Two-Loop QCD Renormalization and Anomalous Dimension of the Scalar Diquark Operator

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Abstract

The renormalization of the scalar diquark operator and its anomalous dimension is calculated at two-loop order in QCD, enabling higher-order QCD studies of diquarks. As an application of our result, the two-loop diquark anomalous dimension in the $\overline{\text{MS}}$ scheme is used to study the QCD renormalization scale dependence of diquark matrix elements of the $\Delta S = 1$ effective weak Hamiltonian.

1 Introduction

Four-quark (or tetraquark) $qq\bar{q}\bar{q}$ states explain the inverted mass hierarchy of the scalar mesons compared to a $q\bar{q}$ nonet in a variety of theoretical approaches [1, 2, 3, 4, 5]. With the inclusion of a gluonium (glueball) state [6], the scalar spectrum below 2 GeV is then understood as mixtures of gluonium, the $q\bar{q}$ nonet, and the $qq\bar{q}\bar{q}$ nonet. The $X(3872)$ [7] and $Y(4260)$ [8] mesons can also be interpreted as four-quark states [9].

Diquark ($qq$) clusters are relevant to the internal structure of hadrons (see e.g., [10, 11]). In particular, Ref. [9] uses constituent models for diquark clusters to study four-quark states. The constituent (scalar) diquark masses that emerge in Ref. [9] are in good agreement with QCD sum-rule analyses of diquarks [12, 13], providing QCD corroboration for the diquark model of four-quark states.

In this paper, we study the renormalization of scalar diquark operators to two-loop order in QCD and thereby obtain the two-loop anomalous dimension of the scalar diquark current. As discussed below, the renormalization of the diquark operator is an essential component of QCD sum-rule analyses, and the anomalous dimension is also necessary for determining the scale dependence of matrix elements of the effective weak Hamiltonian for non-leptonic strange particle decays [14]. Our two-loop results thus enable future QCD studies of diquarks to higher loop order.

The scalar diquark operator in an anti-triplet colour configuration (the “good” diquark in the terminology of Ref. [11]) is given by [12]

$$J_\gamma = \epsilon_{\alpha\beta\gamma}Q^\alpha_i (C\gamma_5)_{ij} q^\beta_j = \epsilon_{\alpha\beta\gamma}Q^T_\alpha C\gamma_5 q_\beta,$$

where the greek and latin indices respectively represent colour and spin degrees of freedom for the quark fields $Q$ and $q$, and $C$ is the charge conjugation operator. The presence of a transposed quark field in (1) implies that the Feynman rule for the three-point function of the diquark operator and $\bar{Q}$, $q$ fields shown in Fig. 1

$$\Gamma^{(0)}_q = -\epsilon_{\alpha\beta\gamma}C\gamma_5,$$

implicitly transposes the external propagator associated with the $Q$ field.

Figure 1: Feynman diagram for the tree-level vertex of the diquark operator with the quark fields $\bar{Q}$ and $q$. The double line represents the $Q$ field that is transposed and the diquark operator is denoted by $\otimes$. This and all subsequent Feynman diagrams were drawn with JaxoDraw [15].
2 One-Loop Renormalization

Although the diquark operator is gauge dependent, the theory of composite-operator renormalization [16] implies that the diquark operator is multiplicatively renormalizable because there are no lower-dimension operators with the same quantum numbers as (1). The one-loop renormalization of the diquark operator can thus be determined by Fig. 2, which results in the following one-particle irreducible (1PI) Green function for a zero-momentum insertion of \( J_\gamma \) in \( D \)-dimensions (dimensional regularization)

\[
\Gamma^{(1)}_d = \frac{g^2}{4} \lambda^a_{\sigma\alpha} \lambda^a_{\tau\beta} \epsilon^{\sigma\tau\gamma} \nu \frac{1}{\nu^{2\epsilon}} \int \frac{d^Dk}{(2\pi)^D} \left( \gamma^\rho \right)^T \frac{(p+k)^T}{(p+k)^2} C \gamma_\gamma \frac{(p+k)^T}{(p+k)^2} \gamma^\mu \left[ -\frac{g_{\mu\rho}}{k^2} + (1 - \xi) \frac{k_\mu k_\rho}{k^4} \right],
\]

