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# On Equations for the Multi-quark Bound States in the Framework of Mean-field Expansion in Nambu – Jona-Lasinio Model and NLO Meson Correction to Quark Mass

Jafarov Rauf Gadir Oglu
Institute for Physical Problems, Baku State University, Z. Khalilov, 23, AZ 1148 Baku, Azerbaijan
e-mail: jafarov@hotmail.com

***Abstract:***
*In present work we review our results of investigation of multi-particle equations in Nambu--Jona-Lasinio model in mean-field expansion. To formulate the mean-field expansion we have used an iteration scheme of solution of the Schwinger-Dyson equation with fermions bilocal source, which has been developed in works by V.E. Rochev et al. We have considered equations for Green functions of Nambu--Jona-Lasinio model up to third step for this iteration scheme, which essentially based on our previous results. The correction of pion amplitudes to a quark mass are calculated, also, in the Nambu-Jona-Lasinio model with 4-dimensional cutoff regularization and dimensional-analytically regularized Nambu-Jona-Lasinio model in mean-field expansion. The analytical calculation shows that in pole approximation of pion amplitude (which corresponds to the leading singularity of pseudo scalar amplitude) the pion contribution to quark mass exactly equals zero. But for the non-pole approximation of pseudo scalar amplitude the situation is different: value of numerical results differ from zero. Coincidence of results in both regularizations one can signify that the zero value of the pion correction to quark mass is the regularization-independent fact of Nambu-Jona-Lasinio model, in the limit of the leading-singularity approximation of amplitude.*

***Keywords***: *NJL model, multi-particle functions, chiral condensate*

## 1. INTRODUCTION AND A BRIEF REVIEW OF NONPERTURBATIV FIELD- THEORETICAL EQUATION

Ever since the success quantum electrodynamics by Feynman et al. ( [1] and references therein), corresponding field-theoretic formulations have been in the forefront of strong interaction dynamics at once, the main strategy being to device various closed form approaches which are represented as appropriate integral equation. One of the earliest efforts in this direction was Tamm-Dancoff equation [2] for 'mesodynamics', which was independent from each other is offered by Tamm and Dancoff, however ( also see [3]), earlier the similar method has been applied to electrodynamics by Fock [4]. The 3-dimensional Tamm-Dancoff equation and 4-dimensional Schwinger-Dayson equation (SDE) have been the source of much wisdom underlying the formulation of many approaches to strong interaction dynamics. To these one should add the Bethe-Salpeter equation (BSE) [5], which is an approximation to SDE for the dynamics of a 4-dimensional two-particle amplitude for the effective nucleon-nucleon interaction, but now adapted to the quark level. A major bottleneck for the BSE approach is its resistance to a probability interpretation, since the logical demands of its 4-dimensional content are incompatible with its approximate nature, which has led to many attempts at 3-dimensional reductions: Logunov-Tavkhelidze quasipotentials [6], Kadyshevsky formalism [7], and remark Georgian theorists et al. efforts [8], (for more review for 3-dimensional reductions see [9] and refs. therein, also).

The multi-particle (three or more particle) generalizations of the BSE have been also studied. A straightforward generalization of two-particle BSE has been intensively studied in sixties-seventies of last century. A best exposition of these studies can be found in the work of Huang and Weldon [10].

An essential imperfection of the original Bethe-Salpeter approach to multi-particle equations was a full disconnection of the approach with the field-theoretical equations for Green functions (which are known as SDE). This imperfection has been eliminated by Dahmen and Jona-Lasinio, which had included the BSE to the field-theoretical Lagrangian formalism with the consideration of functional Legendre transformation with respect to bilocal source of fields [11]. Then this approach has been generalized for multi-particle equations with consideration of Legendre transformation with respect to multi-local sources [12].

However, these theoretical constructions had not solved the principal dynamical problem of quantitative description of real bound states (nucleons, mesons etc.). A solution of BSE-type equations has been founded as a very complicated mathematical tool even for simple dynamical model. There is a main reason of a comparatively small popularity of the method of multi-particle BSE-type equations among the theorists. Much more popular approach to the problem of hadronization in QCD is based on the 't Hooft's conjecture that QCD can be regarded as an effective theory of mesons and glueballs [13]. Subsequently, it was shown by Witten that the baryons could be viewed as the solitons of the meson theory [14]. Futher development of these ideas has been successful and has leaded to the prediction of pentaquark states in baryon spectrum [15].

Nevertheless, the investigations of multi-quark equations are of significant interest due to the much less model assumptions in this approach in comparison with the chiral-soliton models. The solutions of multi-quark equations will provide us almost exhaustive information about the structure of hadrons.

We shall investigate Nambu--Jona-Lasinio (NJL) model with quark content which is one of the most successful effective models of QCD in the nonperturbative region (for review see [16], [17]). In overwhelming majority of the investigations, the NJL model has been considered in the mean-field approximation or in the leading order of $1/n_c$ $-$expansion. However, a number of perspective physical applications of NJL model is connected with multi-quark functions (for example: meson decays, pion-pion scattering, baryons, pentaquarks etc.). These multi-quark functions arise in higher orders of mean-field expansion (MFE) for NJL model. To formulate MFE we have used an iteration scheme of solution of SDE with fermion bilocal source [18].

We have considered equations for Green functions of NJL model up to third step of iterations. The leading approximation and first step of iteration maintain equations for the quark propagator and the two-quark function and also the next-to-leading order (NLO) correction to the quark propagator. The second step of this iterations maintains the equations for four-quark and three-quark functions, and third step of iterations maintains the equations for six-quark and five quark functions [19].

Since the mean-field approximation includes quark loops and the non-renormalizability of the NJL model implement to the indispensable the successful choice a regularization. Most common regularizations for NJL model traditionally entail a four-dimensional cutoff (FDC) regularization or a three-dimensional momentum cutoff regularization. Other regularization schemes also are used for NJL model [20].

In the framework of dimensional-analytically regularization (DAR) in MFE the scalar meson contributions in chiral quark condensate are calculated in [21], [22]. It was shown that sigma-meson contribution in chiral condensate for physical values of parameters is found to be significant and should be taken into account in the choice of the parameter values. Carry out improved fit of parameters of $SU_V(2)\times SU_A(2)$ symmetric NJL model.

In the work [22] has done a systematical comparison of the dimensional-analytically regularized NJL model with the NJL model with FDC regularization. Apart from the corrections to chiral condensate it was calculated also the corrections to quark mass in both regularizations. The numerical calculations at two characteristic values of condensate show that the pion contribution to the quark mass in both regularizations are equal to zero.

The present work, which essentially based on results [19], [23], is devoted to analytical calculations of pion correction to the quark mass.

