Calculating the Decay Constants in Dynamical Symmetry Breaking

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Abstract

Physical quantities accompanying the dynamical chiral symmetry breaking in QCD like theories are examined. Expression for the decay constant (usually called f_π) is given exactly to the ladder approximation by solving the Bethe-Salpeter equation. Numerical calculation is done for Σ(self energy function of fermion), f_π, $\langle\bar{\psi}\psi\rangle_{\mathrm{R}}$, with various β-function of the gauge coupling constant. We find that $f_\pi/\langle\bar{\psi}\psi\rangle_{\mathrm{R}}$, $f_\pi/\Lambda_{\mathrm{QCD}}$ are rather stable against the infrared cutoff parameter, while Σ-function itself depends strongly on the infrared cutoff. In this talk, I will describe very basic points of our framework and present some numerical results.

The spontaneous breaking of chiral symmetry plays essential role in many aspects of elementary particle physics. We have learned much in the strong interaction regime, that is, asymptotically free non-Abelian gauge theory (QCD) breaks chiral symmetry spontaneously. The electroweak theory must spontaneously break some sort of chiral symmetry, and there are various approaches where this symmetry breaking is assumed to occur in a dynamical way like QCD via another strong interactions. Furthermore, unified gauge theories of elementary particles require more fundamental understanding of the chiral symmetry breaking. Here we describe our recent work of calculating the decay constant exactly to the ladder approximation without additional assumptions.[1)]

Basic Framework

Our framework here is summarized as follows. We take a non-trivial solution of the Schwinger-Dyson equation so as to define a vacuum which spontaneously breaks the symmetry. To the ladder approximation in the Landau gauge, the Schwinger-Dyson equation is written as[2)]

$$\Sigma(x) = \frac{\lambda(x)}{4x}\int_0^x \frac{y\Sigma(y)dy}{y+\Sigma^2(y)} + \int_x^{\Lambda^2} \frac{\lambda(y)\Sigma(y)dy}{4(y+\Sigma^2(y))} \quad , \tag{1}$$

where $\Sigma(x)$-function is the self-energy function of fermion, and x is the momentum squared. This is defined by a ratio, $B(x)/A(x)$, where the fermion two point function is defined by

$$\langle T\psi\bar{\psi}\rangle = \frac{i}{\gamma_\mu p^\mu A(x) - B(x)} \quad . \tag{2}$$

The running coupling constant $\lambda(x)$ is normalized as

$$\lambda(x) = \frac{3}{4\pi^2} C_2(F) g^2(x) \quad , \tag{3}$$

where $C_2(F)$ is the second order Casimir of the fermion representation and $g(x)$ is the standard gauge coupling constant.

In the leading logarithmic approximation, λ(x) can be parametrized as

$$\lambda(x) = \frac{\lambda_0}{1+At} \quad , \qquad t = \ln x/\mu^2 \quad . \tag{4}$$

Here we define a parameter B,

$$B \equiv \frac{A}{\lambda_0} = \frac{\beta_0}{12C_2(F)} \quad , \tag{5}$$

where β_0 is the lowest order coefficient of the β-function. For example, in QCD with three color triplet quarks, B is 9/16. With this parameter B we can deal with various QCD-like theories in a unified way, even including the so-called fixed coupling theories as a limit of vanishing B.[3]

Inclusion of 'running' of the coupling constant is some sort of 'improvement' of the original ladder approximation, although the validity is not clear. Furthermore there is some freedom of how we define the renormalization point of the coupling constant in the Schwinger-Dyson kernel. The above method[2] is known to make high energy behavior of Σ(x$\rightarrow\infty$) consistent with the leading renormalization group analysis, which is reasonable if one notes that the leading logarithmic approximation sums up all ladder-type diagrams. Actually, the integral equation (1) is equivalent to the following differential equation with appropriate boundary conditions,

$$\left[\frac{\Sigma'(x)}{(\lambda(x)/4x)'}\right]' = \frac{x\Sigma(x)}{x+\Sigma^2(x)} \quad . \tag{6}$$

The asymptotic behavior of the solution takes a form,

$$\Sigma(x\rightarrow\infty) \sim \frac{1}{x}\left(\ln\frac{x}{\mu^2}\right)^{\frac{1}{4B}-1} \quad , \tag{7}$$

and it is consistent with the operator product expansion result ;[4]

$$\Sigma^{\mathrm{OPE}}(x\rightarrow\infty) \sim \frac{\lambda(x)}{x}\left(\ln\frac{x}{\mu^2}\right)^{\frac{1}{4B}} \quad . \tag{8}$$

The Schwinger-Dyson equation is equivalent to the stationary condition of the effective potential for the bilocal operator $T\bar{\psi}(x)\psi(0)$, and its solution $\Sigma(x)$ (a function) defines a vacuum. In the auxiliary filed method, it is regarded as condensation of a bilocal auxiliary field representing composite states.[5-7]

On the vacuum, there are massless modes, the Nambu-Goldstone modes. These modes are excited by applying the bilinear operator $\bar{\psi}\psi$ and are considered to be composite states of the fermion and the anti-fermion. The wave-function of a composite state

$$\chi(p+q/2, p-q/2) \sim \langle 0|T\psi(x)\bar{\psi}(y)|\mathrm{Ps}\rangle \quad , \tag{9}$$

satisfies the Bethe-Salpeter equation. When solving this equation, information on the vacuum is included via $\Sigma(x)$ function in the equation. The Bethe-Salpeter equation is a linear integral equation with inhomogeneous terms supplied by $\Sigma(x)$, and there must be a unique solution for function χ, otherwise there are infinitely many massless modes in this channel on the vacuum.

Physical Quantities

The spontaneous chiral symmetry breaking is argued with the following physical parameters.

1. β-function of the gauge coupling constant. This is determined by the gauge group and matter contents of the theory. β-function is the only parameter which discriminates the dynamics of theories and controls the running of the coupling constant.