where \( \nu \) is the renormalization scale, the quark mass has been ignored because dimensional regularization is a mass-independent scheme, \( \alpha_s = \frac{g^2}{4\pi} \), colour indices have been explicitly shown for the Gell-Mann matrices \( \lambda^a \), and a covariant gauge with gauge parameter \( \xi \) has been used. Working in normal (or naive) dimensional regularization, we find

\[
\Gamma^{(1)}_d = \frac{8}{3} \left[ -\epsilon^{\alpha\beta\gamma} C \gamma_\gamma \right] \frac{g^2}{4} \frac{1}{\nu^{2\epsilon}} \int \frac{d^Dk}{(2\pi)^D} \gamma^\mu \frac{(p+k)^T}{(p+k)^2} \frac{(p+k)^T}{(p+k)^2} \gamma^\mu \left[ -\frac{g_{\mu\rho}}{k^2} + (1 - \xi) \frac{k_\mu k_\rho}{k^4} \right].
\]

By comparison with the one-loop process determining the renormalization of the scalar current \( J_s \equiv \bar{Q}q \), we see that (4) can be related to the (one-loop) 1PI result for the scalar current \( \Gamma^{(1)}_s \) apart from a numerical factor \( C_d \) representing the ratio of the different colour factors that occur in the two processes

\[
\Gamma^{(1)}_d = \frac{1}{2} \Gamma^{(0)}_d \Gamma^{(1)}_s \equiv C_d \Gamma^{(0)}_d \Gamma^{(1)}_s,
\]

as represented diagrammatically in Fig. 3.

![Figure 2](image-url)  
Figure 2: One-loop Feynman diagram for the renormalization of \( J_\gamma \). As in Fig. 1, the double line represents the (transposed) \( Q \) field and the diquark operator is denoted by \( \otimes \).

![Figure 3](image-url)  
Figure 3: Diagrammatic representation of the relationship (5) between two-point functions with scalar and diquark operator insertions. The scalar operator is denoted by the solid circle.

The renormalized diquark operator \( [J_\gamma]^R \) is defined via the renormalization constant \( Z_d \),

\[
[J_\gamma]^R = Z_d J_\gamma.
\]

Similarly, the well-known renormalization of the scalar operator is

\[
[J_s]^R = Z_m J_s
\]

\footnote{1We are grateful for discussions with John Dixon clarifying this point.} \footnote{2We have chosen to work in normal dimensional regularization (as opposed to, e.g., the 't Hooft-Veltman scheme [17]) because QCD sum-rule analyses of diquarks [12, 14] have used the normal dimensional regularization scheme.}
where $Z_m$ is the quark mass renormalization constant. Using \[5\] it is easy to see that to one-loop order in the minimal-subtraction (MS) and associated schemes
\[Z_d = Z_{2F}^{1/2} Z_m^{1/2},\]
where $Z_{2F}$ is the renormalization constant for the quark fields. Landau gauge ($\xi = 0$) is of particular interest in the QCD sum-rule analysis of diquark currents, because the Schwinger string used for a gauge-invariant formulation of the two-point diquark correlation function vanishes in this gauge \[12\]. Combining the one-loop Landau-gauge result $Z_{2F} = 1$ with \[8\] leads to the one-loop Landau gauge MS-scheme result
\[Z_d = Z_m^{1/2} = 1 + \frac{1}{2} \frac{\alpha}{\pi} \frac{1}{\epsilon},\]
where we use the dimensional regularization convention $D = 4 + 2\epsilon$. Eq. \[9\] agrees with the (one-loop) renormalization and renormalization-group improvement implicitly implemented in Refs. \[12, 14\].

