentum, Eq. (15), is at the level of matrix elements of operators, and there is no a priori operator identity which would make such a connection.

Conclusions. We presented a study of a complete set of relations of T-even leading- and subleading-twist TMDs in the MIT bag model and to what extent they are supported in other quark models. Special attention was paid to the relation of the difference g_1^q and h_1^q to the (1)-moment of pretzelosity and its relation to quark orbital angular momentum. It is interesting to ask whether a quark model relation of type (15) may inspire a way to establish a rigorous connection between TMDs and OAM in QCD? We hope our results will stimulate further studies in quark models.
Sorrily, the limiting frame of the draft does not alow us to dwell on other interesting questions like Lorentz invariant relations, positivity inequalities, the Wandzura-Wilczek-type approximations, which are shown to be all valid, and numerical results in the model. All these questions can be found in paper [1].

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TRANSVERSE MOMENTUM DEPENDENT PARTON DISTRIBUTIONS, GAUGE INVARIANCE AND DUALITY

I.V. Anikin, I.O. Cherednikov, N.G. Stefanis, O.V. Teryaev

Unintegrated (transverse-momentum dependent, TMD) parton distribution functions (PDFs) are important ingredients of the QCD factorization approach to the semi-inclusive high-energy hadronic processes, such as semi-inclusive DIS or the Drell-Yan process—where more than one final or initial hadron is detected and its transverse momentum is observed. The detailed knowledge of these nonperturbative objects is crucial for the analysis of the experimental data obtained in this sort of experiments with polarized and unpolarized hadrons as HERMES, COMPASS, BELLE, BaBar, which have been performed and planned at DESY, GSI, RHIC, LHC and future facilities like EIC and 12 GeV-upgraded JLab. In particular, strong TMD program is developed for the Electron-Ion Collider (EIC). The EIC is envisioned mostly as a high-energy and high-luminosity QCD machine which will address open questions about the dynamics of quarks and gluons: unpolarized and polarized distributions, tomography, propagation in nuclear matter, fragmentation, various aspects of nonperturbative dynamics of partons, nonlinear parton dynamics, etc.

In the years 2009-2010, we continued our study of the theory of TMD PDFs, concentrating mostly on their renormalizability, gauge invariance, evolution equations, factorization and universality. In particular, the renormalization-group properties of TMD PDFs in the light-cone gauge with the Mandelstam-Leibbrandt prescription for the gluon propagator have been considered in detail. An expression for the transverse component of the gauge field at light-cone infinity, which plays a crucial role in the description of the final-/initial-state interactions in the light-cone axial gauge, was obtained. The leading-order anomalous dimension was calculated in this gauge and the relation to the results obtained in other gauges is worked out. It is shown that using the Mandelstam-Leibbrandt prescription the ensuing anomalous dimension does not receive contributions from extra rapidity divergences related to a cusped junction point of the Wilson lines [1].

Another branch of our TMD project is the development of a new framework for the TMD PDFs, based on a generalized conception of gauge invariance which includes into the Wilson lines the spin-dependent Pauli term $\sim F^{\mu\nu}[\gamma_\mu, \gamma_\nu]$. We discussed the relevance of this non-minimal term for unintegrated parton distribution functions, pertaining to spinning particles, and analyze its influence on their renormalization-group properties. It is shown that while the Pauli term preserves the probabilistic interpretation of twist-two distributions—unpolarized and polarized—it gives rise to additional pole contributions to those of twist-three. The anomalous dimension induced this way is a matrix calling for a careful analysis of evolution effects. Moreover, it turns out that the crosstalk between the Pauli term and the longitudinal and the transverse parts of the gauge fields accompanying the fermions induces a constant but process-dependent phase which is the same for leading and subleading distribution functions. Feynman rules for the calculation with gauge links containing the Pauli term are derived [2].

The electromagnetic gauge invariance of the hadron tensor of the Drell-Yan process with one transversely polarized hadron was investigated with the contour gauge for gluon fields. The prescription for the gluonic pole in the twist 3 correlator, being the complementary way to describe TMD Sivers function, is related to causality property for exclusive hard processes. As a result, extra contributions which naively do not have an imaginary phase were found. This enhanced the single spin asymmetry by a factor of two [3].