2. In the literature, $\Sigma(p=0)$ is often referred to as a spontaneously generated mass parameter. Non-vanishing Σ-function definitely indicates the spontaneous breaking, and so it can be used as an order parameter of the symmetry breakdown. However, there is no way of directly observing $\Sigma(0)$.

3. f_π, the decay constant. We generically call the Nambu-Goldstone bosons as π-on and their decay constants as f_π. f_π is directly related to physical quantities. When the symmetry is gauged, f_π determines the mass of gauge bosons,

$$M = gf_\pi/2 \quad . \tag{10}$$

The decay constant is defined as

$$\bar{\psi}\gamma_\mu\gamma_5\psi(x) = f_\pi\partial_\mu\pi(x) + \cdots \quad , \tag{11}$$

where $\pi(x)$ is the renormalized π-on fields. By sandwiching the above equation with the vacuum and the π-on state, we have

$$f_\pi q_\mu = -\int \frac{d^4p}{i(2\pi)^4}\mathrm{Tr}\left[\gamma_\mu\gamma_5\chi(p+q/2, p-q/2)\right] \quad . \tag{12}$$

Thus we can evaluate f_π exactly from the Bethe-Salpeter solution χ.

4. $\langle\bar{\psi}\psi\rangle$. This vacuum expectation value is the most straightforward order parameter of the symmetry. When one consider a small explicit breaking of the initial chiral symmetry, the Nambu-Goldstone boson gets a mass,

$$m_\pi^2 \propto m(\mu)_{\text{breaking}}\left\langle\bar{\psi}\psi\right\rangle_\mu / f_\pi^2 \quad , \tag{13}$$

to the first order approximation. Thus, $\langle\bar{\psi}\psi\rangle$ is a directly observable quantity. It also determines the size of $\Sigma(x)$ at high momentum, which is evaluated with the renormalization group and the operator product expansion technique. $\langle\bar{\psi}\psi\rangle$ is an amplitude of a composite operator, and it needs an additional renormalization, and thus the above formula contains a renormalization point μ. Actually the μ-dependence cancels out each other between the explicit breaking mass parameter and the composite operator. When we evaluate $\langle\bar{\psi}\psi\rangle$ independently, we have to specify a renormalization point μ for it.

There are basically two equivalent ways of evaluating $\langle\bar{\psi}\psi\rangle$ from a solution $\Sigma(x)$ of the Schwinger-Dyson equation. One is that fitting the high energy behavior of $\Sigma(x \to \infty)$ to the operator product expansion result, and extract the proportionality coefficient, which gives $\langle\bar{\psi}\psi\rangle$. This method is not appropriate in the numerical calculation with effective 'explicit breaking' mass terms, where the 'explicit breaking' terms takes the leading asymptotic behavior overwhelming the $\langle\bar{\psi}\psi\rangle$ dependent term.

The other way, which we take here, is to directly evaluate the matrix element $\langle\bar{\psi}\psi\rangle$,

$$\langle\bar{\psi}\psi\rangle = \frac{1}{16\pi^2}\int_0^{\Lambda^2} dx \frac{x\Sigma(x)}{x+\Sigma^2(x)} \quad , \tag{14}$$

and take the value as $\langle\bar{\psi}\psi\rangle$ renormalized at μ=Λ. Then renormalize it by the leading renormalization group formula,

$$\langle\bar{\psi}\psi\rangle_\mu = \langle\bar{\psi}\psi\rangle_\Lambda \left[\frac{\lambda(\Lambda)}{\lambda(\mu)}\right]^{\frac{1}{4B}} \quad . \tag{15}$$

All these renormalization procedures for $\langle\bar{\psi}\psi\rangle$ are totally consistent due to Eq.(7)

Parameters in Theory

We will investigate the above 'physical' quantities in an asymptotically free gauge theory, where the β-function determines the dynamics. At the infrared, the coupling constant diverges. Physically, due to the confinement, such divergence is regarded as irrelevant. To take into account of this, we need some 'infrared' cutoff when calculating quantum loops with the effective coupling constant. There is no definite way of determining this infrared cutoff. However we may claim that if an approximation method is to be reliable, the physical results should depend little on the way of the infrared cutoff. Hence, we rather use the infrared cutoff dependence of various parameters to examine which quantity is 'good', or is properly evaluated in the specific approximation.

Characteristics of chiral symmetry breaking depend on the representation of fermion as well as the β-function. We expect that this dependence of fermion representation may explain hierarchies in various interactions and masses to some extent. We will adopt some different representations for the fermion.[8)]

The decay constant f_π is the central object of chiral symmetry breaking. So far, an approximated formula relying on the so-called dynamical perturbation theory[9)] has been used. Some variations have been proposed for SU(2) breaking case (see below).[10-12)] However, to the ladder approximation (the leading approximation in the large-N like expansion), f_π can be exactly evaluated by solving the Bethe-Salpeter equation to this order without *any* additional approximations.[1)]

Symmetry structures

In QCD, chiral symmetry breaking occurs like (take SU(2) for simplicity),

$$\mathrm{SU(2)_L \times SU(2)_R \to SU(2)_{L+R}} \quad ,$$

that is, conserving the 'vector' part. In this case, we have a common f_π for all triplet Nambu-Goldstone particles.

On the other hand, in models for dynamical electroweak symmetry breaking,[10,13,14)] the symmetry structure takes the following form,

$$\mathrm{SU(2)_L \times U(1)_Y \to U(1)_{QED}} \quad .$$

The corresponding Nambu-Goldstone particles form two irreducible representations of the remaining symmetry group, $\chi^{\pm}, \chi^0$. Therefore there are two independent decay constants $f_\pi^{\pm}, f_\pi^0$, in this pattern of symmetry breaking. This fact that there are two independent mass scales, $f_\pi^{\pm}, f_\pi^0$, is the very reason why we have the famous parameter ρ, which is a ratio of these two. As is well established experimentally, ρ is very close to unity.