## 2. MEAN-FIELD EXPANSION IN BILOCAL-SOURCE FORMALISM FOR NAMBU - JONA-LASINIO MODEL

We consider NJL model with the Lagrangian

$$L = \overline{\psi}\, i\bar{\partial}\psi + \frac{g}{2}\left[\left(\overline{\psi}\psi\right)^2 + \left(\overline{\psi}\, i\gamma_5\tau\psi\right)^2\right],$$

The Lagrangian is invariant under transformations of chiral group $SU_V(2)\times SU_A(2)$, which correspond to up and down quark sector.

A generating functional of Green functions (vacuum expectation values of $T-$products of fields) can be represented as the functional integral with bilocal source:

$$G(\eta) = \int D(\psi,\overline{\psi})\exp i\left\{\int dxL - \int dxdy\,\overline{\psi}(y)\eta(y,x)\psi(x)\right\}.$$

Here $\eta(y,x)$ is the bilocal source of the quark field.

The $n-$th functional derivative of $G$ over source is the $n-$particle ($2n-$point) Green function:

$$\left.\frac{\delta^n G}{\delta\eta(y_1,x_1)\cdots\delta\eta(y_n,x_n)}\right|_{\eta=0} = i^n < 0|T\{\psi(x_1)\overline{\psi}(y_1)\cdots\psi(x_n)\overline{\psi}(y_n)\}|0 > \equiv S_n\begin{pmatrix} x_1 & y_1 \\ \dots & \dots \\ x_n & y_n\end{pmatrix}$$

Translational invariance of the functional-integration measure gives us the functional-differential SDE for the generating functional $G$:

$$\delta(x-y)G + i\hat{\partial}_x\frac{\delta G}{\delta\eta(y,x)} + ig\left\{\frac{\delta}{\delta\eta(y,x)}tr\left[\frac{\delta G}{\delta\eta(x,x)}\right] - \gamma_5\tau\frac{\delta}{\delta\eta(y,x)}tr\left[\gamma_5\tau\frac{\delta G}{\delta\eta(x,x)}\right]\right\} =$$
$$= \int dx_1\eta(x,x_1)\frac{\delta G}{\delta\eta(y,x_1)} \quad . \qquad (1)$$

We shall solve this equation employing the method which proposed in work [21].

Functional $G^{(n)}$ is

$$G^{(n)} = P^{(n)}G^{(0)}$$

where $P^{(n)}$ is a polynomial of $2n-$th order over the bilocal source $\eta$.

### 2.1. LEADING ORDER AND FIRST STEP OF ITERATION: THE EQUATIONS FOR QUARK PROPAGATOR AND TWO-QUARK GREEN FUNCTION

A leading approximation is an approximation of the functional-differential SDE (1) without r.h.s. A solution of the leading approximation is the functional

$$G^{(0)} = \exp\{Tr(S*\eta)\}$$

Here and below $Tr$ is a trace in operator sense, $*$ multiplication operator. Function $S$ is a solution of the equation

$$\delta(x) + i\hat{\partial}S(x) + igS(x)tr[S(0)] = 0$$

The unique connected Green function of the leading approximation is the quark propagator $S$. Other connected Green functions appear in the following iteration steps. The quark propagator in the chiral limit is

$$S_1^{(0)} \equiv S = (m - \hat{p})^{-1},$$

where $m$ is the dynamical quark mass, which is a solution of gap equation in the chiral limit

$$m = -8igmn_c \int \frac{d\tilde{q}}{mi^2 - q^2}, \qquad (2)$$

where $d\tilde{q} = d^4q \big/ (2\pi)^4$.

The leading approximation generates the linear iteration scheme:

$$G = G^{(0)} + G^{(1)} + \cdots + G^{(n)} + \cdots,$$

where $n-$th step functional $G^{(n)}$ is a solution of the equation

$$G^{(n)} + i\bar{\partial}\frac{\delta G^{(n)}}{\delta\eta} + ig\left\{\frac{\delta}{\delta\eta}tr\left[\frac{\delta G^{(n)}}{\delta\eta}\right] - \gamma_5\tau\frac{\delta}{\delta\eta}tr\left[\gamma_5\tau\frac{\delta G^{(n)}}{\delta\eta}\right]\right\} = \eta * \bar{\partial}\frac{\delta G^{(n-1)}}{\delta\eta}. \qquad (3)$$

A solution of first-step equation (i.e. the Eq.(3) at $n=1$)
is functional

$$G^{(1)} = \left\{\frac{1}{2}Tr\left(S_2 * \eta^2\right) + Tr\left(S^{(1)} * \eta^2\right)\right\}G^{(0)}$$

The iteration-scheme equations give us the equation for two-particle function $S_2$:

$$S_2\begin{pmatrix} x & y \\ x' & y' \end{pmatrix} = -S(x-y')S(x'-y) + ig\int dx_1\left\{\left(S(x-x_1)S(x_1-y)\right)tr_u\left[S_2\begin{pmatrix} x_1 & x_1 \\ x' & y' \end{pmatrix}\right] - \right.$$

$$\left. -\left(S(x-x_1)\gamma_5\tau\, S(x_1-y)\right)tr_u\left[\gamma_5\tau\, S_2\begin{pmatrix} x_1 & x_1 \\ x' & y' \end{pmatrix}\right]\right\} \qquad (4)$$

(here $tr_u$ denotes the trace, which includes the upper line of function $S_2$) and the equation for NLO correction to quark propagator $S^{(1)}$

$$S^{(1)}(x-y) = ig\int dx_1 S(x-x_1)\left\{S_2\begin{pmatrix} x_1 & y \\ x_1 & x_1 \end{pmatrix} - \gamma_5\tau^a S_2\begin{pmatrix} x_1 & y \\ x_1 & x_1 \end{pmatrix}\gamma_5\tau^a\right\} +$$

$$+ ig\int dx_1 S(x-x_1)S(x_1-y)trS^{(1)}(0) \qquad (5)$$

The graphical representations of two-quark function see on Fig. 1, where the graphical notations of Fig. 2 are used.

