Masses and Interactions in Quantum Chromodynamics

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Abstract

Correlations of composites corresponding to baryons and mesons are composed within the derivative expansion. The expansion in energy scales permits a quantitative, algebraic description at various energy scales in QCD. The masses in QCD are derived utilizing a proposed line interaction, with explicit checks of the masses up to the baryonic decuplet.
1 Introduction

Correlator calculations in quantum chromodynamics are difficult due to the complicated nature of the diagrammatic expansion. The derivative expansion has been recently developed to simplify these calculations, and in particular, to reduce the complicated integrals to a set of almost free-field ones. Expansions pertaining to colliders are naturally formulated in terms of energies, from lower to higher ones; the derivative expansion is in this spirit. The expansion is equivalent to the usual infinite number of loop graphs, but with the small parameter being the dimensionless ratio of energy scales as opposed to a coupling constant that could be of order unity.

One of the outstanding questions in quantum chromodynamics, and generally in any gauge theory, is the origin of the masses of the gauge invariant states - in particular without resorting to computationally intensive lattice gauge theory. In this work, in addition to formulating the correlator expansions, we derive the set of masses of the mesonic and baryonic tower of QCD masses, using the derivative expansion together with the inclusion of Wilson-like line integrals. The latter operators are suited quite naturally within the derivative expansion.

The masses of the composites containing the (s,d,u) multiplet, i.e. mesons and baryons and etc, are summarized here for convenience. They follow via a formula,

$$\langle O(x_1)O(x_2) \rangle = \sum_{j=1}^{\infty} c_j(x_1,x_2) g^{2j} e^{-m|x_1-x_2|g^2} .$$

(1.1)

At $g \sim 1$, which is realistic for quantum chromodynamics, the 'propagator' or 2-point function for the composites has the usual spatial dependence. It is interesting that the coupling dependence of $g^2$ enters in this manner, which explains the differences in the masses via its flow in energies. The lowest order approximation to these masses (containing the s,d, and u quarks) are found via the term in the exponential in (1.1). It is calculated to contain two terms,

$$m = \sum_{\nu} ^{(j)} + \tilde{f} .$$

(1.2)

The terms are due to the Wilson-like interaction,

$$\tilde{f} = 150(4I - \frac{2}{3})10^6 \text{ eV} ,$$

(1.3)
together with the individual fermionic mass terms in the (s,d,u) at the QCD scale,

\[ \sum m^{(j)}_\psi = 150(N_\psi + N_S)10^6 \text{ eV}. \] (1.4)

There are perturbative corrections to this formula via the power series \( C_j(x_1, x_2) \) in (1.1), but the zeroth order formula in (1.2) agrees well with the known mesonic and baryonic composites.

This work is contained in the context of quantum field theory, and its placement in the context of supergravity is straightforward. In previous work the mass generation of the fermion species content has been explained via gauge and gravitational instantons [1], and when combined with the work here lends a fundamental explanation of the generation of the masses of the physical QCD sector. Also, the gauge theory work presented in the current text may be generalized to finite temperature and supersymmetric field theories, but is not included.

2 Brief review of QCD derivative expansion

Derivative expansions of quantum field and string theories have recently been developed with several goals, one of which is to determine analytically their nonperturbative properties [2]-[9]. This expansion is identical to the usual diagrammatic expansion in loops, but with an expansion in momentum scales as opposed to couplings. As a result, this approach commutes with dualities in supersymmetric theories and is generally nonperturbative in couplings. One facet of this approach is that all integrals may be performed, and theories treated in this expansion have amplitudes that may be determined by a set of algebraic recursive equations, which are almost matrix-like. Gauge theories have been examined briefly in [3] and [4] in this context; we review the description of microscopic correlations describing amplitudes and composite correlations modeling nucleon interactions. (These correlators are quantitatively related, however, for clarity we describe both.)