In the standard Weinberg-Salam model of the electroweak theory (with an SU(2) doublet Higgs condensation), ρ is exactly unity to the tree approximation. Hence, *if the doublet Higgs scheme is true*, the quantum corrections to ρ mainly due to the heaviest quark loop is severely limited, and it gives an *upper bound* of the top quark mass; some 200 GeV typically. Note that this bound comes out because we assume the electroweak theory with ρ=1 (at the tree), without allowing ρ to be an additional free parameter. Up to the analyses with low energy parameters averaged among various processes, the upper limit of the top quark mass seems to disappear, if we add another Higgs structure to cancel out the loop effects. However, more detailed analyses on process by process may not allow a large top quark mass even with a free ρ_{tree}. This type of free ρ analyses are to be done more elaborately.

Recent models for the electroweak symmetry breaking with four fermion interactions, where the Higgs mode is supplied by a bound state of the top anti-top, give effectively almost the same low energy structure as that of the doublet Higgs standard model. In those models, however, the ρ=1 constraint should be argued carefully, since it is not easy to correctly evaluate the quantum corrections to it. Thus, it is quite important to evaluate $f_\pi^{\pm}, f_\pi^0$ exactly to the ladder approximation.

Formula for f_π

Here we describe only the case of QCD-like theories. The case with four-fermion interactions generating SU(2) breaking goes similarly, and will be published elsewhere.

According to the previous discussion, f_π is evaluated, *by definition*, as Eq.(12). We define invariant amplitudes for the Bethe-Salpeter wave-function as follows:

$$\begin{aligned}\chi(p+q/2,p-q/2) =&\gamma_5 S(p,q) + \gamma_\mu\gamma_5(P(p,q)(p\cdot q)p^\mu + Q(p,q)q^\mu) \\ &+ \sigma_{\mu\nu}\gamma_5 T(p,q)(q^\mu p^\nu - p^\mu q^\nu) \ . \end{aligned} \tag{16}$$

We also define the same form of invariant functions for the amputated Bethe-Salpeter wave-function:

$$\hat{\chi}(p,q) \equiv (\not{p} - \Sigma(p^2))\chi(p,q)(\not{q} - \Sigma(q^2)) \ , \tag{17}$$

$$\begin{aligned}\hat{\chi}(p+q/2,p-q/2) =&\gamma_5 \hat{S}(p,q) + \gamma_\mu\gamma_5(\hat{P}(p,q)(p\cdot q)p^\mu + \hat{Q}(p,q)q^\mu) \\ &+ \sigma_{\mu\nu}\gamma_5 \hat{T}(p,q)(q^\mu p^\nu - p^\mu q^\nu) \ . \end{aligned} \tag{18}$$

All invariant functions can be taken as even-functions of $(p\cdot q)$. By substituting the above expression into the definition of f_π (12), we get

$$f_\pi = \frac{N}{2}\int_0^\infty \frac{x dx}{16\pi^2}(4Q(x) - xP(x)) \ , \tag{19}$$

where factor N comes from SU(N) of the gauge interactions.

The Bethe-Salpeter equation is a homogeneous equation for χ, and we needs an additional normalization condition. Taking a limit of $q \to 0$ in the Ward-Takahashi identity,

$$-q^\mu \Gamma_{5\mu}(p-q/2,p+q/2) = S_F^{-1}(p-q/2)\gamma_5 + \gamma_5 S_F^{-1}(p+q/2) \ , \tag{20}$$

we have the normalization condition,

$$f_\pi \hat{S}(x) = 2\Sigma(x) \ . \tag{21}$$

For simplicity of equations, we rescale the Bethe-Salpeter equation with a factor of $f_\pi/2$, that is, we normalize $\hat{S}(x)$ to be $\Sigma(x)$. Under this normalization condition, we finally get the *ladder exact formula for* f_π:

$$\frac{1}{2}f_\pi^2 = \frac{N}{2}\int_0^\infty \frac{x dx}{16\pi^2}(4Q(x) - xP(x)) \ , \tag{22}$$

and the ladder Bethe-Salpeter equation as follows:

$$\begin{pmatrix}\hat{Q}(x)\\ \hat{P}(x)\end{pmatrix} = -\lambda \begin{pmatrix} \int_0^x dy \frac{y(2x+y)}{2x^2} + \int_x^{\Lambda^2} dy \frac{3}{2} & \int_0^y dy \frac{-y^2(3x+y)}{12x^2} + \int_x^{\Lambda^2} dy \frac{5x-9y}{12} \\ \int_0^x dy \frac{2y(y-x)}{x^3} & \int_0^x dy \frac{y^2(3x-2y)}{6x^3} + \int_x^{\Lambda^2} dy \frac{1}{6} \end{pmatrix} \begin{pmatrix} Q(y) \\ P(y)\end{pmatrix} \tag{23}$$

Definition (17) gives relations between invariant amplitudes,

$$Q(y) = \frac{1}{(y+\Sigma^2(y))^2}\left[(\Sigma^2(y) - y)\hat{Q}(y) + \Sigma^2(y)\right] \ , \tag{24}$$

$$P(y) = \frac{1}{(y+\Sigma^2(y))^2}\left[-2\hat{Q}(y) + (y+\Sigma^2(y))\hat{P}(y) + 2\Sigma(y)\Sigma'(y)\right] \ . \tag{25}$$

The Bethe-Salpeter equation derived above is a linear integral equation for $\hat{Q}(x)$

and $\hat{P}(x)$,

$$(1+K[\lambda,\Sigma])\begin{pmatrix}\hat{Q}\\ \hat{P}\end{pmatrix}=C[\lambda,\Sigma] \quad , \tag{26}$$

where the kernel and the inhomogeneous term are local functionals of $\Sigma(x)$. Thus, given a $\Sigma(x)$ (which determines a vacuum), we have a unique solution for $\hat{Q}(x)$ and $\hat{P}(x)$. Then, with Eq.(22), we evaluate f_π.