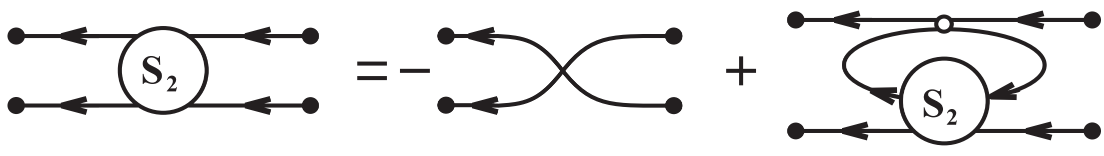


Figure 1. The equation for two-quark function

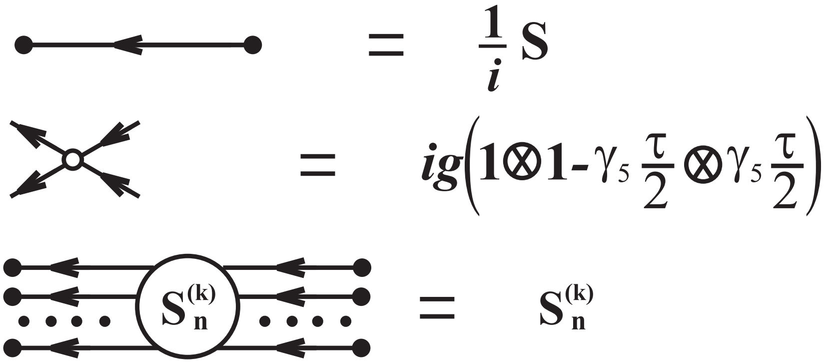


Figure 2. Diagram rules

These equations reduces in the momentum space to a system of simple algebraic forms. The NLO mass operator $\sum^{(1)} = S^{(-1)} * S^{(1)} * S^{(-1)}$ (where $S^{(1)}$ NLO quark propagator), is defined in $x-$space by equation:

$$\sum^{(1)}(x) = S(x)A_\sigma(x) + 3S(-x)A_\pi(x) + ig\delta(x) tr S^{(1)}(0) \; . \tag{6}$$

The Eqs. (4)-(5) can be easily solved, and the solutions contain singlet scalar quark-antiquark bound state with mass $2m$ (sigma-meson) and massless (in the chiral limit) pseudo scalar bound states (pion). To describe the solution of the equation for two-quark function and for future purposes we introduce the composite meson propagators by following way:

a) Let us define scalar-scalar function

$$S_\sigma(x-x') \equiv tr\left[ S_2\begin{pmatrix} x & x \\ x' & x' \end{pmatrix}\right] \sim <\bar{\psi}\psi(x)\bar{\psi}\psi(x')> . \tag{7}$$

From the Eq.(4) for two-quark function we obtain (in momentum space)

$$S_\sigma(p) = \frac{1}{ig}\left(1 - i\Delta_\sigma(p)\right) \; . \tag{8}$$

Here we define the following function, which we call sigma-meson propagator

$$\Delta_\sigma(p) = \frac{Z(p)}{4m^2 - p^2} \quad , \tag{9}$$

where

$$Z_\sigma(p) = \frac{I_0(4m^2)}{I_0(p^2)}$$

and

$$I_0(p) = \int d\,\bar{q}\, \frac{1}{\left(m^2 - (p+q)^2\right)\left(m^2 - q^2\right)} \quad . \tag{10}$$

The integral (10) is divergent, and it should be considered as a regularization.

b) Pseudoscalar-pseudoscalar function is defined as

$$S_\pi^{ab}(x-x') \equiv tr\left[ S_2\begin{pmatrix} x & x \\ x' & x' \end{pmatrix}\gamma_5 \frac{\tau^a}{2}\gamma_5 \frac{\tau^b}{2}\right] \sim <\bar{\psi}\gamma_5 \frac{\tau^a}{2}\psi(x)\bar{\psi}\gamma_5 \frac{\tau^b}{2}\psi(x')> \quad . \tag{11}$$

From the Eq.(4) for two-quark function we obtain (in momentum space):

$$S_\pi^{ab}(p) = -\frac{1}{ig}\left(\delta^{ab} - i\Delta_\pi^{ab}(p)\right). \tag{12}$$

Here we define the pion propagator

$$\Delta_{\pi}^{ab}(p) = -\frac{\delta^{ab} Z(p)}{p^2} \tag{13}$$

where $Z_{\pi}(p) = \frac{I_0(0)}{I_0(p^2)}$.

## 2.2. SECOND STEP: THE EQUATIONS FOR FOUR-QUARK AND THREE-QUARK GREEN FUNCTIONS

Second-step generating functional is

$$G^{(2)}[\eta] = \left\{ \frac{1}{4!} Tr\left(S_4 * \eta^4\right) + \frac{1}{3!} Tr\left(S_3 * \eta^3\right) + \frac{1}{2} Tr\left(S_2^{(1)} * \eta^2\right) + Tr\left(S^{(2)} * \eta\right) \right\} G^{(0)}.$$

The equations for four-quark and three-quark functions see on Figs. 3 and 4.

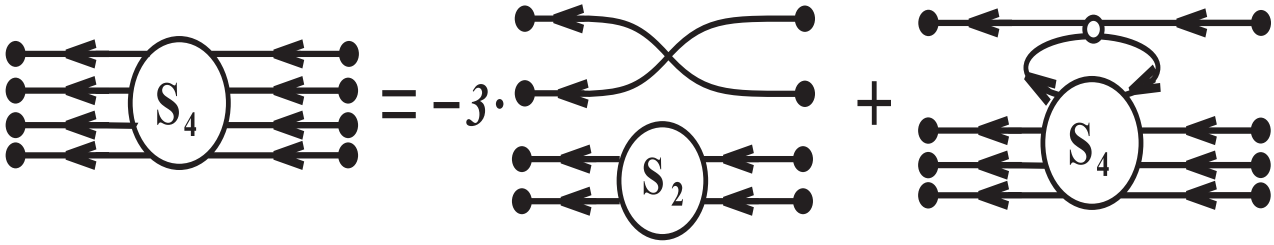


Figure 3. The equation for four-quark function

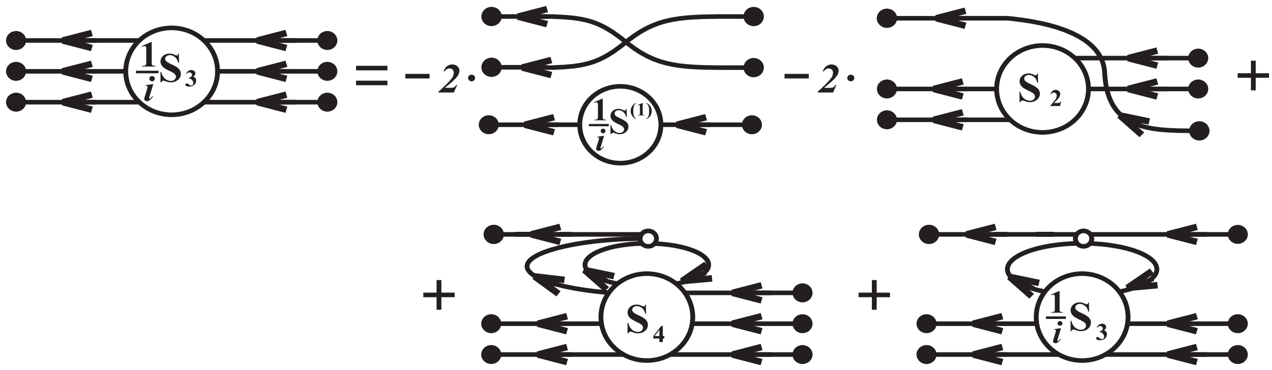


Figure 4. The equation for three-quark function

The equations for the four-quark function $S_4$ and for the three-quark functions $S_3$ are new. In the second step we obtain also the equations for NLO two-particle function $S_2^{(1)}$ and next-to-next-to-leading order (NNLO) correction to propagator $S^{(2)}$. These equations have the same form as the corresponding first step equation except of the inhomogeneous terms. The inhomogeneous term of NNLO quark propagator contains NLO two-quark function $S_2^{(1)}$ and the inhomogeneous term of equation for NLO two-quark function $S_2^{(1)}$ contains three-quark function $S_3$, also NLO quark propagator $S^{(1)}$ (it is naturally, i.e. the derivatives of first step functional to form the inhomogeneous term, also (see Eq. (3) ) ).

