CURRENT s - QUARK MASS CORRECTIONS TO THE FORM FACTORS OF D - MESON SEMILEPTONIC DECAYS

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Abstract

The infinite mass effective theory, when a heavy quark mass tends to infinity, and Chiral perturbation theory at the quark level, based on the extended Nambu – Jona – Lasinio model with linear realization of chiral $U(3) \times U(3)$ symmetry, are applied to the calculation of current $s$ – quark mass corrections to the form factors of the $D \to \bar{K} e^+ \nu_e$ and $D \to \bar{K}^* e^+ \nu_e$ decays. These corrections turn out to be quite significant, of the order of $7-20\%$. The theoretical results are compared with experimental data.

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1 Introduction

In our recent publications [1,2] we calculated in the chiral limit the form factors of the semileptonic $D \to \bar{K}^* e^+ \nu_e$ and $D \to \bar{K} e^+ \nu_e$ decays. For the description of $D$-mesons we applied the infinite mass effective theory (IMET) [3,4], when the $c$-quark mass $M_c$ tends to infinity. IMET the long-distance physics we describe within Chiral perturbation theory at the quark level (CHPT) [5], based on the extended Nambu – Jona – Lasinio (ENJL) model with linear realization of chiral $U(3) \times U(3)$ symmetry [6].

In this paper we apply IMET and (CHPT)$_q$ to the calculation of the fine structure of the form factors of the $D \to \bar{K} e^+ \nu_e$ and $D \to \bar{K}^* e^+ \nu_e$ decays at the first order in current $s$–quark mass expansion. Within IMET and (CHPT)$_q$ the first order current – quark – mass corrections to the mass spectra of charmed pseudoscalar and vector mesons and charmed pseudoscalar – meson leptonic constants have been calculated in [7]. The amplitude of the $D \to h e^+ \nu_e$ decay can be determined as follows

$$M(D \to h e^+ \nu_e) = -\frac{G_F}{\sqrt{2}} V_{cs}^* < h(Q) | \bar{s}(0) \gamma_\mu (1 - \gamma_5) c(0) | D(p) > \ell^\mu ,$$

where $h = \bar{K}$ or $\bar{K}^*$, $G_F = 1.166 \times 10^{-5}$ GeV$^{-2}$ is Fermi weak constant, $|V_{cs}| = 0.975$ is the CKM–mixing matrix element, $s(0)$ and $c(0)$ are the $s$– and $c$–current quark fields with $N$ colour degrees of freedom each, and $\ell^\mu = \bar{u}(k_\nu_e) \gamma^\mu (1 - \gamma_5) v(k_{e^+})$ is the weak leptonic current.

We shall seek the hadronic matrix element

$$M_\mu (D \to h) = < h(Q) | \bar{s}(0) \gamma_\mu (1 - \gamma_5) c(0) | D(p) >$$

in the form of an expansion in powers of the current $s$–quark mass upto first order terms

$$M_\mu (D \to h) = M_\mu^{(0)} (D \to h) + M_\mu^{(1)} (D \to h).$$

Here we have denoted

$$M_\mu^{(0)} (D \to h) = < h(Q) | \bar{s}(0) \gamma_\mu (1 - \gamma_5) c(0) | D(p) >_{\text{ch.l.}}$$

$$M_\mu^{(1)} (D \to h) =$$

$$= -i m_{0s} \int d^4 x < h(Q) | T(\bar{s}(x) s(x) \bar{s}(0) \gamma_\mu (1 - \gamma_5) c(0)) | D(p) >_{\text{ch.l.}}.$$

The matrix element $M_\mu^{(0)} (D \to h)$ describes the $D \to h$ transition calculated in the chiral limit (ch.l.) while $M_\mu^{(1)} (D \to h)$ is the first order correction in the current $s$–quark mass expansion. The matrix elements $M_\mu^{(0)} (D \to h)$ for $h = \bar{K}^*$ and $\bar{K}$ have been calculated in [1,2]. In this paper we shall calculate $M_\mu^{(1)} (D \to h)$.