If we take a vanishing λ in the above equation, leaving Σ unchanged, $K[\lambda,\Sigma]$ and $C[\lambda,\Sigma]$ both vanish, and we have a trivial solution,

$$\hat{Q}(x)\equiv 0 \quad , \quad \hat{P}(x)\equiv 0 \quad . \tag{27}$$

Then f_π is evaluated as,

$$f_\pi^2=\frac{N}{4\pi^2}\int_0^\infty x dx\frac{\Sigma(x)(\Sigma(x)-x\Sigma'(x)/2)}{(x+\Sigma^2(x))^2} \quad . \tag{28}$$

This formula is nothing but the Pagels-Stoker approximation. Thus we understand that the approximation might be applicable in the weak λ regime. As we see later, due to the asymptotic freedom of QCD-like theories, the Pagels-Stoker approximation gives rather good estimates.

Evaluating f_π

In numerical calculations, we discretize the Bethe-Salpeter equation, and it is now equivalent to a multi-dimensional linear equation. Inverting a coefficient matrix, we get a solution. Note that the convergence of f_π defined by Eq.(22) is not trivial. Analyzing the asymptotic behavior of solutions, we have

$$\hat{Q}(x\to\infty)\to\frac{c}{x} \quad , \quad \hat{P}(x\to\infty)\to\frac{-2c}{x^2} \quad , \quad c\text{: constant} \quad , \tag{29}$$

and this assures the ultraviolet convergence of the integral(22) for f_π to this order of approximation.

Actual calculations are in progress, and here we will present only some preliminary results. As is described before, we introduce an infrared cutoff, t_{IF}, below which the coupling constant ceases to grow. With this running coupling constant $\lambda(x)$, we solve the Schwinger-Dyson equation, and get solution $\Sigma(x)$ for each set of parameters B and t_{IF}. We have to calculate the ultraviolet behavior of $\Sigma(x\to\infty)$ precisely so that our renormalization procedure for $\langle\bar{\psi}\psi\rangle$ might work.

The number of discretization points for Bethe-Salpeter equation is taken as 150, that is, 150 for $\hat{Q}(x)$ and $\hat{P}(x)$ each. Thus we solve 300-dimensional linear equation. Here we should be careful to get high precision, since the kernel is a rapidly varying function, which results in a rather peaked integrand for f_π.

In the following, we show some preliminary results of numerical calculations. We set the ultraviolet cutoff to be

$$\ln\left(\Lambda^2/\Lambda_{\rm QCD}^2\right) = 21 \ , \quad \ln\left(\Lambda^2/\Sigma^2(0)\right) = 30 \ \text{ for } \ B = 0 \ , \tag{30}$$

where $\Lambda_{\rm QCD}$ represents the scale parameter of the leading logarithmic running coupling constant, that is, at $\Lambda_{\rm QCD}$ the coupling constant diverges. The infrared cutoff is parametrized by $t_{\rm IF}$, below which the coupling constant stops diverging, and tends to be a constant. Avoiding possible discontinuity in $\Sigma'(x)$, we set $\lambda'(x)$ continuously goes down towards 0, linearly in $\ln x$, up to a fixed point μ_0 defined by

$$\ln\left(\mu_0^2/\Lambda_{\rm QCD}^2\right) = -2 \ . \tag{31}$$

Below this point μ_0, the coupling constant does not run at all.

A schematic view of the Bethe-Salpeter kernel is in Fig. 1, where we show two cases of $B = 9/16$ (three triplets QCD) and $B = 0$ (non-running coupling constant). The corresponding $\Sigma(x)$, solutions $\hat{Q}(x)$, $\hat{P}(x)$ and the integrands for f_π, and $\langle\bar{\psi}\psi\rangle$, are shown in Fig. 2a and 2b. As for the integrand for f_π, we plot both our ladder exact integrand and the Pagels-Stoker integrand (f_π-PS). Both are strongly peaked at near the peak of $|\Sigma'(x)|$.

Changing the infrared cutoff $t_{\rm IF}$, we get Fig. 3a, where we take cases of $B = 9/16$ (three triplets QCD) and $B = 23/120$ (single sextet QCD). Fig. 3b is a zoom-up of $B = 9/16$ case. Lower $t_{\rm IF}$ drives the infrared coupling constant larger. In any case, we set a renomalization condition, f_π (ladder exact) = 94MeV, which is shown as a thick line of every plots. The reason of taking a specific value of 94MeV is simply for us to easily catch relative scales of various parameters. One sees the following results of our ladder exact calculation.

1. The Pagels-Stoker approximation is not so bad. We understand that it is due to the asymptotic freedom of the theory, that is, after the most essential part of dynamics is taken as a solution of the Schwinger-Dyson equation $\Sigma(x)$, the rest is controlled by the rather 'weak' coupling λ, as is seen in Eq.(26).
2. However, the ladder exact results give some corrections to the Pagels-Stoker formula, and its amount can be very important when used in case of electroweak symmetry breaking where ρ parameter is strongly constrained. It might cause larger deviation of ρ from unity, which is under calculation.
3. The renormalization of $\langle\bar{\psi}\psi\rangle$ works excellently. We get an almost stable value of the renormalized $\langle\bar{\psi}\psi\rangle_{\rm R}$, stable against change of the infrared cutoff parameter $t_{\rm IF}$, while the unrenormalized counter part $\langle\bar{\psi}\psi\rangle_{\rm U}$ depends much on $t_{\rm IF}$.
4. On the other hand, $\Sigma(0)$ depends quite a lot on $t_{\rm IF}$. In many articles, $\Sigma(0)$ is used as a mass scale of the spontaneous symmetry breaking. One should note that it does depend on the infrared structure of the running coupling constant, and thus it is not a good measure of physics.
5. As for $\Lambda_{\rm QCD}$, it depends on $t_{\rm IF}$. However, the dependence is negligible for lower $t_{\rm IF}$ region, while in such region, $\Sigma(0)$ diverges.