The equation for the four-quark function has a simple exact solution which is the product of two-quark functions (see Fig. 5). As is seen from this solution, the $\pi\pi-$scattering in NJL model is suppressed, i.e. in the second step of iterations this scattering is absent, and it perhaps arises in the third step of the iteration scheme only.

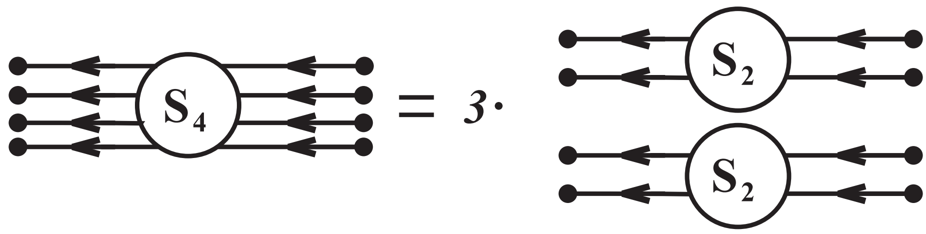


Figure 5. The solution of equation for four-quark function

**2.3. VERTEX** $\sigma\pi\pi$

The existence of the exact solution for the four-quark function gives us a possibility to obtain a closed equation for the three-quark function. As a first step in an investigation of this rather complicated equation we shall solve a problem of definition of $\sigma\pi\pi$ – vertex with composite sigma-meson and pions. Let us introduce a function

$$W^{ab}_{\sigma\pi\pi}(xx'x'') \equiv tr\left[S_3\begin{pmatrix} x & x \\ x' & x' \\ x'' & x'' \end{pmatrix}\gamma_5\frac{\tau^a}{2}\gamma_5\frac{\tau^b}{2}\right] \sim <\bar{\psi}\psi(x)\bar{\psi}\lambda_5\frac{\tau^a}{2}\psi(x')\bar{\psi}\lambda_5\frac{\tau^b}{2}\psi(x'')> ,$$

and define:

a) scalar vertex

$$V_\sigma(xx'x'') \equiv tr\left[S(x-x')S_2\begin{pmatrix} x' & x \\ x'' & x'' \end{pmatrix}\right] = 2in_c\int dx_1\upsilon_S(xx'x_1)\Delta_\sigma(x_1-x'') \, . \qquad (14)$$

Here: $\upsilon_S(xx'x'') = tr_a[S(x-x')S_0(x'-x'')S(x''-x)]$ is the triangle diagram.

b) pseudo scalar vertex

$$V^{ab}_\pi(xx'x'') \equiv tr\left[S(x-x')\gamma_5\frac{\tau^a}{2}S_2\begin{pmatrix} x' & x \\ x'' & x'' \end{pmatrix}\gamma_5\frac{\tau^b}{2}\right] = 2in_c\int dx_1\upsilon_p(xx'x_1)\Delta^{ab}_\pi(x_1-x''). \qquad (15)$$

Here:

$$\upsilon_p(xx'x'') = tr_a[S(x-x')\gamma_5 S(x'-x'')\gamma_5 S(x''-x)].$$

With definitions (7)-(15), for vertex function $W^{ab}$ we obtain the following equation:

$$W^{ab}_{\sigma\pi\pi}(xx'x'') = W^{ab}_0(xx'x'') + 2ign_c\int dx_1 l_S(x-x_1)W^{ab}_{\sigma\pi\pi}(x_1x'x''),$$

where $l_S(x) \equiv tr_a[S(x)S(-x)]$ is the scalar quark loop. Inhomogeneous term $W^{ab}_0$ is:

$$W^{ab}_0(xx'x'') = V^{ab}_\pi(xx'x'') + V^{ab}_\pi(xx''x') + 4ig\int dx_1V^{a_1a}_\pi(xx_1x')S^{a_1b}_\pi(x_1-x'') +$$
$$+ 4ig\int dx_1V^{a_1b}_\pi(xx_1x'')S^{a_1a}_\pi(x_1-x') - ig\int dx_1\left(V_\sigma(xx_1x_1) - 4V^{a_1a_1}_\pi(xx_1x_1)\right)S^{ab}_\pi(x'-x'')$$

Using definitions (7)-(15) we have:

$$\left[W^{ab}_0(xx'x'')\right]^{con} = -2n_c\int dx_1dx_2\upsilon_p(xx_1x_2)\left[\Delta^{a_1a}_\pi(x_2-x')\Delta^{a_1b}_\pi(x_1-x'') + \Delta^{a_1ab}_\pi(x_2-x'')\Delta^{a_2a}_\pi(x_1-x')\right]$$

The equation for $W^{ab}$ can be easy solved in the momentum space and the solution is

$$W^{ab}_{\sigma\pi\pi}(pp'p'') = i\Delta_\sigma(p)W^{ab}_0(pp'p'')$$

where $p$ is $\sigma -$ mesons momentum, and $p'$, $p''$ are pion momentum: $p = p' + p''$.

The connected part of $W^{ab}$ is the decay amplitude $\sigma \to \pi\pi$. It has the following form:

$$\left[W^{ab}_{\sigma\pi\pi}(pp'p'')\right]^{con} = \frac{2n_c}{i}\Delta_n(p)\left[\upsilon_p(pp'p'') + \upsilon_p(pp''p')\right]\Delta^{aa_1}_{\pi}(p')\Delta^{a_1 b}_{\pi}(p'')$$

(See also Fig.6.).