In accordance to the procedure expounded in [1,2,7] we reduce (6) to the expression

$$M_\mu^{(1)} (D \to h) = g_D m_{0s} i \int d^4 x \int_{-\infty}^{\infty} d z_0 \theta (-z_0) \times$$

$$\times < h(Q) | T(\bar{s}(x) s(x) s(0) \gamma_\mu \left( \frac{1 + \gamma^0}{2} \right) \gamma^5 q(z_0, 0)) | 0 >_{\text{ch.l.}}.$$
obtained in leading order in the large $N$ and $M_c$ expansion, $q = u$ or $d$ for $D^0$ or $D^+$, respectively. The coupling constant $g_D$ has been calculated in [8]

$$g_D = \frac{2\sqrt{2}\pi}{\sqrt{N}} \left( \frac{M_D^2}{M_e \bar{v}'} \right)^{1/2}, \quad (7)$$

where $\bar{v}' = 4\Lambda = 2.66$ GeV and $\Lambda$ is the cut-off in 3-dimensional Euclidean momentum space. $\Lambda$ is connected to the scale of spontaneous breaking of chiral symmetry (SBCS) $\Lambda_\chi$ by the relation $\Lambda = \Lambda_\chi/\sqrt{2} = 0.66$ GeV at $\Lambda_\chi = 0.94$ GeV [5].

The r.h.s. of (8) involves only the light – quark fields. Therefore for the evaluation of (8) we can apply (CHPT)$_q$ [5,7]. Since the leading order of the r.h.s. of (8) in current – quark – mass expansion is fixed by the factor $m_{0s}$, so the matrix element $< h (Q)| T(\ldots)| 0 >$ has to be calculated in the chiral limit (ch.l.).

By applying the formulas of quark conversion (Ivanov [5]) we can present the matrix element $M^{(1)}_\mu (D \rightarrow h)$ in terms of constituent – quark – loop diagrams [7]. The momentum representation of these diagrams reads

$$M^{(1)}_\mu (D \rightarrow h) = i m_{0s} g_D g_h \frac{\bar{v}}{4m} \left( \frac{N}{16\pi^2} \right) \int \frac{d^4k}{\pi^2 i} \text{tr} \left[ \frac{1}{m-k} \bar{\Gamma}_h \times \gamma_\mu (1-\gamma^5) \left( \frac{1+\bar{v}}{2} \right) \gamma^5 \frac{1}{k \cdot v + i0} \right]. \quad (8)$$

The appearance of the factor $\bar{v}/4m$ is due to the contribution of the diagram with the light scalar $\sigma_s$ - meson exchange [8]. Here $\bar{v} = -<0|\bar{q}q|0>/F_0^2 = 1.92$ GeV, $F_0 = 0.092$ GeV and $m = 0.33$ GeV are the PCAC constant of light pseudoscalar mesons and the constituent quark mass calculated in the chiral limit [5]. The coupling constants $g_h$ describe the interaction between light constituent quarks and light mesons $\bar{K}$ and $\bar{K}^*$, that is $g_K = 2\pi/\sqrt{N}$ and $g_{K^*} = \pi \sqrt{6}/\sqrt{N}$ [7,8] such that $g_{K^*}/g_K = \sqrt{3}/2$ [5]. $\Gamma_h$ is either $\Gamma_K = i\gamma^5$ or $\Gamma_{K^*} = \gamma^5 \epsilon^{\nu\epsilon\nu}(Q)$ depending on whether $h = \bar{K}$ or $h = \bar{K}^*$. Now we can proceed to the calculation of the current $s$ - quark mass corrections to the form factors.

2. The $D \rightarrow \bar{K} e^+ \nu_e$ Decay

For the $D \rightarrow \bar{K}$ the matrix element $M^{(1)}_\mu (D \rightarrow \bar{K}^+)$ reads

$$M^{(1)}_\mu (D \rightarrow \bar{K}) = i m_{0s} g_D \frac{2\pi}{\sqrt{N}} \frac{\bar{v}}{4m} \left( \frac{N}{16\pi^2} \right) \int \frac{d^4k}{\pi^2 i} \text{tr} \left[ \frac{1}{m-k} i\gamma^5 \times \gamma_\mu (1-\gamma^5) \left( \frac{1+\bar{v}}{2} \right) \gamma^5 \frac{1}{k \cdot v + i0} \right] = \quad (9)$$

$$= m_{0s} g_D \frac{\sqrt{N}}{4\pi} \int \frac{d^4k}{\pi^2 i} \frac{k_\mu}{[m^2-k^2-i0][m^2-(k+Q)^2-i0]} \frac{1}{k \cdot v + i0} + \ldots .$$