The deeply virtual Compton scattering off a spin-one particle was considered. We discuss the role of twist three contributions for restoring the QED gauge invariance of the amplitude. We consider both kinematic and dynamic sources of twist three generalized parton distributions [4].
We compare the proposed higher twist factorization method with the covariant method formulated in the coordinate space based on the operator product expansion. We prove the equivalence of two approaches by computing the impact factor for the transition of virtual photon to transversally polarized rho-meson up to the twist 3 accuracy [5].
We study the phenomenon of duality in hard exclusive reactions to which QCD factorization applies. In the QCD case, the appearance of duality is sensitive to the particular nonperturbative model applied and can, therefore, be used as an additional adjudicator [6].

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GENERALIZED PARTON DISTRIBUTIONS AND HARD ELECTROPRODUCTION

S.V. Goloskokov, P. Kroll

We analyze light meson electroproduction at intermediate energies where the quark contributions are essential. Our calculations are carried out on the basis of the handbag factorization approach. The scattering amplitude is factorized at large photon virtualities Q^2 into hard meson electroproduction off partons and Generalized Parton Distributions (GPDs). It was shown that this approach works well for the processes which include ϕ, ρ^0, ω, ρ^0 production determined by the standard GPDs. The production of charged mesons which are determined in terms of transition GPDs like $p \to n$ or $p \to \Sigma$ can be studied in our approach too. GPDs are modeled with the help of the double distribution representation. In our model [1], we modify the leading twist amplitude by including the quark transverse degrees of freedom and the Sudakov corrections in the hard subprocess amplitude. Taking into account the quark transverse momentum gives a possibility to calculate the higher twist TT amplitude which is essential in the description of spin effects.

The analysis of light meson electroproduction in the model was carried out. We consider the gluon, sea and valence quark GPD contributions [1]. This permits us to analyze vector meson production from moderate energies typical of HERMES and COMPASS up to HERA energies [2]. It was found that the quark contribution is mostly essential at $W \sim 5$GeV. At lower energies the quark contribution decreases, as well as the gluon and sea one, and the ρ cross section falls at energies $W \leq 5$GeV [3]. This is in contradiction with CLAS results which show essential growth of σ_ρ in this energy range. At the same time, the model describes fine ϕ production at CLAS [4]. This means that we have a problem only with the valence quark contribution at low JLAB energies.

We estimated cross sections and spin observables for various vector mesons. Our results are in good agrement with HERMES experiments. Predictions for physical observables at COMPASS energies were made [3].

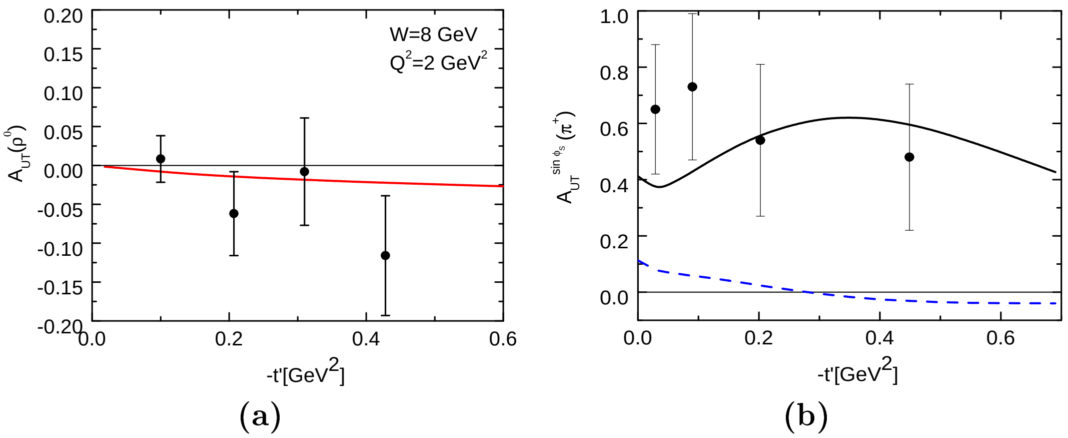

Fig. 1: **(a)** Predictions for A_{UT} asymmetry $W = 8$GeV and $Q^2 = 2$GeV^2. Preliminary data are from COMPASS. **(b)** The $sin\phi_s$ moment of A_{UT} asymmetry π^+ production at HERMES. The dashed line is obtained by neglecting the twist-3 contribution.