The Lagrangian considered is

\[ \mathcal{L} = \int d^4x \left( -\frac{1}{4} F^2 + \psi^a \not{D}\psi_a \right) \] (2.1)

quantum chromodynamics; the non-perturbative properties via coherent state Wilson loops and instantons are also examined. The effective theory, expanded in derivatives is found from all possible combinations of gauge invariant operators \( \mathcal{O}^{(j)}(x) \),
\[ S = \int d^4 x \sum_{j=1}^{\infty} h_j(g, \theta) \mathcal{O}^{(j)}(x) \] (2.2)

and \( h_j(g, \theta) \) contains the full coupling dependence. Example gauge invariant operators are \( \text{Tr} F^2 \frac{1}{\not{D}} F^2 \), and \( \frac{1}{m_{\psi}^4} \). In the derivative expansion, self-consistency of the effective action with unitarity, implemented via sewing, allows for a determination of the functions \( h_j \). The action is next examined with respect to both logarithmic modifications of the terms and regulator dependencies.

In addition to the hard dimension labeling the operator, logarithms also in general modify the form of the generating function, through, for example,

\[ \text{Tr} F^2 \left[ \ln^{n_1}(\Box) \ln^{n_2}(\Box) \ldots \ln^{n_m}(\Box) \right] \Box^2 F^2, \] (2.3)

with covariantized boxes. The presence of logarithms is required by unitarity and are generic in loop integrals; there are generically \( L \) multiplicative log terms at loop order \( L \) in the loop expansion. These terms may be computed either in a direct sense via their inclusion in the effective theory, or may be determined by unitarity. The logarithms are required via \( \Im S = S^\dagger S \) and may be computed from the analytic terms after their coefficients are determined.

The form of the series expansion in terms of the operators depends on how the gauge field is regularized, in string theory with the string inspired regulator and dimensional reduction. In the former there is a dimensional parameter \( (\alpha') \) acting effectively as a cutoff; there may in general be other geometric scales depending on the model that may serve in the same role as \( \alpha' \) in the following. As the generating function contains one-particle reducible graphs, there must be inverse powers of derivatives, which are local in the sense that in momentum space these terms simply model the propagator \( 1/(k^2 + m^2) = 1/m^2 \sum (-k^2/m^2)^n \); in the massless case the \( 1/\Box \) occurs, and in gauge theory their universal form, in the sense of independence of the number of external particles in a correlator, is expected based on collinear and soft factorizations. Last, on-shell gauge theory amplitudes have infrared singularities; in the x-space expressions these singularities are absent as the lines are effectively off-shell.

The regulator dependence of an \( \alpha' \), or other dimensional parameter such as a mass term not depending on \( \alpha' \) or a geometric parameter, follows in a straightforward sense by allowing their powers to occur in the expansion, i.e. \( \alpha^m m^{-p} \). This occurs in quantum field theory via the decoupling of massive states, and in low-energy effective field
theory as an expansion about an ultra-violet cutoff. The general term in the effective action we consider is determined by including all gauge invariant terms discussed in the previous paragraph together with these dimension parameters. (In a dimensionally regularized theory the effective ultra-violet cutoff is absent and only the mass terms, with any other dimensionful parameters, occurs in the derivative expansion.) Parameters such as Yukawa couplings in a spontaneously broken context occur in a polynomial sense as dictated by perturbation theory.

The gauge coupling expansions of $h_j$ follow from the usual expansion of the gauge theory amplitudes,

$$h_j(g, \theta) = \sum_{n=0}^{\infty} a_j^{(n)} g^{2+2n}, \quad (2.4)$$

and a series of non-perturbative terms,

$$\tilde{h}_j(g, \theta) = \sum_{n=1}^{\infty} \tilde{a}_j^{(n)} e^{n \left( -\frac{4\pi}{g^2} + i \frac{\theta}{\pi} \right)}. \quad (2.5)$$

The coefficients $a_j^{(n)}$ are determined via the sewing relations. The instantons in the background field method generate $\tilde{a}$; potentially these contributions are redundant with the exponentiated gauge field integrals.