6. The case with a sextet, $B = 23/120$, gives more stable results against the infrared cutoff. This is because the running of the coupling constant is slower than the three triplet case ($B = 9/16$). Thus, the dynamical structure is controlled by the B parameter.

7. Case with a sextet shows another differences. For example, take a ratio of f_π or $\langle\bar{\psi}\psi\rangle_{\mathrm{R}}$ to Λ_{QCD},

$$\frac{(f_\pi/\Lambda_{\mathrm{QCD}})_{\mathrm{sextet}}}{(f_\pi/\Lambda_{\mathrm{QCD}})_{\mathrm{triplet}}} \sim 2.5 \ , \tag{32}$$

$$\frac{(\langle\bar{\psi}\psi\rangle_{\mathrm{R}}/\Lambda_{\mathrm{QCD}})_{\mathrm{sextet}}}{(\langle\bar{\psi}\psi\rangle_{\mathrm{R}}/\Lambda_{\mathrm{QCD}})_{\mathrm{triplet}}} \sim 2.8 \ . \tag{33}$$

This is a reasonable enhancement of the symmetry breaking mass scale in case of higher (larger $C_2(F)$) representation fermion, although quantitative comparison with other calculations needs more analyses. Furthermore, the difference of the above two ratios, 2.5 vs 2.8 might have some physical significance.

References

1. K-I. Aoki, M. Bando, T Kugo and H. Nakatani, in preparation.
2. K. Higashijima, Phys. Rev. D29 (1984) 1228.
3. K-I. Aoki, M. Bando, K. Hasebe, T Kugo and H. Nakatani Prog. Theor. Phys. 81 (1989) 866; Prog. Theor. Phys. 82 (1989) 1151.
4. H. D. Politzer, Nucl. Phys. B117 (1976) 397.
5. T. Kugo, Phys. Lett. 76B (1978) 625.
6. T. Morozumi and H. So, Prog. Theor. Phys. 77 (1987) 1434;
7. H. So and M. Suwa, NIIG-DP-89-4.
8. Talk by I. Watanabe at this Workshop; K. Fukazawa, M. Inoue, T. Muta, J. Saito, I. Watanabe and M. Yonezawa, preprint, HUPD-8917.
9. H. Pagels and S. Stoker, Phys. Rev. D20 (1979) 2947.
10. V.A. Miransky, M. Tanabahi and K. Yamawaki, Mod. Phys. Lett. A4 (1989) 1043; Phys. Lett. B221 (1989) 177.
11. Bob Holdom, Phys. Lett. B226 (1989) 137.
12. T. Appelquist, T. Takeuchi, M. B. Einhorn and L.C.R. Wijewardhana, Phys. Lett. B232 (1989) 211.
13. Y. Nambu, preprint, EFI-89-08; EFI-89-39.
14. W. A. Bardeen, C. T. Hill, M. Lindner, preprint, FERMILAB-PUB-89/127-T.