Following [9,10] we have kept only divergent contributions. The dots denote the contributions of convergent integrals. The integration over $k$ gives [1]
\[
\int \frac{d^4k}{\pi^2i} \frac{k_\mu}{[m^2 - k^2 - i0][m^2 - (k + Q)^2 - i0]} \frac{1}{k \cdot v + i0} = v_\mu 2 \ln \left(1 + \frac{\bar{v}'}{4Q_0}\right) + Q_\mu \frac{2}{Q_0} \left[1 - \ln \left(1 + \frac{\bar{v}'}{4Q_0}\right)\right]
\]

(10)

where \( Q_0 = (M_D^2 - q^2)/2M_D \) is the energy of the massless \( K \) meson in the rest frame of the \( D \) meson. The appearance of the \( q^2 \) dependence is due to the \( q^2 \) dependence of \( Q_0 \). The matrix element \( M^{(1)}_\mu (D \rightarrow \bar{K}) \) can be expressed in terms of two form factors

\[
M^{(1)}_\mu (D \rightarrow \bar{K}) = f^{(1)}_+ (q^2) (p + Q)_\mu + f^{(1)}_- (q^2) , (p - Q)_\mu
\]

(11)

where

\[
\begin{align*}
f^{(1)}_+ (q^2) &= \frac{m_{0s}}{M_*} \frac{\bar{v}}{4m} \frac{M_D^2}{M_D^2 - q^2} \left[1 - \frac{M_D^2 + q^2}{2M_D^2} \ln \left(1 + \frac{M_D^2}{M_D^2 - q^2}\right)\right], \\
f^{(1)}_- (q^2) &= -\frac{m_{0s}}{M_*} \frac{\bar{v}}{4m} \frac{M_D^2}{M_D^2 - q^2} \left[1 - \frac{3M_D^2 - q^2}{2M_D^2} \ln \left(1 + \frac{M_D^2}{M_D^2 - q^2}\right)\right].
\end{align*}
\]

(12)

Here we have denoted \( 2M_* = \sqrt{2M_D \bar{v}'} \). It should be stressed that the formulae (12) are valid in the physical region only, i.e. \( 0 \leq q^2 \leq (M_D - m_K)^2 \). At \( q^2 = 0 \) the form factors \( f^{(1)}_+ (0) \) and \( f^{(1)}_- (0) \) read

\[
\begin{align*}
f^{(1)}_+ (0) &= \frac{m_{0s}}{M_*} \frac{\bar{v}}{4m} \left[1 - \frac{1}{2} \ln \left(1 + \frac{M_D^2}{M_D^2}\right)\right] = 0.09, \\
f^{(1)}_- (0) &= -\frac{m_{0s}}{M_*} \frac{\bar{v}}{4m} \left[1 - \frac{3}{2} \ln \left(1 + \frac{M_D^2}{M_D^2}\right)\right] = -0.02.
\end{align*}
\]

(13)

In the chiral limit the quantity \( f_+ (0) \) has been calculated in [8] (see also [2])

\[
f^{(0)}_+ (0) = \frac{1}{\sqrt{2}} \left(\frac{\bar{v}'}{2M_c}\right)^{1/2} = 0.6.
\]

(14)

The numerical value is estimated at the equality \( M_c = M_D = 1.86 \) GeV accepted in our approach [9]. By adding the current \( s \) quark mass correction (13) we get the total value of \( f_+ (0) \)

\[
f_+ (0) = f^{(0)}_+ (0) + f^{(1)}_+ (0) = 0.69
\]

(15)

which is good compared with the experimental data \( |f_+ (0)|_{\exp} = 0.7 \pm 0.1 \) [11]. Our result \( f_+ (0) = 0.69 \) agrees well with the theoretical prediction by Dominguez and Paver [12] obtained within the QCD sum rule approach. We find the current \( s \)-quark mass correction to be about 15%.
3 The $D \to \bar{K}^* e^+ \nu_e$ decay

The matrix element $M^{(1)}_\mu (D \to \bar{K}^*)$ can be expressed in terms of four form factors [1]

$$M^{(1)}_\mu (D \to \bar{K}^*) = \begin{aligned} &i a^{(1)}_1 (q^2) e^*_\nu (Q^2) - i a^{(1)}_2 (q^2) (e^*(Q) \cdot p) (p + Q)_\mu - \\ &i a^{(1)}_3 (q^2) (e^*(Q) \cdot p) (p - Q)_\mu - \\ &2 b^{(1)} (q^2) \varepsilon_{\mu \nu \alpha \beta} e^{\nu \nu} (Q) p^\alpha Q^\beta, \end{aligned} \quad (16)$$