For purposes of reviewing we formulate the four-point scattering of gauge bosons. The effective vertices to be inserted into the derivative diagrams are the interactions $(A, A^m\psi^k)$, $(A^2, A^{m-1}\psi^k)$, and $(A^3, A^{m-2}\psi^k)$. They are found by variation of the effective action,

$$v_{1,2,3}^{\mu,\nu;m,n}.$$ (2.6)

The unitarity relation that generates the full four-point amplitude function, in $k$-space, is,

$$\int \prod_{q=1}^{m+n} d^dq_j \quad v_1^{m,n}(k_1; q_i) \prod_m \Delta_A \prod_n \Delta_{\psi} \quad v_3^{m-2,n}(q_i; k_2 + k_3 + k_4) \quad (2.7)$$

$$+ v_2^{m-1,n}(k_1 + k_2; q_i) \prod_m \Delta_A \prod_n \Delta_{\psi} v_2^{m-1,n}(q_i; k_3 + k_4) \quad (2.8)$$
in which the full derivative dependence has been implied in the vertices, and should be expanded termwise. The integrals are easier to evaluate in $x$-space, as in [5]. The indices of the fields have been suppressed for notational purposes. The propagators are indexed by $A$ and $\psi$. In order to generate the full amplitude at the four-point, including the coefficients at general order $g^{2+L}$, we have to include the sewing relations that generate all of the other vertices; the infinite tower of unitarity relations are coupled and together generate the complete effective action corresponding to the loop expansion. In general this appears complicated; however, to a finite order in coupling only a finite number of vertices are involved, which is small at low orders in coupling. Furthermore, all of the integrals may be performed (including the massless ones using analytic methods as outlined in [4] and [5]). If interested in computing to high orders in $g_{YM}$ then the method is amenable via direct calculation or in a computer implementation.

An explicit evaluation of the terms and integrals have been performed in massive scalar field theory, in an arbitrary dimension. The reader is referred to [5] to see the simplest implementation.

Next we examine the composite operator correlations; the composite operators $O_j$ model the bare nucleons in terms of free particle states. Flow of momentum amongst the various free-particle states in the composite operator is general. A schematic is illustrated in figure 2. In the correlations involving the composite operators, and in order to make contact with the parton model in perturbation theory, the internal lines of the operators (nucleons) are connected to a) full interaction vertices or b) from one nucleon to another. In other words, the vertices (depicted in figure 4) are one particle reducible so that the perturbative contributions to the interactions in the usual loop expansion is obtained. The interactions are depicted in figure 3 for a sample collision of three $\psi^3$ hadrons. Figure 4 illustrates a usual interaction graph with that in the derivative expansion.

3 Exponential insertions and Masses

In this section integrals along 1-cycles and 2-cycles are included in the expansion, via the gauge field $A$ and the curvature $F$. In doing so the masses of the nucleons are derived. Their inclusion in the correlations described in the previous section depart from the parton model in that the realistic masses of the nucleons are obtained.
Figure 1: The sewing relation illustrated at 4-point. Permutations are not included.
Figure 2: Momentum flow of diagram and comparison with the usual parton picture.
Figure 3: A sample composite operator correlation corresponding to three baryon interaction.
Figure 4: A sample relation between the loop graph and derivative graph, found by expanding the integral.
First the exponentials are obtained, together with their interacting exponentiated relatives, mediated by the microscopic theory (interacting in gauge coupling $g$); both are depicted in Figure 5. The exponentials resemble flux tubes.

The exponential operator we consider is

$$\frac{1}{2} e^{-\alpha \oint A P_G},$$

with end-points fixed at the location of the composite operators. Another operator that may be considered involves a curvature term,

$$e^{\tilde{\alpha} \oint F \tilde{P}_G},$$

with an integral taken over compact Riemann surfaces attached at the operator locations. The line integral we take to be oriented in accord with the action of the isospin operator, the projection $P_G$ is an isospin operator acting on the composite operators, with an explicit factor of a half inserted to agree with the orientation of the integral; Its eigenvalue on the nucleon made up of $n$-fermions with maximal isospin $I$ is

$$\lambda = 4I - \frac{2}{3}.$$ (3.3)