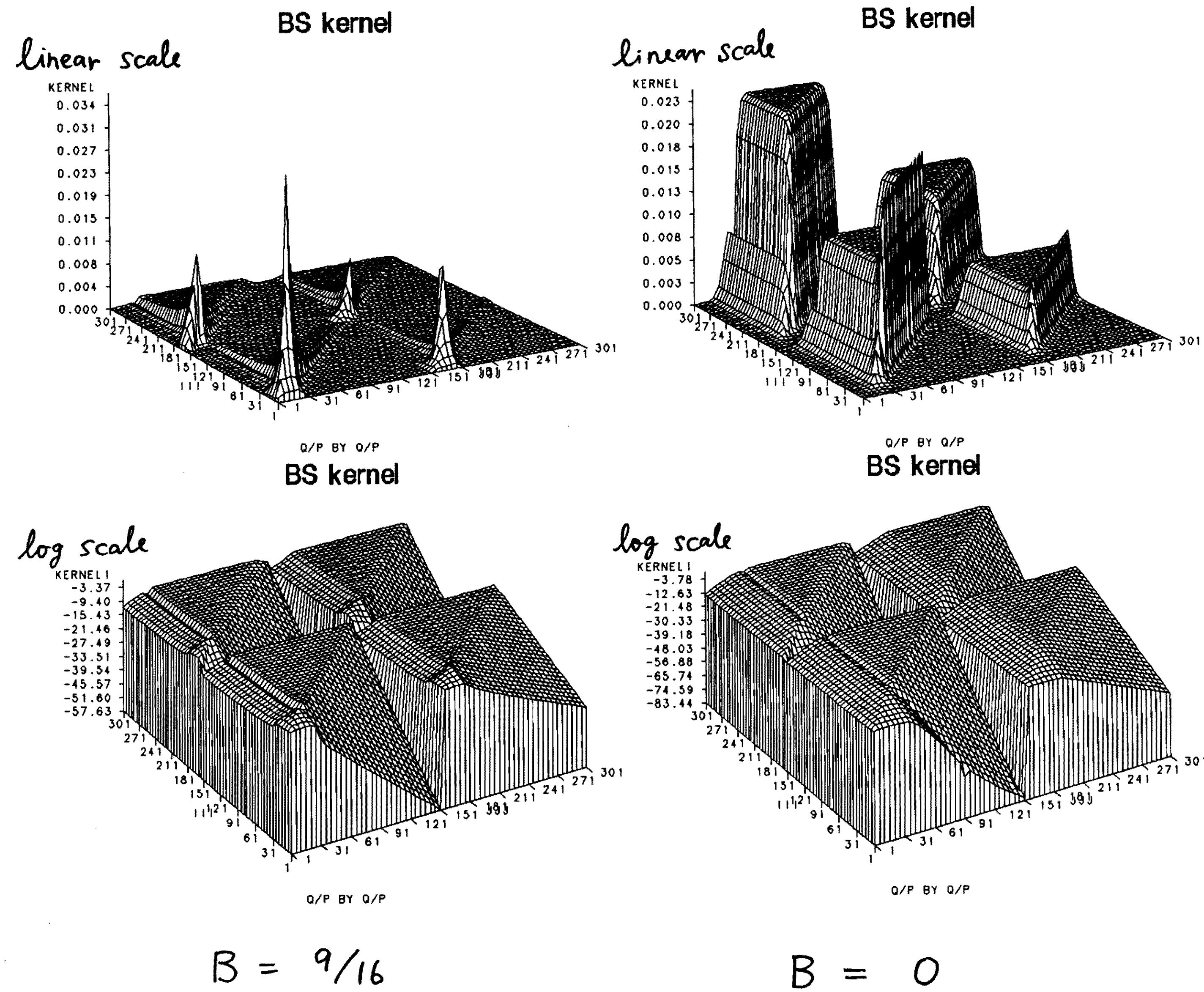

Fig. 1 Schematic View of Bethe Salpeter Kernel

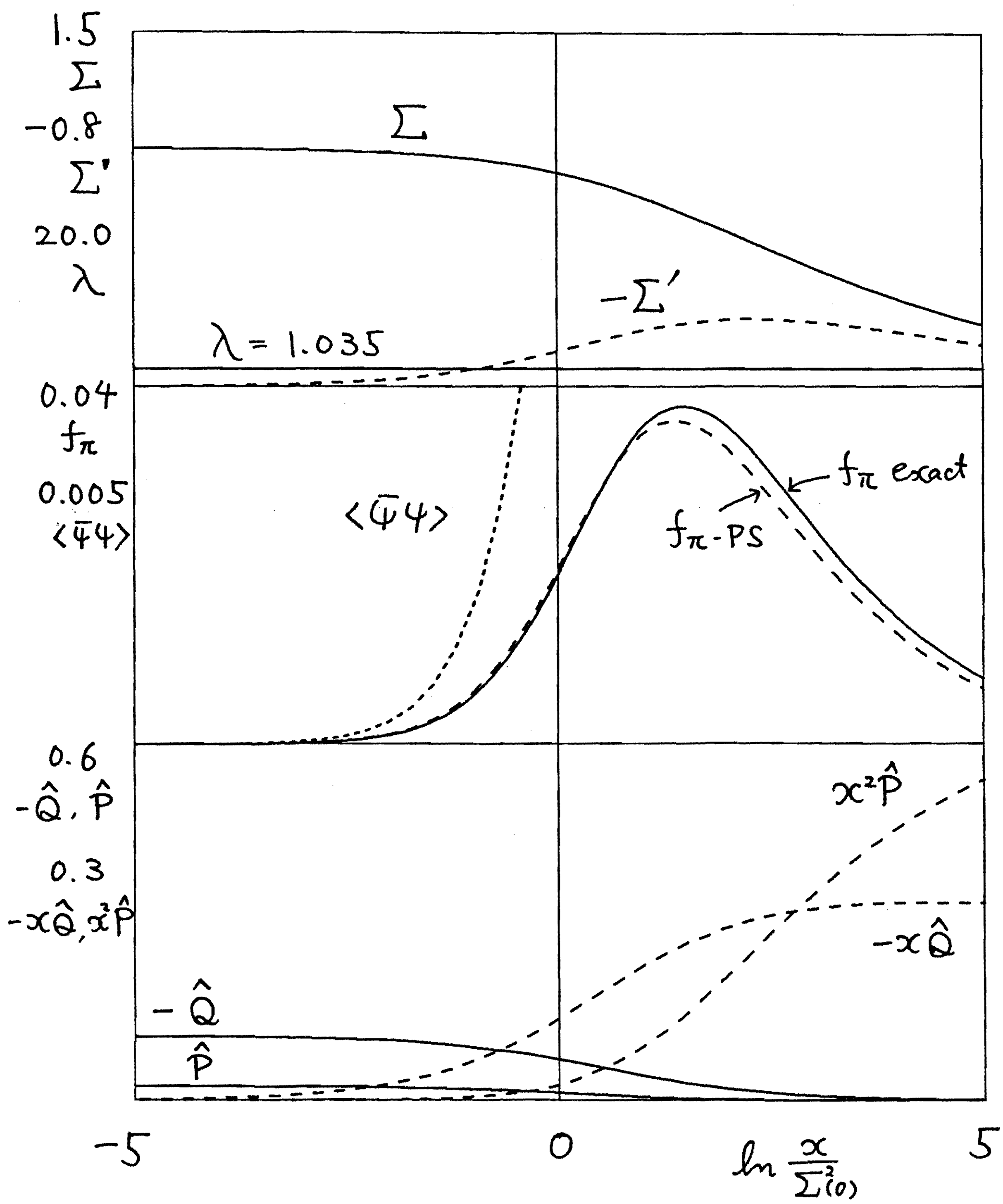

Fig. 2a Solutions of Schwinger-Dyson and Bethe-Salpeter Equations
Case of $B = 0$ (Fixed Coupling Constant)