### 3 Two-Loop Renormalization

The two-loop diagrams for the renormalization of the diquark operator are shown in Fig. 4. As in the one-loop analysis and shown in Fig. 3, each diagram is given by a colour factor $C$ and corresponding scalar diagram $\Gamma_{s,i}$ in terms of the renormalization functions $Z_d$. The diagrams for the renormalization of the diquark operator are shown in Fig. 4. As in the one-loop analysis and shown in Fig. 3, each diagram is given by a colour factor $C_d$ multiplying the bare diquark vertex and the equivalent diagram with a scalar current. The divergent parts for each of the two-loop diagrams in Fig. 4 are expressed in Table \[1\] in terms of the corresponding scalar diagram $\Gamma_{s,i}^{(2)}$ in the modified minimal-subtraction (MS) scheme
\[\Gamma_{s,i}^{(2)} = \left( \frac{\alpha_b}{\pi} \right)^2 \left[ A_i + B_i \right], \quad i \in \{1, 2, \ldots, 11\},\]
where $n_f$ is the number of active quark flavours and $\alpha$, $\xi$ are the bare coupling and gauge parameter. A number of the Feynman diagrams are clearly related by the exchange of $Q$ and $q$ fields, and hence Table \[1\] exhibits anticipated symmetries $\Gamma_4 = \Gamma_6$, $\Gamma_7 = \Gamma_8$ and $\Gamma_9 = \Gamma_{11}$. Note that the colour factors $C_d$ that relate the scalar and diquark diagrams are not universally equal to the one-loop result $C_d = 1/2$, implying that one cannot expect the simple pattern of the one-loop result \[9\] to persist at two-loop order. The diagrams that are the exception to the one-loop pattern ($\Gamma_5$ and $\Gamma_{10}$) require multiple applications of colour algebra identities unique to the Feynman rule \[2\]; all other diagrams contain a single application of these identities combined with standard colour algebra factors occurring in the renormalization of the scalar operator \[1\].

The two-loop renormalization procedure first involves the replacement of $\alpha_b$ and $\xi_b$ with their (one-loop) renormalized expressions (see, e.g., Ref. \[21\])
\[Z_\alpha = 1 + \frac{\alpha}{\pi} \left[ \frac{33 - 2n_f}{12\epsilon} \right], \quad \alpha_b = Z_\alpha \alpha;\]
\[Z_\xi = 1 + \frac{\alpha}{\pi} \left[ \frac{4n_f - 39 + 9\xi}{24\epsilon} \right], \quad \xi_b = Z_\xi \xi.\]
in the two-loop 1PI Green function
\[\Gamma_d = \Gamma_d^{(0)} + \Gamma_d^{(1)} + \Gamma_d^{(2)}.\]

For consistency at two-loop level, \[15\] requires inclusion of the finite parts of the one-loop calculation \[5\]
\[\Gamma_5^{(1)} = \frac{1}{3} \frac{\alpha_b}{\pi} \left[ - \frac{4}{3} \frac{\xi_b}{\epsilon} + 2(2 + \xi_b) - L(3 + \xi_b) \right], \quad L = \log \left[ -\frac{p^2}{\nu^2} \right].\]
The renormalization constant $Z_d$ is then constrained by the requirement that it cancel the divergences in
\[Z_d Z_{2F} \left[ \Gamma_d^{(0)} + \Gamma_d^{(1)} + \Gamma_d^{(2)} \right] = 0,\]
where the two-loop MS quark field renormalization constant is \[22\]
\[Z_{2F} = 1 + \frac{\alpha}{\pi} \frac{\xi}{3\epsilon} + \left( \frac{\alpha}{\pi} \right)^2 \left[ \frac{\xi (27 + 17\xi)}{144\epsilon^2} + \frac{201 - 12n_f + 72\xi + 9\xi^2}{288\epsilon} \right].\]

\[3\]In the previous version of this paper the Table \[1\] colour factor for diagram 10 in Fig. 4 was erroneous \[20\].
As a benchmark to ensure accuracy in our calculations in Table 1, we have verified that our results for the scalar diagrams lead to the required two-loop $\overline{\text{MS}}$ result $Z_s = Z_m$ [29].

The final QCD result for the two-loop $\overline{\text{MS}}$ diquark renormalization constant is

$$Z_d = 1 + \frac{\alpha}{\pi} \left[ \frac{3 - \xi}{6\epsilon} \right] + \left( \frac{\alpha}{\pi} \right)^2 \left[ \frac{1}{\epsilon^2} \left( \frac{1545 - 40n_f}{2880} - \frac{\xi}{8} - \frac{\xi^2}{64} \right) + \frac{1}{\epsilon^2} \left( \frac{234 - 12n_f}{288} - \frac{17\xi}{96} - \frac{5\xi^2}{288} \right) \right].$$

(18)

The cancellation of the $L/\epsilon$ terms in $Z_d$ that are generated by (14) provides another consistency check on our calculation. Note that the two-loop Landau gauge result does not uphold the one-loop ($\xi = 0$) pattern $Z_d = Z_m^{1/2}$.