The contraction of a product of free-particle Wilson lines between two points $x = x_1 - x_2$, by dimensional grounds and finiteness of the integrals,

$$\prod_{n} e^{-\alpha \oint A} \rightarrow e^{n \alpha^2/2 \ln(x^2)},$$

with $\alpha$ a general coupling constant, taken for example as $g_{QCD} \sim 1$ near the QCD scale. The summation over an arbitrary number of closed loops, without taking into account interactions between them is,

$$\sum_{n=0}^{\infty} \frac{(-\lambda)^n}{n!} e^{\alpha^2/2 \ln(x^2/\mu^2)} = \sum_{n=0}^{\infty} \frac{(-\lambda)^n}{n!} \left(\frac{x^2}{\mu^2}\right)^n \alpha^2/2 = e^{-\lambda(x^2/\mu^2)\alpha^2/2}. $$ (3.5)
one the exponential has the form to model a mass term in a propagator. The coupling in QCD is of order unity, and as a result these contributions naturally model the mass of a nucleon.

A nucleon is not a fundamental particle, and as a result one does not expect a propagator in the sense of $\partial \partial \Delta = -\delta$ to model its dynamics. The free particle composite of $n$ fermions, as in the parton model, has the form $\Delta^n(x)$, containing $\exp(-xnm_\psi)$. The mass of the nucleonic state is found by computing the correlation $\langle O(x_1)O(x_2) \rangle$,

$$\langle O(x_1) \left[ \text{Tr} \prod e^{-\alpha \oint A P_G} \right] O(x_2) \rangle \sim C(x)e^{-f(x^2/\mu^2)^{1/2}} \quad (3.6)$$

$$C(x)e^{-\sum m_j^2 e^{-\tilde{f} x}} \sim 1; \quad \tilde{f} = f/\mu \quad (3.7)$$

The projection operator is taken to act in both directions along the line integral of $\oint A$. The variable $\tilde{f}$ is,

$$\tilde{f} = 150(4I - \frac{2}{3})10^6 \quad \text{eV} \quad (3.8)$$

and the mass sum for fermions in the (s,d,u) at the QCD scale is approximately,

$$\sum m_j^2 = 150(N_\psi + N_S)10^6 \quad \text{eV} \quad (3.9)$$

One can check that the mass formula agrees quite well with the masses in the baryon octet and decuplet, and the meson vector nonet (the bare masses of the fermions are taken as approximately 150 MeV and 300 MeV for the (u,d) and s quarks at the QCD scale.) This approximation is in the free-field point of view, and resummations of gauge interactions could modify the 'mass' of the nucleon - in quotes because an interacting nucleon is not really a particle. The meson octet does not nearly agree as well as the rest; possibly this is due to the odd parity of these states and electroweak interactions. If the quarks were massless, then a 2-quark state is differentiated from $n > 2$ because of infra-red divergences, as can be seen by the first order gluonic correction to the two meson correlation. Furthermore, it would be interesting to attempt to derive the subtle mass difference between the proton and neutron, or other degenerate states; the perturbative corrections to the mass calculation are desired.
To compare, we list the known meson and hadron masses for the first few multiplets. The quark content and masses of the baryonic octet are,

\[
\begin{pmatrix}
  I_3 : & -1 & -
  \frac{1}{2} & 0 & 
  \frac{1}{2} & 1 \\
  N(939) : & n & udd & \quad & p & uud
\end{pmatrix},
\]

\[
\begin{pmatrix}
  \Sigma(1193), \Lambda(1116) : & \Sigma^- & dds & \quad & \Sigma^0 & \Lambda & uds & \quad & \Sigma^+ & uus
  \\
  \Theta(1318) : & \Theta^- & dss & \quad & \Theta^0 & uus
\end{pmatrix} \quad (3.10)
\]

and the same for the baryonic decuplet,

\[
\begin{pmatrix}
  I_3 : & -
  \frac{3}{2} & -1 & -
  \frac{1}{2} & 0 & 
  \frac{1}{2} & 1 & 
  \frac{3}{2} \\
  \Delta(1232) : & \Delta^- & ddd & \quad & \Delta^0 & ddu & \quad & \Delta^+ & duu & \quad & \Delta^{++} & uuu
  \\
  \Sigma(1384) : & \Sigma^- & dds & \quad & \Sigma^0 & dus & \quad & \Sigma^+ & uus
  \\
  \Theta(1533) : & \Theta^- & dss & \quad & \Theta^0 & uus
  \\
  \Omega(1672) : & \Omega^- & sss
\end{pmatrix} \quad (3.11)
\]

These masses agree very well with the generated mass formula.