In order to obtain the form factors $a^{(1)}_i (q^2)$ ($i = 1, 2, 3$) and $b^{(1)} (q^2)$ we have to calculate the following momentum integral

$$\mathcal{M}_{\mu \nu} = \int \frac{d^4k}{\pi^2 i} \text{tr} \left[ \frac{1}{m - k} \gamma_\nu \frac{1}{m - Q - k} \frac{1}{m - Q - k} \times \right.$$  
$$\times \gamma_\mu (1 - \gamma^5) \left( \frac{1 + \hat{v}}{2} \right) \gamma^5 \frac{1}{k \cdot \hat{v} + i 0} \right]. \quad (17)$$

By keeping only divergent contributions [9,10] and using the integrals given in the Appendix of [1], we get

$$\mathcal{M}_{\mu \nu} = - 4 \ln \left( \frac{\bar{v}'}{4 m} \right) g_{\mu \nu} + \frac{8}{M^2_D - q^2} \left[ 1 - \ln \left( 1 + \frac{M^2}{M^2_D - q^2} \right) \right] Q_\mu p_\nu -$$
$$- \frac{8 i}{M^2_D - q^2} \left[ 1 - \ln \left( 1 + \frac{M^2}{M^2_D - q^2} \right) \right] \varepsilon_{\mu \nu \alpha \beta} p^\alpha Q^\beta. \quad (18)$$

The appearance of the $q^2$ dependence is due to the quantity $Q_0 = (M^2_D - q^2)/2 M_D$, being the energy of the massless $\bar{K}^*$ meson in the rest frame of the $D$ meson. The neglect of the $\bar{K}^*$ meson mass in the r.h.s. of (17) is in accordance with the prescription of (CHPT)$_q$ which incorporates the Vector Dominance approach [5,13], admitting the smooth dependence of low-energy hadronic matrix elements on the masses of low-lying vector mesons ($\rho$, $\omega$, $\varphi$, $K^*$) [1,13,14].

By using (17), one can calculate the following chiral corrections to the form factors of the $D \to \bar{K}^*$ transition

$$a^{(1)}_1 (q^2) = \frac{\sqrt{3}}{2} \frac{m_0 s}{M^2} \frac{\bar{v}}{4 m} M_D \ln \left( \frac{\bar{v}'}{4 m} \right)$$
$$a^{(1)}_2 (q^2) = - a^{(1)}_3 (q^2) = b^{(1)} (q^2)$$
$$b^{(1)} (q^2) = \frac{\sqrt{3}}{2} \frac{m_0 s}{M^2} \frac{\bar{v}}{4 m} \frac{M_D}{M^2_D - q^2} \left[ 1 - \ln \left( 1 + \frac{M^2}{M^2_D - q^2} \right) \right]. \quad (19)$$

In the chiral limit the form factors of the $D \to \bar{K}^*$ transition have been calculated in [1]
These numerical results obtained by taking into account the first order current mass corrections confirm the results found in [1]. It is because in [1] we expressed the form factors of the $D \rightarrow \bar{K}^*$ transition in terms of the form factor of the $D \rightarrow K$ transition $f_{+}(0)$. There for the numerical estimate we used the value $f_{+}(0) = 0.7$, which we obtained in present paper only at the first order in current $s$–quark mass.

\begin{align}
a_1^{(0)}(q^2) &= \sqrt{\frac{3}{8}} M_* \\
a_2^{(0)}(q^2) &= \sqrt{\frac{3}{8}} M_* \left[ \frac{q^2}{M_D^2 - q^2} + \frac{M_D^2 - q^2}{M_*^2} \left( 1 - \frac{2mM_D}{M_D^2 - q^2} \right) \ln \left( 1 + \frac{M_*^2}{M_D^2 - q^2} \right) \right], \\
a_3^{(0)}(q^2) &= -\sqrt{\frac{3}{8}} M_* \left[ \frac{2M_D^2 - q^2}{M_D^2 - q^2} - \frac{M_D^2 - q^2}{M_*^2} \left( 1 - \frac{2mM_D}{M_D^2 - q^2} \right) \ln \left( 1 + \frac{M_*^2}{M_D^2 - q^2} \right) \right], \\
b^{(0)}(q^2) &= \sqrt{\frac{3}{8}} \frac{1}{M_*} \ln \left( 1 + \frac{M_*^2}{M_D^2 - q^2} \right). \tag{20}
\end{align}