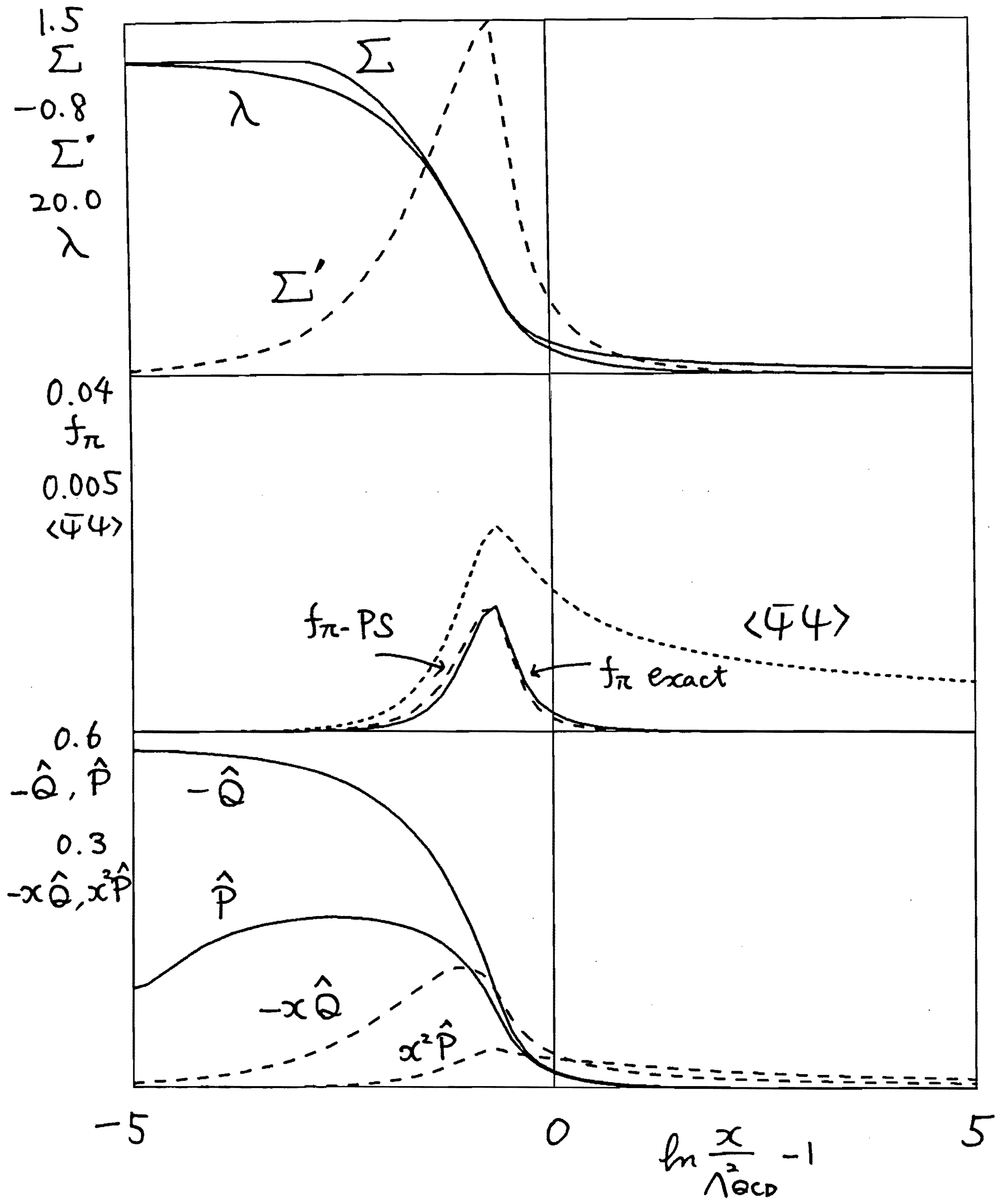

Fig. 2b Solutions of Schwinger-Dyson and Bethe-Salpeter Equations
Case of $B = 9/16$ (three triplets QCD)