The pseudoscalar mesons have masses: \( \pi^\pm \) \((u\bar{d})\) 140 MeV, \( \pi^0 \) \((d\bar{d} − u\bar{u})/\sqrt{2}\) 135 MeV; \( K^\pm(u\bar{s})\) 494 MeV, \( K^0(d\bar{s})\) 498 MeV; \( \eta_8 \) \((d\bar{d} + u\bar{u} − 2s\bar{s})/\sqrt{6}\) 549 MeV; \( \eta_0 \) \((d\bar{d} + u\bar{u} + s\bar{s})/\sqrt{3}\) 958 MeV. The masses of the vector nonet multiplet are: \( \rho \) 776 MeV; \( K^\pm \) 892 MeV; \( \omega(s\bar{s}) \) 783 MeV; \( \phi \) \((u\bar{u} + d\bar{d})/\sqrt{(s)}\) 1019 MeV. The mass formula agrees well with the vector nonet and requires improvement with regards to the pseudoscalars.

A primary aspect of the mass formula is that it exhibits approximate Reggeization. The mass found in this approximations lay on \( R^2 \) parameterized by \( I \) and fermion number \( N_\psi + N_S \), illustrated in figure 6.

The full 2-point correlator is obtained via inserting the vertex between the two free nucleonic states, and also summing the exponentiated gauge interactions, depicted in figure 7. The vertex contracts any number of sets of lines, and there may be disconnected components connecting the lines between \( x_1 \) and \( x_2 \), as explained in section 2. The correlator has a power series expansion,

\[
\langle O(x_1)O(x_2) \rangle = \sum_{j=1}^{\infty} c_j(x_1, x_2) g^{2j} e^{-m|x_1−x_2|^2/2}, \quad (3.12)
\]

at \( g \sim 1 \), with corrections when the coupling is away from unity. The vertex is obtained via the method in the previous section. The potential soft dimensional terms in \( c_j \) naively could resum to alter the mass term obtained from the interacting exponential gauge terms.
Figure 5: The graphical interpretation of the $e^{-\alpha f^A}$ inclusion.
Figure 6: Mass patterns as a function of the maximal isospin representation and fermion number.
Figure 7: Sample exponential terms contributing to the 2-point correlator.
4 Nucleon interactions

The interactions follow via the interactions of the composite operators as described in section 2; these reproduce on a microscopic setting the parton model. The additional interactions to be included are those of the Wilson loops. Contrasted with the simplest two operator correlations, these exponentiated paths may join points $x_j$ in a number of ways for a given number of exponentials. As a nucleon is not a point particle, the latter interactions model an effective mass of the constituent composite nucleons.

5 Discussion of general gauge theory

Two primary differences between QCD and a general gauge theory are: the flow of the coupling constant and the energy scale of the theory. These two properties change the mass formula and the scattering of the composite states. For example, the mass formula has a $x^g$ in the exponential and depending on the UV properties of the theory (from the microscopic theory) the properties of the two-point correlator change; as usual, the composites may break into free particle constituents with an effective mass containing only the bare fermions. The flow equations of the coupling constant dictate the fixed points of the theory.

6 Conclusion

The correlations of nucleonic states are examined within the derivative expansion; these correlations are identical to those of quantum chromodynamics. Additional exponential interactions, i.e. Wilson loops, are added to the interactions. Masses of the nucleons are derived, and excellent agreement is found with the observed parameters. They have a non-trivial dependence on the coupling constant, and in an asymptotically free limit, degenerate into the bare masses of the quark content without the gluon contribution. Small differences in the masses at given strangeness number within the multiplets are potentially found via perturbative corrections, for example, the mass splittings between the neutron and proton as well as other sets. The fundamental masses of the quarks have also been analyzed and derived in the context of M theory via gravitational (and gauge) instantons.

Explicit calculations and diagrams are presented that explain the interactions and methodology. In general, gauge theories may be examined in the same approach;
the running of the coupling constant is governed by the perturbative expansion. The
derivation of the masses and their role in the dynamics may be found in the usual
loop expansion. Generalization to gauge theory with matter at finite temperature
could be explained in the context presented in this work (with testable predictions
for example at the RHIC collider).