The GPD E, which is responsible for proton helicity flip, is not well known as yet. We constructed the GPD E from double distributions and constrained it by the Pauli form

factors of the nucleon, positivity bounds and sum rules. This give a possibility to estimate spin effects on the transversally polarized target, like the A_{UT} asymmetry. The predictions for COMPASS on t- dependence of A_{UT} asymmetry at $W = 8$GeV are shown in Fig. 1.a and are in good agreement with preliminary COMPASS data. Our predictions for the asymmetry at $W = 5$GeV and $W = 10$GeV were given for the ρ^0, ω, ρ^+, K^{*0} mesons [5]. The predicted A_{UT} asymmetry in the ω production is negative and not small, about -10% at HERMES and COMPASS. This is caused by the fact that GPD E^u and E^d do not compensate each other in this reaction. The predictions for ρ^+ asymmetry are positive and rather large $\sim 40\%$.
Information on the parton angular momenta can be obtained from the Ji sum rules. In our model, we found not small angular momenta for u quarks and gluons [5]

$$< J_v^u >= 0.222, \quad < J_v^d >= -0.015, \quad < J^g >= 0.214, \tag{1}$$

which are not far from the lattice results.
In the pion electroproduction the amplitude with longitudinally polarized photons dominates. These amplitudes are calculated using a hard subprocess and the the pion contribution which is treated with the fully experimentally measured electromagnetic form factor of the pion [6]. Using this model approach we calculate all amplitudes with exception of $\mathcal{M}_{0-,++}$. The last amplitude should be constant at small momentum transfer. However, in the handbag approximation it vanished at $t = 0$. This problem was solved by including a twist-3 contribution to the amplitude $\mathcal{M}_{0-,++}$. In order to estimate this effect, we use a mechanism that consists of the helicity-flip GPD H_T and the twist-3 pion wave function.
Our results [6] on the cross section and six moments of spin asymmetries for the polarized target are in good agrement with HERMES experimental data. It was found that twist-3 effects were very essential in the description of the partial π^+ cross section and in the spin asymmetries Fig. 1.b [3]. The obtained results are important in analyses and interpretation of HERMES and COMPASS data in terms of GPDs.

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RADIATIVE CORRECTIONS AND PARTICLE PRODUCTION AT THE COLLIDERS

A.B. Arbuzov, V. V. Bytev, E. A. Kuraev, E. Tomasi-Gustafsson, and Y. M. Bystritskiy

Radiative corrections In 2009, we calculated [4] the QED and QCD radiative corrections to the charged lepton energy distributions in the dominant semileptonic decays of the top quark $t \to bW^+ \to b(\ell^+\nu_\ell)$ $(\ell = e, \mu, \tau)$ in the Standard Model(SM), and for the decay $t \to bH^+ \to b(\tau^+\nu_\tau)$ in an extension of the SM having a charged Higgs boson $H^\pm$ with its mass $m_{H^\pm} < m_t - m_b$. The QCD corrections are calculated in the leading and next-to-leading logarithmic approximations, but the QED corrections are considered in the leading logarithmic approximation only. These corrections are numerically important for precisely testing the universality of the charged current weak interactions in t-quark decays. As the τ^+ leptons arising from the decays $W^+ \to \tau^+\nu_\tau$ and $H^+ \to \tau^+\nu_\tau$ are predominantly left- and right-polarized, respectively, influencing the energy distributions of the decay products in the subsequent decays of the τ^+, we work out the effect of the radiative corrections on such distributions in the dominant (one-charged prong) decay channels $\tau^+ \to \pi^+\bar{\nu}_\tau, \rho^+\bar{\nu}_\tau, a_1^+\bar{\nu}_\tau$ and $\ell^+\nu_\ell\bar{\nu}_\tau$. The inclusive π^+ energy spectra in the decay chains $t \to b(W^+, H^+) \to b(\tau^+\nu_\tau) \to b(\pi^+\bar{\nu}_\tau\nu_\tau + X)$ are calculated, which can help in searching for the induced $H^\pm$ effects at the Tevatron and the LHC.
We also considered radiative corrections to the $K^\pm \to \pi^+\pi^-e^\pm\nu$ decay [5]. The final state interaction of pions in this decay allows one to obtain the value of the isospin and angular momentum zero pion-pion scattering length a_0^0. Basing on the lowest order results and the factorization hypothesis, we get the expressions for RC in the leading and next-to-leading logarithmical approximation. It is shown that the decay width dependence on the lepton mass m_e through the parameter $\sigma = \frac{\alpha}{2\pi}\left(\ln\frac{M^2}{m_e^2} - 1\right)$ has a standard form of the Drell–Yan process and is proportional to the Sommerfeld–Sakharov factor.
In paper [6], we showed the possibility to measure the deviation of the cross section of the small angle electron(positron)–ion elastic scattering from the Rutherford formula due to multiple virtual photon exchange in heavy ion-lepton collisions.