The anomalous dimension for the diquark operator defined by

$$\gamma_d = \nu \frac{dZ_d}{d\nu},$$

(19)

is easily extracted from (18) to obtain the two-loop $\overline{\text{MS}}$ QCD anomalous dimension for the diquark operator.

$$\gamma_d(\alpha) = \gamma_1 \left( \frac{\alpha}{\pi} \right) + \gamma_2 \left( \frac{\alpha}{\pi} \right)^2,$$

(20)

$$\gamma_1 = 1 - \frac{\xi}{3}, \quad \gamma_2 = \frac{1545 - 40n_f}{720} - \frac{\xi}{2} - \frac{\xi^2}{16}.$$

(21)

In the extraction of the anomalous dimension we have verified that the two-loop coefficients of $Z_d$

$$Z_d = 1 + \frac{Z_{d,1}}{\epsilon} + \frac{Z_{d,2}}{\epsilon^2} + \ldots$$

(22)

satisfy the renormalization-group constraint

$$2\alpha \frac{\partial Z_{d,2}}{\partial \alpha} = \left[ \gamma_d(\alpha) - \beta(\alpha)\alpha \frac{\partial}{\partial \alpha} - \delta(\alpha, \xi)\xi \frac{\partial}{\partial \xi} \right] Z_{d,1},$$

(23)
where we are working in the conventions of [21] with the (one-loop) $\beta$ function and anomalous dimension $\delta$ of the gauge parameter given by

$$\beta(\alpha) = \beta_1 \frac{\alpha}{\pi}, \quad \beta_1 = -\frac{11}{2} + \frac{n_f}{3}$$

$$\delta(\alpha, \xi) = \delta_1 \frac{\alpha}{\pi}, \quad \delta_1 = \frac{1}{4} (13 - 3\xi) - \frac{n_f}{3}.$$  

Confirmation of this renormalization-group constraint provides another verification of the accuracy of our results given in Table 1.

### 4 Application and Conclusions

It has previously been noted that at leading-order, the renormalization scale dependence cancels between the QCD perturbative contributions to the diquark decay constants and the $\Delta S = 1$ effective weak Hamiltonian, although there remains some residual scale dependence from non-perturbative terms [14]. As an application of our two-loop results, we can explore this scale dependence at next-to-leading order. Following Ref. [14], we consider the combination

$$c_-(\mu)g_+(\mu)g_+(\mu)$$

where $c_-(\mu)$ represents the renormalization scale dependence of the Wilson coefficient in the $\Delta S = 1$ effective weak Hamiltonian [24] and $g_+(\mu)$ is the scale-dependent scalar diquark decay constant emerging from QCD sum-rules [14]. The renormalization-group (RG) factor arising from $c_-$ is [24]

$$c_-(\mu) \sim \exp \left[ -\int \frac{\gamma_-(\alpha)}{\beta(\alpha)} \frac{d\alpha}{\alpha} \right],$$

### Table 1: Results for the two-loop diagrams in Fig. 4.

| $i$ | $C_i$ | $A_i$ | $B_i$ |
|----|------|------|------|
| 1  | $\frac{n_f(2-L)}{6}$ | $-\frac{n_f}{12}$ | $\frac{n_f}{12}$ |
| 2  | $\frac{(2L-5)(1+\xi^2)}{32}$ | $\frac{1+\xi^2}{32}$ | $\frac{1+\xi^2}{32}$ |
| 3  | $\frac{\xi^2+4\xi^2e_{\pi}-44-L(\xi^2+6\xi_{\pi}-25)}{16}$ | $\frac{25-6\xi_{\pi}-\xi^2}{32}$ |
| 4  | $\frac{(3+\xi_{\pi})(2L(3+\xi_{\pi})-11-5\xi_{\pi})}{18}$ | $\frac{3(3+\xi_{\pi})}{18}$ |
| 5  | $\frac{(3+\xi_{\pi})(2L(3+\xi_{\pi})-11-5\xi_{\pi})}{18}$ | $\frac{3(3+\xi_{\pi})}{18}$ |
| 6  | $\frac{3L(\xi^2+4\xi_{\pi}+3)-5\xi^2_{\pi}-17\xi_{\pi}-24}{16}$ | $\frac{3(\xi^2+4\xi_{\pi}+3)}{32}$ |
| 7  | $\frac{3L(\xi^2+4\xi_{\pi}+3)-5\xi^2_{\pi}-17\xi_{\pi}-24}{16}$ | $\frac{3(\xi^2+4\xi_{\pi}+3)}{32}$ |
| 8  | $\frac{(3+\xi_{\pi})(2L(3+\xi_{\pi})-11-5\xi_{\pi})}{18}$ | $\frac{3(3+\xi_{\pi})}{18}$ |
| 9  | $\frac{3-6\xi_{\pi}-\xi^2}{144}$ | $\frac{0}{144}$ |
| 10 | $\frac{3-6\xi_{\pi}-\xi^2}{144}$ | $\frac{0}{144}$ |
| 11 | $\frac{(3+\xi_{\pi})(2L(3+\xi_{\pi})-11-5\xi_{\pi})}{18}$ | $\frac{3(3+\xi_{\pi})}{18}$ |