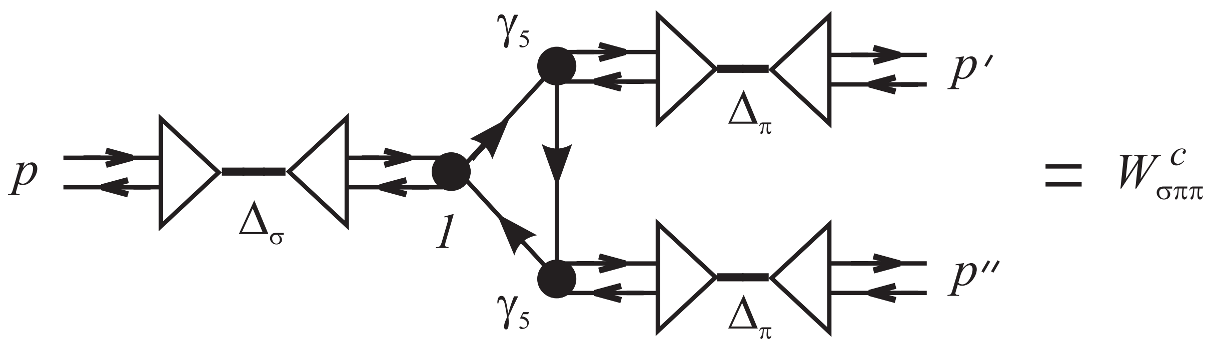


Figure 6. The connected part of sigma-pion-pion

**2.4. THIRD STEP: THE EQUATIONS FOR SIX-QUARK AND FIVE-QUARK GREEN FUNCTIONS**

The third-step generating functional is

$$G^{(3)}[\eta] = \left\{\frac{1}{6!}Tr\left(S_6 * \eta^6\right) + \frac{1}{5!}Tr\left(S_5 * \eta^5\right) + \frac{1}{4!}Tr\left(S^{(1)}_4 * \eta^4\right) + \frac{1}{3!}Tr\left(S^{(1)}_3 * \eta^3\right) + \right.$$

$$\left. + \frac{1}{2}Tr\left(S^{(2)}_2 * \eta^2\right) + Tr\left(S^{(3)} * \eta\right)\right\} G^{(0)}.$$

The graphical representations of equations for six-quark function and for five-quark function see on Figs. 7 and 8.

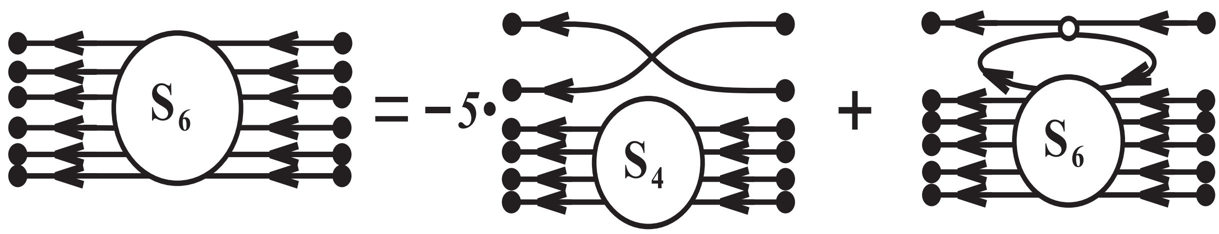


Figure 7. The equation for six-quark function

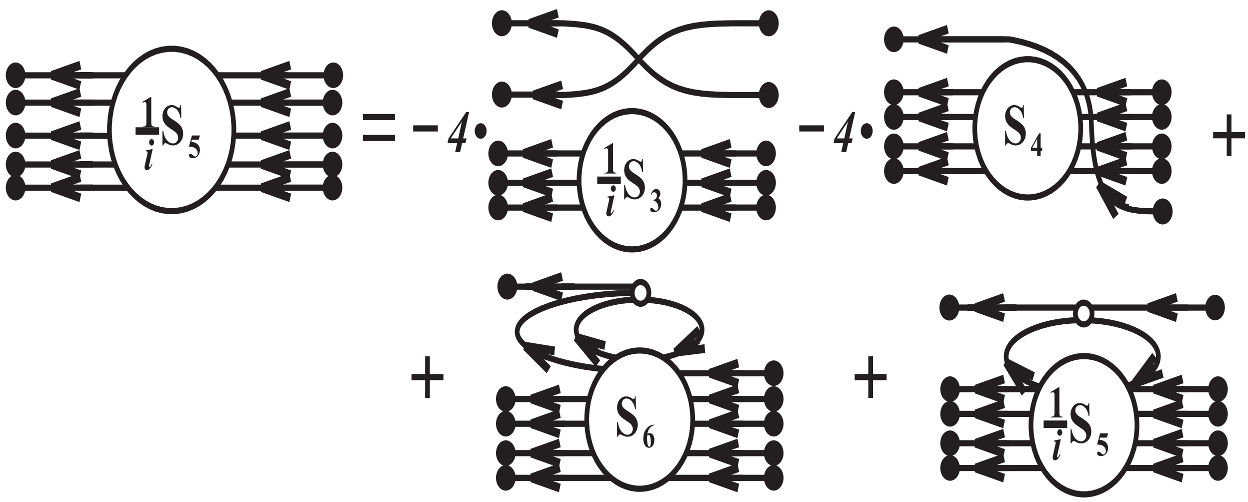


Figure 8. The equation for three-quark function

The equations for the six-quark function and for the five-quark function in our iteration scheme are new. The third step gives us the equations for NLO four-quark $S_4^{(1)}$ and three-quark $S_3^{(1)}$ functions, also NNLO two-quark function $S_2^{(2)}$ (which have the same form as the second-step equations except of the inhomogeneous term). The inhomogeneous term of the NLO four-quark function $S_4^{(1)}$ contains five-quark $S_5$ and $S_3$ functions, and NLO two-quark function $S_2^{(1)}$. Into inhomogeneous term of equation for NLO three-quark function $S_3^{(1)}$ enter: NLO four-quark function $S_4^{(1)}$, NLO two-quark function $S_2^{(1)}$ and NNLO one-particle function $S^{(2)}$. The equation for NNLO two-quark function $S_2^{(2)}$, which kernel is same form as the second- and first-steps equations, and it inhomogeneous term contains NLO 3-quark function $S_3^{(1)}$ and NNLO one-particle function $S^{(2)}$. The third step gives us also to the next-to-next-to-next-to-leading order (NNNLO) correction to quark propagator $S^{(3)}$. A inhomogeneous term of equation for NNNLO quark propagator contains the NNLO two-quark function $S_2^{(2)}$ (note that the inhomogeneous term of analogous equation for NNLO quark propagator $S^{(2)}$contain NLO two-quark function $S_2^{(1)}$and a inhomogeneous term of first-step equation for NLO quark propagator maintain two-quark Green function $S^{(2)}$), and all equations for quark propagators have the analogous form, except of the inhomogeneous term.

## 3. TWO-PARTICLE AMPLITUDE AND MESON CONTRIBUTIONS IN CHIRAL CONDENSATE AND NLO MASS FUNCTIONS

It follows from two-particle equation (4) leading order (LO) two-particle amplitude $A$ (connected part of amputated two-particle function) consist of two parts: pseudo scalar amplitude (pion) $A_\pi$ and scalar amplitude ($\sigma-$meson) $A_\sigma$. In momentum space these amplitudes of the NJL model depend on a momentum $p$ only, where $p$ is a sum of quark and antiquark momentum. They have the form (see [21], [22] for detail):

$$A_\pi = \frac{1}{4n_c I_0\left(p^2\right)p^2}, \qquad (16)$$

$$A_\sigma = \frac{1}{4n_c I_0\left(p^2\right)\left(4m^2 - p^2\right)}. \tag{17}$$

Here $I_0\left(p^2\right)$ is two-loop integral, which has the form (10), and it should be considered as a regularization.

