Heavy Quark Effective Fields as Operator-valued Distributions *

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Abstract

We look at effective fields defined in the heavy-quark effective theory as operator-valued generalized functions on Minkowski space-time to be averaged with physically suitable smoothing functions (Gaussians of typical width the hadronic size) leading to operator-valued distributions in Hilbert space. One-heavy-quark states are thus represented by normalizable wave packets displaying a particle-like behaviour at the characteristic hadronic time scale $\Lambda_{QCD}^{-1}$. We examine some consequences relative to the average kinetic energy of heavy quarks in hadrons to avoid inconsistencies within this formalism.

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1 Introduction

In recent years the physics of hadrons containing a heavy quark has been analyzed by means of an effective theory for the strong interaction (HQET) based on an expansion in the inverse heavy quark mass \( m_Q^{-1} \). This approach shows the existence of new symmetries at leading order, dealing systematically with the corrections arising from higher order terms, and allowing numerous fruitful phenomenological applications among which the extraction of the Cabibbo-Kobayashi-Maskawa matrix element \( |V_{cb}| \). Nevertheless, such a powerful framework still suffers from some internal ambiguities. Specifically, typical parameters like the average kinetic energy of the heavy quark in a hadron (needed for a precise determination of \( |V_{cb}| \) from \( B \) decays) are under intense discussion, and whose precise physical definiteness has been questioned (see \([2, 3]\) and references therein).

In constructing HQET a cornerstone is the observation that massive quarks in heavy-light hadrons exchange momenta with the light degrees of freedom typically of order \( \Lambda_{QCD} \), the characteristic scale of the strong interaction, much smaller than the heavy quark mass \( m_Q \). Therefore and much contrary to light quarks, a heavy quark should propagate along a particle-like trajectory in space - with an almost constant velocity - until the decay of the hadron takes place. This physical motivation (with vivid analogies in early literature for the heavy quark such as baseball \([4]\) or even cannonball \([5]\)) shows that HQET resembles more a classical approximation (e.g. no heavy-quark pair production is allowed) than a non-relativistic one, at least as far as the long-distance aspects of the theory are concerned.

These remarks suggest the split of the heavy quark four-momentum into two pieces:

\[
p_Q = m_Q v + k_Q
\]

where the first term in the rhs represents the large mechanical part as \( v \) is the hadron’s velocity (although this is not necessary because of the so-called reparametrization invariance of the theory \([\tilde{4}]\) and \( k_Q \) denotes the “residual momentum” arising from the predominantly soft interaction of the heavy quark with gluons.

Squaring \( p_Q \) we shall write

\[
p_Q^2 = m_Q^2 + \Delta
\]

with \( \Delta \) denoting a measure of the off-shellness of the heavy quark, i.e.

\[
\Delta = p_Q^2 - m_Q^2 = 2m_Q v \cdot k_Q + k_Q^2
\]

where in the limit \( m_Q \to \infty, \Delta/m_Q^2 \to 0 \).

As a simple but illustrative exercise, let us suppose at a given instant a heavy quark on-shell (with a particular choice for \( m_Q \)) and at rest in the hadron rest frame, absorbing a soft on-shell gluon whose four-momentum is \((\Lambda_{QCD}, \vec{\Lambda}_{QCD})\). Looking at the interaction as a real process, energy-momentum conservation implies that after the absorption the heavy quark mass \( m_Q^* \) would shift according to

\[
m_Q^* = m_Q + 2 \Lambda_{QCD}
\]

that is

\[
m_Q^* \approx m_Q + \Lambda_{QCD}
\]

while its kinetic energy should be \( \simeq \Lambda_{QCD}^2/2m_Q \). Notice, however, that the absorbed energy by the heavy quark actually is \( \Lambda_{QCD} \). (A similar reasoning applies to the emission of a soft gluon.)
According to the energy-time uncertainty principle it is possible to interpret that during a characteristic time interval $\Delta t$ given by

$$\Delta t \simeq \frac{3}{2} \Lambda_{QCD}^{-1}$$

the heavy quark propagates virtually free over a distance of at least its Compton wavelength, after which its (total) energy is changed by an amount $\simeq \Lambda_{QCD}$, varying both its mass and kinetic energy accordingly. In fact, this assumption is on the grounds for the use of the static approximation in computing the quark-antiquark (QCD-inspired) potential for heavy quarkonium, where $\Lambda_{QCD}^{-1} \approx 10^{-23}$s is the typical time required by the gluonic fields to adjust themselves to the movement of the heavy quarks. Furthermore, $\Lambda_{QCD}^{-1}$ is as well the characteristic time scale for building up heavy-light or heavy-heavy hadrons subsequently to a heavy-flavour production process.

As a consequence of many successive soft interactions between the heavy quark and the light constituents the former acquires a mean kinetic energy while its mass “fluctuates” about $m_Q$ with typical deviation $\Lambda_{QCD}$ (that is $|\Delta| \simeq 2m_Q\Lambda_{QCD}$). Certainly this is an idealization since we have ignored the existence of short-distance physics. Nonetheless, hard gluons play little role in the structure of low-lying hadronic states and, in particular, the HQET operator of the non-relativistic kinetic energy is not multiplicatively renormalized as a consequence of the reparametrization invariance already mentioned. We shall turn to this point in Section 3, however.