Production of mesons and jets at colliders A few papers [7, 8] were devoted to reactions of production of scalar and vector mesons in proton–proton and electron-positron collisions. In [7], the process of proton–proton collisions with the exchange of all types of forces (scalar, pseudoscalar, vector and axial vector) and with the production of neutral and charged scalar mesons $a_0(980)$, $a_+(980)$, $f_0(980)$, $\sigma(600)$ was considered. The estimation for the facilities of moderately high energies such as PANDA and NICA is presented. Similar analysis is given for processes of charged and neutral Higgs boson production at high energy proton-proton colliders such as Tevatron, RHIC and LHC. The possible signal of neutral Higgs boson decay into two oppositely charged leptons of different kinds is discussed.
In paper [8], we consider the vector mesons ϕ, ω, and J/Ψ with the $J^{PC} = 1^{--}$ production electron–positron collisions at high energy. The scattering and annihilation mechanisms are considered. The annihilation channel contribution is enhanced by the initial hard photon emission mechanism. It occurs in the kinematic region where the final particles are emitted

at large angles and plays the role of background. Differential distributions of the energy fraction and of the transverse component of the vector meson are calculated.
We also participated in performing calculations for the LHC scientific program. In particular, in paper [9], an eikonalized elastic proton-proton and proton-antiproton scattering amplitude $F(s,t)$, calculated from QCD as a finite sum of gluon ladders, is compared with the existing experimental data on the total cross section and the ratio $\rho(s,0) = ReF(s,0)/ImF(s,0)$ of the real part to the imaginary part of the forward amplitude.

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PRECISION SPECTROSCOPY OF LIGHT ATOMS AND MOLECULES

V.I. Korobov, J.-Ph. Karr, and L. Hilico

It is known that binding corrections of one, or two-loop diagram contributions are some series expansions in terms of α, fine structure constant. It is a big challenge to get higher order terms in these expansions. Say, in the case of the self-energy one-loop contribution, the $\ln^2\alpha$ and $\ln\alpha$ terms of an order of $m\alpha(Z\alpha)^6$ for the hydrogen atom S-state were known since the early 60s of the last century [1]. But the first accurate value for the nonlogarithmic term in this order was obtained only in 1992 [2]!

Recently, a general formula for the one-loop self-energy correction in order $m\alpha(Z\alpha)^6$ has been obtained [3]:

$$\begin{aligned}\Delta E^{(7)}_{\rm se} = \frac{\alpha}{\pi}\Bigg\{&(Z\alpha)^6\mathcal{L}+\left(\frac{5}{9}+\frac{2}{3}\ln\left[\frac{1}{2}(Z\alpha)^{-2}\right]\right)\left\langle\left[\boldsymbol{\nabla}^2V\right]\,Q(E-H)^{-1}Q\,H_R\right\rangle\\&+\frac{1}{2}\left\langle\sigma^{ij}\left[\nabla^iV\right]p^j\,Q(E-H)^{-1}Q\,H_R\right\rangle+\left(\frac{779}{14400}+\frac{11}{120}\ln\left[\frac{1}{2}(Z\alpha)^{-2}\right]\right)\left\langle\boldsymbol{\nabla}^4V\right\rangle\\&+\left(\frac{23}{576}+\frac{1}{24}\ln\left[\frac{1}{2}(Z\alpha)^{-2}\right]\right)\left\langle 2\mathrm{i}\sigma^{ij}p^i\left[\boldsymbol{\nabla}^2V\right]p^j\right\rangle\\&+\left(\frac{589}{720}+\frac{2}{3}\ln\left[\frac{1}{2}(Z\alpha)^{-2}\right]\right)\left\langle\left[\boldsymbol{\nabla}V\right]^2\right\rangle+\frac{3}{80}\left\langle\mathbf{p}^2\left[\boldsymbol{\nabla}^2V\right]\right\rangle-\frac{1}{2}\left\langle\mathbf{p}^2H_{so}\right\rangle\Bigg\}\end{aligned} \tag{1}$$

where $\mathcal{L}$ is the relativistic Bethe logarithm; Q is a projection operator and