The basic order parameter, which defines a degree of DCSB is a quantity $\chi$ of chiral condensate. Meson contribution to the chiral condensate can be calculated in the next-to-leading term of the MFE. A free quark propagator and gap equation (2) for NJL model give us to regularization-independent formula for LO condensate

$$\chi = itrS(0) = -\frac{m}{g}. \tag{18}$$

Let us defining NLO condensate:

$$\chi^{(1)} = itrS^{(1)}(0) \quad . \tag{19}$$

For the ratio

$$r = \frac{\chi^{(1)}}{\chi} \tag{20}$$

of NLO condensate (19) to the LO condensate (18) we obtain the formula [22]

$$r = -\frac{8ign_c}{1-8ign_c J}\int \frac{d\,\tilde{p}\;d\,\tilde{q}}{\left(m^2 - p^2\right)^2\left[m^2 - (p-q)^2\right]}\left[\left[m^2 - p^2 + 2(pq)\right]A_\sigma(q) - 3\left[m^2 - p^2 + 2(pq)\right]A_\pi(q)\right], \tag{21}$$

where

$$J = \int d\tilde{q}\,\frac{p^2 + m^2}{\left(m^2 - p^2\right)^2} \quad . \tag{22}$$

The integral (22) is divergent, and it should be considered as a regularization. It follows from equation (21), that the ratio $r$ of NLO condensate to the LO condensate consist of two parts: pion corrections (due to pseudo scalar amplitude $A_\pi$) and corrections due to scalar amplitude $A_\sigma$: $r = r_\sigma + r_\pi$.

### 3.1. NLO MASS FUNCTION

Eq.(5) for NLO quark propagator $S^{(1)}$ and Eq.(4) for two-amplitude $S_2$ give us a possibility for define the meson corrections to quark mass. First-order equations for iterations define corrections to quark propagator (see [22]). NLO mass operator $\Sigma^{(1)}$ has the form (6).

Let us to introducing dimensionless NLO mass functions $a^{(1)}$ and $b^{(1)}$:

$$\Sigma^{(1)} \equiv a^{(1)}\hat{p} - b^{(1)}m \tag{23}$$

and using the formula (20) for ratio of NLO condensate to the LO condensate in pion channel, we obtain from (6) the expressions for $a^{(1)}$ and $b^{(1)}$ in momentum space the following formulas:

$$p^2 a^{(1)}\left(p^2\right) = \int d\tilde{q}\,\frac{p^2 - (pq)}{m^2 - (p-q)^2}\left[A_\sigma(q) - 3A_\pi(q)\right], \tag{24}$$

$$b^{(1)}\left(p^2\right) = r - \int d\tilde{q}\,\frac{1}{m^2 - (p-q)^2}\left[A_\sigma(q) + 3A_\pi(q)\right]. \tag{25}$$

Using expression (23) for the NLO mass operator, we may to rewrite inverse quark propagator

$$S^{-1} = m - \hat{p} - \Sigma^{(1)} \tag{26}$$

as the form:

$$S^{-1} = \left(1 + b^{(1)}\right)m - \left(1 + a(1)\right)\hat{p} \tag{27}$$

(where, according gap equation (2), $m$ is LO quark mass).

Suppose the propagator has a pole in point $p^2 = m_r^2$, which corresponds to a particle with mass $m_r$. According Eq.(27) we obtain following equations for quark mass $m_r$:

$$b(m_r^2) = m_r a(m_r^2).$$

Since $a^{(1)}$and $b^{(1)}$ are small additions ($a^{(1)}$<<1, $b^{(1)}$<<1 ), we can to expand $a^{(1)}(m_r^2)$ and $b^{(1)}(m_r^2)$ near the point $m$ and to obtain the formula for the quark-mass correction $\delta m \cong m_r - m$:

$$\frac{\delta m}{m} \cong b^{(1)}\left(m^{(2)}\right) - a^{(1)}\left(m^{(2)}\right). \tag{28}$$

## 4. NJL MODEL WITH DIMENSIONAL-ANALYTICAL REGULARIZATION AND FOUR-DIMENSIONAL CUTOFF REGULARIZATION

The prediction of the model, however, are intimately compromising with the specific strategy adopted to handle the ultraviolet divergences given the nonrenormalizable nature of the model. Consequently, the quite essential point of the model is a regularization. Practitioners of the NJL model have followed the attitude of using it as a regularization-dependent model, considering the regularization procedure part of definition of the model, for example see [24].

Most common regularizations for NJL model traditionally entail a four-dimensional cutoff in Euclidean momentum or a three-dimensional momentum cutoff. Other regularization schemes (Pauli- Villars regularization or non-local Gauss form-factors) also are used for the NJL model. Even more often there are a works, which use a dimensional regularization in NJL model. In [25] the dimensional regularization is modified (which based on ideas of Wilson and Collins) to keep four dimensional properties of the nonrenormalizable theory as much as possible. To achieve this goal the dimensional regularization is applied to only the radial part in loop integrals. This is one of the analytic regularization. The meson loop contribution to the chiral symmetry breaking is also analyzed in the NJL model with the modified dimensional regularization [21]. In this treatment all calculations are made in four-dimensional Euclidean momentum space, and the regularization parameter is treated as a power of a weight function, which regularizes divergent integrals. It should be stressed that in this treatment of dimensional regularization, the regularization parameter is not at all a deviation in the physical dimension of space. In [26] it was studied characteristic features of the NJL model in the dimensional regularization. As usual the dimensional regularization is applied to momentum integrals for internal fermion lines. Since the model is not renormalizable, the authors ([26]) cannot take the four dimensional limit and they evaluated some physical properties of the model in the space-time dimensions less than four. The authors ([26]) take notice of that only the radiative corrections should be evaluated in the space-time dimensions less than four to keep the four-dimensional properties in the real world. In [22] the meson loop contribution to the chiral symmetry breaking is also analyzed in the NJL model with the modified dimensional regularization.

The calculating of the gap equation (2) in dimensional-analytical regularization (DAR) and four-dimensional cutoff (FDC) regularization lead to (in detail see [21], [22]):

$$1 = k\Gamma(\xi)\left(\frac{4\pi M^2}{m^2}\right)^{1+\xi}, \tag{29}$$

$$1 = k_\Lambda \left( 1 - \frac{1}{x} \log(1+x) \right), \tag{30}$$

respectively. Here, $k_\Lambda = \dfrac{gn_c \Lambda^2}{2\pi^2}$, $k = \dfrac{gn_c m^2}{2\pi^2}$, $x = \dfrac{\Lambda^2}{m^2}$, $\xi = \dfrac{2-D}{2}$ ($D-$dimension of space) and $\Lambda$ are regularizations parameters in DAR and in FDC regularization, respectively.