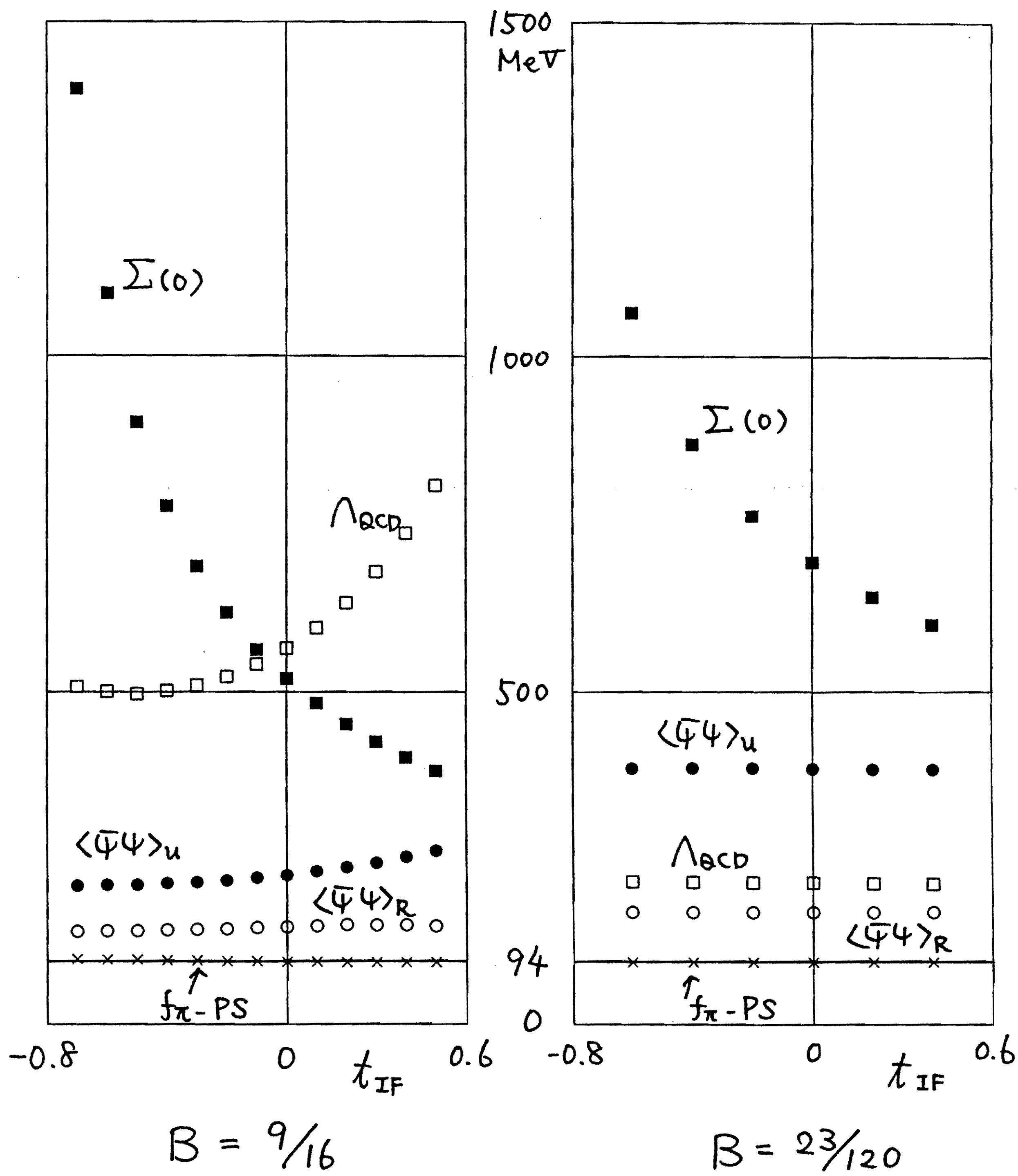

Fig. 3a Value of Physical Quantities
Case of $B = 9/16$ (three triplets QCD)
and $B = 23/120$ (single sextet)

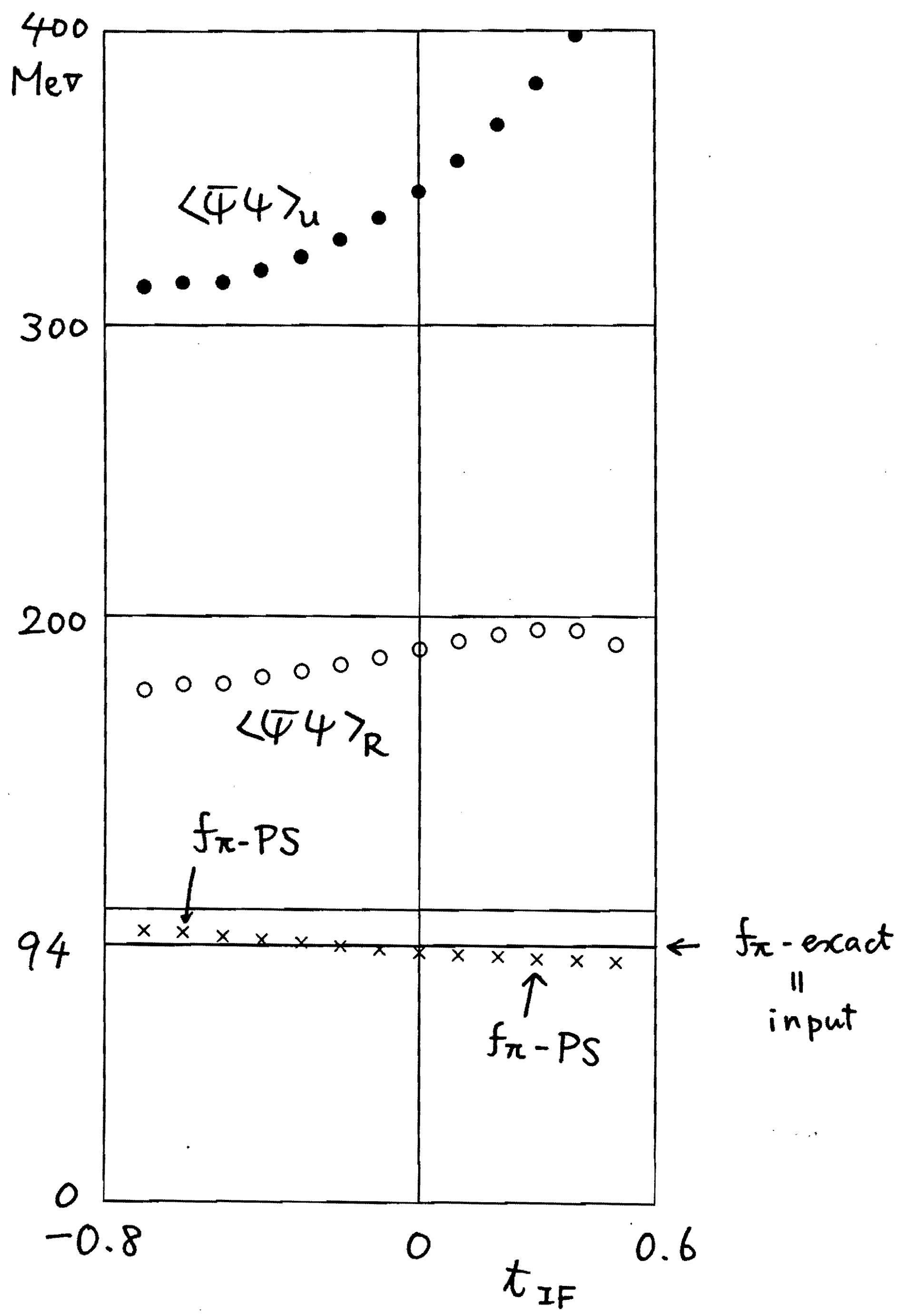

Fig. 3b Value of Physical Quantities
Case of $B = 9/16$(three triplets QCD)