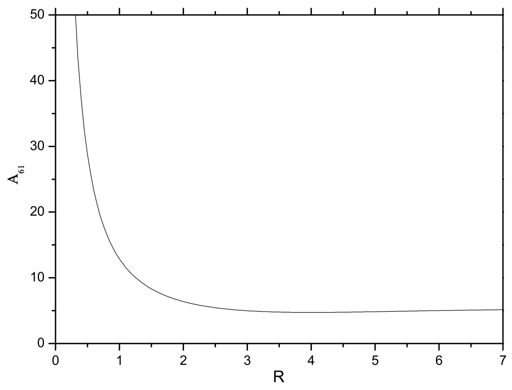
$$\begin{cases} H_{so} = \dfrac{1}{4}\,\sigma^{ij}\nabla^iVp^j = Z\alpha\dfrac{[\mathbf{r}\times\mathbf{p}]\,\boldsymbol{\sigma}}{4r^3}\,,\\ H_R = -\dfrac{p^4}{8}+\dfrac{\pi}{2}Z\alpha\delta(\mathbf{r})+H_{so}.\end{cases}$$

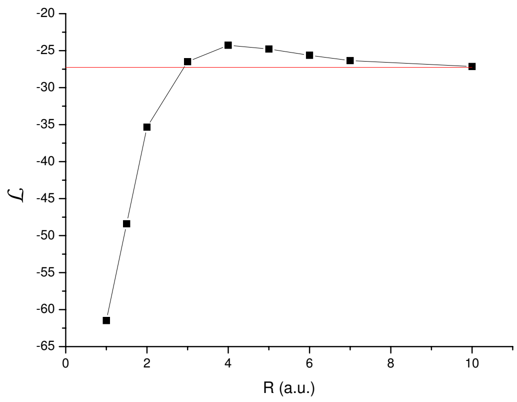
The major deficiency of this formula is that it is valid only for the normalized difference of S-states in hydrogen, $\Delta_n \equiv n^3\Delta E(nS)-\Delta E(1S)$. Thus, it is correct only to some additive contribution of the δ-function type.

In our recent studies we have found this additional δ-function type term. That would allow us to generalize application of Eq. (1) and to extend it to the two Coulomb center problem by introducing two finite value distributions:

$$\begin{aligned}\mathcal{Q} &= \lim_{r_0\to0}\left\{\left\langle\frac{1}{4\pi r^3}\right\rangle_{r_0}+(\ln r_0+\ln(Z\alpha)+\gamma_E)\left\langle\delta(\mathbf{r})\right\rangle\right\}\\ \mathcal{R} &= \lim_{r_0\to0}\left\{\left\langle\frac{1}{4\pi r^4}\right\rangle_{r_0}-\left[\frac{1}{r_0}\left\langle\delta(\mathbf{r})\right\rangle+(\ln r_0+\ln(Z\alpha)+\gamma_E)\left\langle\delta'(\mathbf{r})\right\rangle\right]\right\}.\end{aligned}$$

The first one is the well-known Araki-Sucher term [4]. The mean values for all operators in Eq. (1) can be then re-expressed in terms of these two distributions and finite value operators. For the two-center problem we elaborated the numerical way to obtain the "effective" potentials of the $\alpha(Z\alpha)^6$ order one-loop SE contribution for the $\ln\alpha$ term, $A_{61}(R)$, and the relativistic Bethe logarithm, $\mathcal{L}(R)$ [5]. The latter is an especially challenging numerical task even for the leading $m\alpha^5$ order. The results are shown in the figure below.





The numerical work is not yet fully completed, but eventually when it will be done, the relative uncertainty of theoretical prediction for ro-vibrational transition energies will be below 10^{-10} and would become an ultimate precision for determination of the electron atomic mass.
At present, the atomic mass deduced from the comparison of the antiprotonic helium theory [6] and two-photon spectroscopy experiment [7] is

$$A_r(e) = 0.000\,548\,579\,909\,1(7)[1.4\times10^{-9}].$$

That is still less accurate then the one obtained from the bound electron g factor measured in $^12C^{5+}$ ion, but it is about a factor of 2 better than from a direct comparison of the cyclotron frequencies in the Penning trap spectroscopy [8].

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Dynamics of charged particles produced by confining environment

S.I. Vinitsky

The adiabatic method for analysis of a channelling problem for charged particles with a transversal confining environment induced by an axial homogeneous magnetic field, a crystal lattice or a quantum well and wire was developed [1, 2, 3, 4, 5]. The proposed schemes were applied to analyze spectral characteristics of models of spheroidal quantum dots in the effective mass approximation [6]. Spectral and optical characteristics of hydrogen-like impurities in quantum-dimensional semiconductor nanostructures (quantum dots, quantum wires and quantum wells) were analyzed depending on their structure and external fields [7]. The effects of resonance transmission and total reflection for the Coulomb scattering with the confining environment due to interference of quasistationary states imbedded in the continuum were revealed. These effects can increase the rate of recombination processes of ions in a magneto-optical trap, or in the axis channeling in crystals and quantum wires [1, 8].

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