The calculating of the integral (10) in both regularizations (in *DAR* and in *FDC* regularization) lead us:

$$\left[ I_0\left(p^2\right) \right]^{DAR} = \frac{i}{(4\pi)^2} \frac{\xi}{k} F\left( 1+\xi, 1; 3/2; \frac{p^2}{4m^2} \right), \tag{31}$$

$$\begin{aligned} \left[ I_0\left(p^2\right) \right]^{FDC} = \frac{1}{(4\pi)^2} \Bigg[ & \log(1+x) - \frac{x}{1+x} F\left( 1,1; 3/2; \frac{p^2}{4m^2(1+x)} \right) - \\ & - \frac{p^2}{6m^2(1+x)} F\left( 1,1; 5/2; \frac{p^2}{4m^2(1+x)} \right) + \frac{p^2}{6m^2} F\left( 1,1; 5/2; \frac{p^2}{4m^2} \right) \Bigg] \end{aligned}. \tag{32}$$

Here $F(a,b;c;z)$ is the Gauss hypergeometric function.

## 5. MESON CONTRIBUTION TO QUARK MASS

### 5.1. MESON CONTRIBUTION TO QUARK MASS IN POLE APPROXIMATION OF SCALAR AND PSEUDOSCALAR AMPLITUDES

Pseudo scalar amplitude $A_\pi$ naturally is associated with the pion, which in the chiral limit is a massless Goldstone particle. In both regularizations under consideration we can define a pion propagator as a pole term of $A_\pi^{pole}$, which leads to the singularity of pseudo scalar amplitude (see [22]):

$$A_\pi^{pole} = \frac{1}{4n_c I_0(0) p^2}, \tag{33}$$

where $I_0(0)$ is defined by equation (31) for DAR and by (32) for FDC regularization and has the following forms

$$\left[ I_0(0) \right]^{DAR} = \frac{i}{(4\pi)^2} \frac{\xi}{k}, \tag{34}$$

$$\left[ I_0(0) \right]^{FDC} = \frac{i}{(4\pi)^2} \left[ \log(1+x) - \frac{x}{1+x} \right], \tag{35}$$

correspondingly.

For the scalar amplitude the situation is different. In both regularizations function $I_0(p)$ possesses a cut which originates in the point $p^2 = 4m^2$. Nevertheless, for FDC is possible to define a scalar sigma-meson propagator as [22]

$$A_\sigma^{pole} = \frac{1}{4n_c I_0\left(4m^2\right)\left(4m^2 - p^2\right)}, \tag{36}$$

since the integral $\left[I_0\left(4m^2\right)\right]^{FDC}$ (35) is a finite quantity. In a different way the matter is for DAR. Quantity $\left[I_0\left(4m^2\right)\right]^{DAR}$ is finite only at $\xi < -1/2$:

$$\left[I_0\left(4m^2\right)\right]^{DAR}\Big|_{\xi<-1/2} = -\frac{i}{8gn_c m^2}\frac{\xi}{1+2\xi}.$$

For the interpretation of the sigma-meson as a particle in the NJL model with DAR we can do the following trick: since in the region $-1 < \xi < -1/2$ integral $I_0$ converges we use the above value in the point $p^2 = 4m^2$ as a foundation an analytical continuation of the pole part of the amplitude on parameter $\xi$ the physical region $0 < \xi < 1$. Then the sigma-meson propagator for DAR will be [22]

$$\left[A_\sigma^{pole}(p)\right]^{DAR} = \frac{2igm^2(1+2\xi)}{\left(4m^2 - p^2\right)\xi}. \tag{37}$$

This expression was used for a calculation of the sigma-meson contribution in chiral condensate in work [21]. Surely, such procedure of definition of sigma-meson propagator seems to be a somewhat artificial. A more consistent procedure is a separation of a leading singular part of amplitude in the region of physical values of regularization parameter $\xi$ (see for more detail [21], [22]).

For the pseudo scalar amplitude the separation of leading singularity near point $p^2 = 0$ leads to same result (33), i.e. the pion in DAR possesses all properties of usual observable particle. For the scalar amplitude is not so. At $p^2 -> 4m^2$ in region $0 < \xi < 1$ [22]:

$$I_0^{DAR} \cong \frac{i\sqrt{\pi}\,\Gamma(\xi+1/2)}{16gn_c m^2\Gamma(\xi)}\cdot\left(\frac{4m^2}{4m^2-p^2}\right)^{\xi+1/2}$$

and, correspondingly, the leading singularity, i.e. leading term in an expansion on powers of $4m^2 - p^2$ is the expression

$$\left[A_\sigma^{LS}\right]^{DAR} \cong -\frac{ig\Gamma(\xi)}{\sqrt{\pi}\,\Gamma(\xi+1/2)}\cdot\left(\frac{4m^2}{4m^2-p^2}\right)^{1/2-\xi}. \tag{38}$$

## 5.1.1. PION CONTRIBUTION TO QUARK MASS

According the Eqs. (33)-(35) the expressions for pion amplitudes in pole approximation in both regularizations (DAR and FDC regularization) has the forms

$$\left(A_\pi^{pole}\right)^{DAR} = \frac{1}{12p^2 I_0^{DAR}(0)} = -\frac{2igm^2}{\xi p^2}\quad, \tag{39}$$

$$\left(A_\pi^{pole}\right)^{FDC} = \frac{1}{12p^2 I_0^{FDC}(0)} = -i\frac{4\pi^2}{3\left(\log(1+x) - \frac{x}{1+x}\right)p^2}. \tag{40}$$

The pion contribution NLO condensate in pole approximation of pion amplitude in both regularizations (DAR Eq.(39) and FDC regularization Eq.(40)) is calculated by Eqs. (21)-(22) in pion channel

$$r = \frac{24ign_c}{1-8ign_c J}\int d\tilde{p}d\tilde{q}\,\frac{\left[m^2 - p^2 + 2(pq)\right]A_\pi(q)}{\left(m^2-p^2\right)^2\left[m^2-(p-q)^2\right]}.$$

where $J$ has the form (22). All integrals over $dp$ and $dq$ can be calculated in closed form, and the results in both regularizations are the very simple expressions [21], [22]:

$$r_{\pi}^{DAR} = \frac{1}{8\xi}, \tag{41}$$

$$r_{\pi}^{FDC} = -\frac{\log(1+x)}{8\left(\log(1+x) - \frac{x}{1+x}\right)} \quad .$$

(42)