Here we have used the relation (14). The numerical values of the form factors at $q^2 = 0$ read

\begin{align}
a_1(0) &= a_1^{(0)}(0) + a_1^{(1)}(0) = 0.96 + 0.14 = 1.10 \text{ (GeV)}, \\
a_2(0) &= a_2^{(0)}(0) + a_2^{(1)}(0) = 0.14 + 0.03 = 0.17 \text{ (GeV)}^{-1}, \\
a_3(0) &= a_3^{(0)}(0) + a_3^{(1)}(0) = -0.42 - 0.03 = -0.45 \text{ (GeV)}^{-1}, \\
b(0) &= b^{(0)}(0) + b^{(1)}(0) = 0.21 + 0.03 = 0.24 \text{ (GeV)}^{-1}. \tag{21}
\end{align}

One sees that the first order current $s$–quark mass corrections are between 7 and 20%. The form factors $a_i(q^2)$ ($i = 1, 2, 3$) and $b(q^2)$ are connected with the standard form factors $A_i(q^2)$ ($i = 1, 2, 3$) and $V(q^2)$ via the relations [1]

\begin{align}
A_1(q^2) &= \frac{1}{M_D + M_{K^*}} a_1(q^2)|_{q^2=0} = 0.40 \\
A_2(q^2) &= (M_D + M_{K^*}) a_2(q^2)|_{q^2=0} = 0.47 \\
A_3(q^2) &= (M_D + M_{K^*}) a_3(q^2)|_{q^2=0} = -1.24 \\
V(q^2) &= (M_D + M_{K^*}) b(q^2)|_{q^2=0} = 0.66, \tag{22}
\end{align}

where $M_{K^*} = 0.89$ GeV is the mass of the $\bar{K}^*$–meson [11]. The theoretical values compare reasonably with recent experimental data [15]

\begin{align}
A_1(0)_{exp} &= 0.46 \pm 0.05 \pm 0.05, \\
A_2(0)_{exp} &= 0.38 \pm 0.1 \pm 0.07, \\
V(0)_{exp} &= 0.92 \pm 0.1 \pm 0.12. \tag{23}
\end{align}

These numerical results obtained by taking into account the first order current $s$–quark mass corrections confirm the results found in [1]. It is because in [1] we expressed the form factors of the $D \rightarrow \bar{K}^*$ transition in terms of the form factor of the $D \rightarrow K$ transition $f_{+}(0)$. There for the numerical estimate we used the value $f_{+}(0) = 0.7$, which we obtained in present paper only at the first order in current $s$–quark mass.
expansion (13). Recall that in the chiral limit we have \( f_+(0) = 0.6 \). This overlap of results underscores the self-consistency of the current \( s \)-quark mass corrections to the form factors of the transitions \( D \to \bar{K}^* \) and \( D \to \bar{K} \) calculated within IMET and (CHPT)\( _q \).

4 Conclusion

We have applied IMET and (CHPT)\( _q \) for the computation of the current \( s \)-quark mass corrections to the form factors of the semileptonic decays of the non-strange charmed \( D \)-mesons, \( D \to \bar{K} e^+ \nu_e \) and \( D \to \bar{K}^* e^+ \nu_e \). We have obtained non-zero contributions for the first order corrections in current \( s \)-quark mass expansion to the form factors of the \( D \to \bar{K} e^+ \nu_e \) decays. This result contradicts the Ademollo–Gato theorem for the form factors of the semileptonic decays of \( K \)-mesons [16]. Within (CHPT)\( _q \) the Ademollo–Gato theorem has been analyzed in [17]. The observed contradiction can be explained as an effect of IMET. Indeed IMET is based on the infinite limit \( M_c \to \infty \) which violates chiral \( SU(4) \times SU(4) \) symmetry, a necessary condition for the validity of the Ademollo–Gato theorem for the \( D \to \bar{K} e^+ \nu_e \) decays. The current \( s \)-quark mass corrections to the form factors of the \( D \to \bar{K}^* e^+ \nu_e \) decays are consistent with the corrections calculated for the form factors of the \( D \to \bar{K} e^+ \nu_e \) decays.

Note that we have kept to the first order corrections in current \( s \)-quark mass expansion calculated at the tree–meson level. Of course, the one–meson–loop corrections can be taken into account too. The consistent procedure for meson–loop chiral corrections within (CHPT)\( _q \) has been developed in Ref.[5]). This procedure can also be applied to charmed meson physics.

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