The quantity $L = \log(-p^2/\nu^2)$ and the notations for $A_i$ and $B_i$ are defined in Eq. (10).
where in the normal dimensional regularization scheme with $n_f = 3$, the anomalous dimension $\gamma_-(\alpha)$ is

$$\gamma_-(\alpha) = \frac{\gamma_1}{\pi} + \frac{\gamma_2}{\pi} \left( \frac{\alpha}{\pi} \right)^2 \quad (28)$$

$$\tilde{\gamma}_1 = -2 \, , \quad \tilde{\gamma}_2 = -\frac{50}{48} \quad (29)$$

Similarly, the anomalous dimension for the diquark operator leads to the following RG factor for the (scalar) diquark decay constants

$$g_+(\mu)g_+(\mu) \sim \exp \left[ -2 \int \frac{\gamma_2(\alpha)}{\beta(\alpha)} \frac{d\alpha}{\alpha} \right] \quad (30)$$

As mentioned above, QCD sum-rule calculations with diquark currents extract gauge-invariant information from the two-point correlation function through the insertion of a Schwinger string, which becomes trivial for a line geometry in Landau gauge [12]. Thus for applications to RG behaviour of the diquark decay constants, we use (21) with $n_f = 3$ and $\xi = 0$:

$$\gamma_1 = 1 \, , \quad \gamma_2 = \frac{95}{48} \quad (31)$$

The resulting RG behaviour of (26) is

$$c_-(\mu)g_+(\mu)g_+(\mu) \sim \exp \left[ \int \frac{4}{9} \frac{1 + \frac{\gamma_1}{\gamma_1} \frac{\alpha}{\pi}}{1 + \frac{\gamma_2}{\beta_1} \frac{\alpha}{\pi}} \frac{d\alpha}{\alpha} \right] \exp \left[ - \int \frac{4}{9} \frac{1 + \frac{\gamma_1}{\gamma_1} \frac{\alpha}{\pi}}{1 + \frac{\gamma_2}{\beta_1} \frac{\alpha}{\pi}} \frac{d\alpha}{\alpha} \right] = 1 - \frac{35 \, \alpha(\mu)}{54 \, \pi} \quad (32)$$

Thus the leading-order cancellation of scale dependence in (26) for the perturbative contributions to $g_+$ does not persist to second order. However, the residual scale dependence associated with (32), which decreases with increasing $\alpha(\mu)$, does have the right qualitative behaviour to counter the residual scale dependence encountered in Ref. [14]. A more detailed analysis of the residual scale dependence is beyond the scope of this paper because it would require a full next-order sum-rule analysis of the diquark decay constants.

In conclusion, we have determined the MS renormalization constant and associated anomalous dimension for the scalar diquark operator at two-loop order in QCD in an arbitrary covariant gauge for normal dimensional regularization. This result enables future QCD sum-rule studies of diquarks to higher-orders in perturbation theory. For example, the divergent terms in the diquark renormalization constant [18] combined with lower-loop $O(\epsilon)$ and $O(\epsilon^2)$ terms generate finite parts corresponding to renormalization-induced physical contributions to the diquark correlation function. Furthermore, the anomalous dimension of the diquark operator appearing in the renormalization-group equation governing scale dependence of the diquark correlation function is an essential feature of QCD Laplace sum-rule analyses [25]. Given the relative size of the one- and two-loop terms in (18) and (31), these renormalization-induced and anomalous dimension effects could be significant in higher-loop extensions of [14].

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*Note that we have converted the expressions in [24] into our conventions.
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