According the Eqs. (24)-(25) the NLO mass functions $a^{(1)}$ and $b^{(1)}$ in pion channel are defined by the following equations:

$$p^2 a_{\pi}^{(1)} = \left(p^2\right) = -3\int \frac{d\tilde{q}}{m^2 - (p-q)^2} A_{\pi}^{pole}(q) \ , \tag{43}$$

$$b_{\pi}^{(1)}\left(p^2\right) = r_{\pi} - 3\int \frac{d\tilde{q}}{m^2 - (p-q)^2} A_{\pi}^{pole}(q) \ . \tag{44}$$

Using the leading singularity approximation for $\left(A_{\pi}^{pole}\right)^{DAR}$ (39) and $\left(A_{\pi}^{pole}\right)^{FDC}$ (40) in (43) and (44) after calculating the integrals in DAR and FDC regularization we obtain for the pion corrections to the quark mass in next expressions according to (28)

$$\left(\frac{\delta m_{\pi}}{m}\right)^{DAR} = r_{\pi}^{DAR} - \frac{1}{8\xi}, \tag{45}$$

$$\left(\frac{\delta m_{\pi}}{m}\right)^{FDC} = r_{\pi}^{FDC} + \frac{\log(1+x)}{8\left(\log(1+x) - \frac{x}{1+x}\right)}. \tag{46}$$

From (45) and (46) it follows that, the pion contribution in quark mass is equal zero, according to (41) and (42).

### 5.1.2. SIGMA-MESON CONTRIBUTION TO QUARK MASS

Consider a contribution of scalar amplitude in pole approximation $A_{\sigma}^{pole}$ in quark mass. In correspondence with Eqs.(24) and (25) we have

$$p^2 a_{\sigma}^{(1)}\left(p^2\right) = \int d\tilde{q} \frac{p^2 - (pq)}{m^2 - (p-q)^2} A_{\sigma}(q), \tag{47}$$

$$b_{\sigma}^{(1)}\left(p^2\right) = r_{\sigma} - \int d\tilde{q} \frac{1}{m^2 - (p-q)^2} A_{\sigma}(q).$$

(48)

To calculate this contributions we use the leading-singularity approximation for amplitudes:

$$A_{\sigma}^{pole} = \frac{1}{4n_c I_0\left(4m^2\right)\left(4m^2 - p^2\right)\Big|_{p-4m^2}}, \tag{49}$$

with

$$\left[I_0\left(4m^2\right)\right]^{FDC} = \frac{i}{(4\pi)^2}\left[\log(1+x) - \sqrt{x}\arctan\frac{1}{\sqrt{x}}\right] -$$

according Eq. (32) in $p^2 = 4m^2$, for FDG regularization and Eq. (38) for DAR.

From Eqs. (21)-(22) we obtain the quantity $(r_\sigma)^{DAR}$ in DAR. A computation gives us the following values for sigma-meson contribution: $\xi = 0{,}25$ we obtain $(r_\sigma)^{DAR} = -0{,}033$; at $\xi = 0{,}4$ we obtain $(r_\sigma)^{DAR} = -0{,}01$. The sigma-meson contribution is small in comparison of the contribution and possesses the opposite sign, i.e. it decrease the common contribution (see [22], note, that this result is qualitatively the same as result of work [21], in which was used a pole approximation $A_\sigma$. Thus, all conclusions of work [21] about the part of the meson contributions stand also for the more exact leading-singularity approximation, which is used in [22], or in present work).

For FDC regularization the leading-singularity approximation for $A_\sigma$ coincides with the pole approximation (49). The ratio in sigma-meson channel $r_\sigma$ calculated by Eqs. (21)-(22). This quantity for FDC is a function of $x \equiv \Lambda^2 / m^2$. The values of $r_\sigma^{FDC}(x)$ for two characteristic values of ratio: at $x = 3$ (which corresponds to value $c^{(0)} = -210 MeV$ of the LO condensate) - $r_\sigma^{FDC}(3) = -0{,}007$. At $x = 19$, which corresponds to value $c^{(0)} = -250 MeV$ of the LO condensate, the ratio is: $r_\sigma^{FDC}(19) = -0{,}116$. In contrast to the DAR, the sign of sigma contribution for FDC is the same as for pion contribution.

A sigma-correction to quark mass for DAR given by formula [22]

$$\left(\frac{\delta m_\sigma}{m}\right)^{DAR} = r_\sigma^{DAR} - \frac{\cos \pi\xi}{4^{1+\xi} n_c \pi}(1/2 - \xi)$$

and attains: at $\xi = 0{,}25$: $\delta m_\sigma^{DAR} = -0{,}086m$, at $\xi = 0{,}4$: $\delta m_\sigma^{DAR} = -0{,}056m$. Since a pion correction to quark mass in this regularization equal zero (see above), these values are full corrections to quark mass for DAR [22].

$$\left(\frac{\delta m_\sigma}{m}\right)^{FDC} = r_\sigma^{FDC} - \frac{4\log(1 + x/4) - \log(1 + x)}{8n_c\left[\log(1 + x) + \sqrt{x}\arctan\sqrt{1/x}\right]} .$$

At $x = 3$: $\delta m_\sigma^{FDC} = -0.022m$; ; at $x = 19$: $\delta m_\sigma^{FDC} = -0.158m$ [22].

### 5.2. PION CORRECTION TO QUARK MASS IN NON-POLE APPROXIMATION OF AMPLITUDE

In sub subsection 5.1.1, using the pole approximation for the pseudo scalar amplitude in both regularization ( in FDC regularization and DAR) we obtain the values for $\pi-$meson contribution in quark mass equal to zero. However, since the model is not renormalizable in four space-time dimensions, the physical results and parameters depend on the regularization method. This lead us to calculate the correction to quark mass beyond the non-pole approach of the amplitude. Using the expressions of pion amplitude (16) and the integral (31) in DAR, we can to calculate the ratio in pion sector $r_\pi$. Also, having calculated in DAR NLO mass functions by the Eqs. (24) and (25), according the formula (28) of NLO quark mass correction we obtain:

$$\left[\left(\frac{\delta m_\pi}{m}\right)^{DAR}\right]^{non-pole} = \left[r_\pi^{DAR}\right]^{non-pole} - \frac{1}{2\sqrt{\pi}\Gamma(2-\xi)}\sum_{k=0}^{\infty} 4^k\,\Gamma(1-\xi-k)\Gamma(3/2-k)F(1+\xi+k,\,1+k;\,2-\xi;\,1)$$

where

$$\left[r_\pi^{DAR}\right]^{non-pole} = -2\frac{\sin(\pi\xi)}{\pi\xi}\int_0^\infty dz\,\frac{z^{-1-\xi}}{F(1+\xi,1;3/2;-z/4)}\int_0^1 du\,\frac{1-u}{[1+u(1-u)z]^{1+\xi}} \times \left[1 - \xi + (1+\xi)\frac{1+u(u-2)z}{1+u(1-u)z}\right]$$

From here is clears that the result differ from zero.

This means, that the zero value of the pion correction to quark mass is independent from regularization choice in NJL model in leading singularity approach of pseudo scalar amplitude.

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