In the derivative expansion the calculations are much simplified over the usual
perturbative loop approach. The former is well suited to extend quantitative QCD
work well into the lower energy regimes and to have a variety of quantitative appli-
cations, with latitude, at various energies.

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Accurate Determination of Hadron and Baryon Masses

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Abstract

In previous work a mass formula was generated within the derivative expansion that models the masses of the gauge invariant composites, such as mesons, hadrons, and baryons. A possible improvement based on uniform corrections is given in this work that models the masses to an approximate .1 percent.
In (1) QCD was reformulated in the derivative expansion. In there, a mass formula was generated that models the masses of the gauge invariant composites,

\[ M = 150(N_\psi + N_s) + 150\left(4I - \frac{2}{3}\right) \]

(6.1)

in terms of \(10^6\) eV. The first term is found from the bare masses of the fermions, and the second term is modeled from a superposition of 'Wilson' lines. There is a similar mass formula (2) for the fundamental fermions, of the form,

\[ 10^{n-4} \pm 2^m 5^i \times 10^{-3} \text{ GeV} \]

or in MeV,

\[ 10^{n-1} / -2^m 5^i \times 10^{-3} \text{ MeV} \quad n = 1, 2, \ldots 6 \]

(6.3)

which has a leading term plus a subleading term, which is quantized.

It turns out that the addition of

\[ M' = 18N \times 10^6 \text{ eV} \]

(6.4)

to the formula in (6.1), with \(N\) a positive or negative integer models the masses to an approximate .1 percent. The formula pertains to the hadronic and baryonic multiplets.

The states and masses of the baryonic and hadronic multiplets are,

\[
\begin{pmatrix}
I_3 : & -1 & -\frac{1}{2} & 0 & \frac{1}{2} & 1 \\
N(939) : & n & udd & p & uud \\
\Sigma(1193), \Lambda(1116) : & \Sigma^- & dds & \Sigma^0, \Lambda & uds & \Sigma^+ & uus \\
\Theta(1318) : & \Theta^- & dss & \Theta^0 & uus \\
\end{pmatrix}
\]

(6.5)

and the same for the baryonic decuplet,

\[
\begin{pmatrix}
I_3 : & -\frac{3}{2} & -1 & -\frac{1}{2} & 0 & \frac{1}{2} & 1 & \frac{3}{2} \\
\Delta(1232) : & \Delta^- & ddd & \Delta^0 & ddu & \Delta^+ & duu & \Delta^{++} & uuu \\
\Sigma(1384) : & \Sigma^- & dds & \Sigma^0 & dus & \Sigma^+ & uus \\
\Theta(1533) : & \Theta^- & dss & \Theta^0 & uus \\
\Omega(1672) : & \Omega^- & sss \\
\end{pmatrix}
\]

(6.6)
These masses agree very well with the formula in (6.1).

The differences between the masses in (6.1) and the 'actual' masses are

\[
\begin{align*}
N(939) &\quad - 18 + 7 \\
\Sigma(1193), \Lambda(1116) &\quad 5(18) + 3, 18 - 2 \\
\Theta(1318) &\quad 4(18) + 2 \\
\Delta(1232) &\quad - 18 \\
\Sigma(1384) &\quad - 18 + 2 \\
\Theta(1533) &\quad - 18 + 1 \\
\Omega(1672) &\quad - 18 .
\end{align*}
\]

Some of the excited states also follow the same pattern, including, for example,

\[
\begin{align*}
\Xi(1321) &\quad 15(18) + 1 \\
\Xi(1351) &\quad 15(18) + 7 .
\end{align*}
\]

The corrections are normalized with the hadron masses.

The subleading terms of 1, 2, (7) are indicative of another uniform correction. Perhaps the factors of 18 and the prefactors are indicative of instanton-like or non-perturbative corrections. If the factors of 18 are indeed of the correct form, which is especially indicative in the baryonic multiplet, then the masses are found to an approximate .1 percent. The sub-leading terms of 18 × N appear to justify the formula in (6.1), essentially in the baryonic decuplet, and the derivation in [2].
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