All the above considerations suggest writing a plane-wave Fourier expansion for the fermionic field $Q_v(x)$ (positive frequencies) corresponding to an almost on-shell heavy quark moving inside a hadron with four-velocity $v$

$$Q_v(x) = \int d\mu(p) \sum_r b_r(\vec{p}) u_r(\vec{p}) e^{-ip \cdot x}$$

where we have required that all the Fourier components satisfy the same off-shellness condition imposed by Eq. (2), that is

$$p^2 = m_Q^2 + \Delta$$

Hence $d\mu(p)$ stands for the invariant measure although $p$ does not exactly lay on the upper mass-shell hyperboloid (corresponding to the “assumed” $m_Q$ value) since $p^0 = \sqrt{m_Q^2 + \Delta + \vec{p}^2}$. On the other hand, $b_r(\vec{p})$ is the usual annihilation particle operator, for a heavy quark in this case.

Let us redefine the momenta of the Fourier components in (4) according to HQET as

$$p = m_Q v + k$$

where $k$ represents the Fourier residual four-momentum. Let us stress that in our notation the latter must not be confused with the residual four-momentum $k_Q$ characterizing a particular hadronic state as an expectation value (see Ref. [9]).

On the other hand, velocity-dependent effective fields in HQET are introduced as:

$$h_v(x) = e^{im_Q v \cdot x} P_+ Q_v(x)$$

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1. We keep all numerical factors - they cancel out at the end of the calculation anyway - although, of course, only the order of magnitude is significant. Factor 3 is associated with the fact that we are considering three-dimensional motion.

2. Heavy anti-quarks would be treated in the same way. We also omit any reference to colour and flavour.
\[ H_v(x) = e^{i m_Q v \cdot x} P_- Q_v(x) \]  

where the projector operators \( P_{\pm} = (1 \pm \phi)/2 \) leave only the upper (lower) components in the hadron rest frame. As is well-known, the “small” components \( H_v \) can be removed by means of the (classical) field equations yielding the following long-distance Lagrangian at order \( 1/m_Q \) expressed in standard notation as

\[ L_{\text{eff}} = \overline{\eta} v \cdot D h_v + \frac{1}{2 m_Q} \overline{\eta} v(iD_\perp)^2 h_v + \frac{g}{4 m_Q} \overline{\eta} \sigma_{\alpha\beta} G^{\alpha\beta} h_v \]  

where the second term in the rhs stands for the kinetic energy arising from the off-shell residual motion of the heavy quark. The third term describes the chromomagnetic interaction of the heavy quark with the gluon field \( \Phi \).

As far as the heavy quark could be considered (quasi)-free during the tiny time interval \( \Delta \), from comparison between \( (4) \) and \( (7) \) the \( h_v \) field can be written as

\[ h_v(x) = 1 + \frac{1 + \phi}{2} \int d\mu(\vec{k}) \sum b_r(\vec{v}, \vec{k}) u_r(\vec{k}) e^{-i \vec{k} \cdot \vec{x}} \]  

where \( b^*_r(\vec{v}, \vec{k}) \) denotes the annihilation operator for a heavy quark with residual spatial momentum \( \vec{k} \) in a hadron moving with four-velocity \( v \); \( d\mu(\vec{k}) \) stands again for the invariant measure now expressed in terms of the Fourier residual momentum. (The zero component of the residual four-momentum \( k \) in the hadron rest frame must satisfy: \( k^0 = \sqrt{m_Q^2 + k^2 + \Delta} - m_Q \) corresponding to the small off-shellness condition derived from \( (5) \) \( k^2 + 2 m_Q v \cdot k = \Delta \), for positive frequencies.)

Observe that, as written in Eq. \( (10) \), the \( \overline{\eta}(x) \) field creates a heavy quark at point \( x \) of space-time as a superposition of single-particle states with momenta ranging over the entire \( \vec{k} \) domain. Nevertheless, one may introduce an ultraviolet cut-off eliminating those components with residual momenta greater than the heavy quark mass \( m_Q \). This corresponds to ignore those Fourier components whose wavelengths are smaller than the Compton wavelength of the heavy quark, amounting to a position uncertainty of this order as is well-known from Quantum Theory \[ \text{[1]} \].

Moreover, in constructing the low energy effective theory for heavy quarks one has still to add a further restriction to the Fourier expansion of the fields, limiting the range of the \( \vec{k} \) components to values of the order of \( \Lambda_{QCD} \). Note that this condition amounts to a new constraint not completely equivalent to the small virtuality already imposed by means of Eq. \( (5) \). Indeed, the smallness of \( \vec{k} \) implies the almost on-shellness condition but the converse is not necessarily true. In fact, removing those components with large residual spatial momentum in the Fourier expansion of the fields is similar to integrating away the high-velocity states in the functional integral formulation of Ref. \[ \text{[12]} \].

Short-distance effects involving large virtual momenta \( \vec{k} \) can be included in the effective theory in a perturbative way using renormalization group techniques in a procedure called matching \[ \text{[1]} \].

\[ ^3 \text{In one-particle relativistic theory this amounts to a zigzag motion traditionally known as zitterbewegung \[ \text{[1]} \]. This oscillatory motion, involving negative-energy states, present even in the hypothetical case of a free quark, should be distinguished from the residual motion due to the soft interaction with the light constituents of the hadron.} \]

\[ ^4 \text{As stressed in \[ \text{[2]} \] the normalization scale } \mu \text{ should lie somewhere in the range } \Lambda_{QCD} << \mu << m_Q \text{ in the context of the Wilsonian (operator product expansion) approach to the heavy-quark theory. In practice } \mu \simeq \text{ several units } \times \Lambda_{QCD} \]
2 Effective Fields as Operator-valued Distributions

In HQET, hadron states are usually identified with the eigenstates of the leading term in the Lagrangian (9) (corresponding to \( m_Q \to \infty \)) supplemented with the standard QCD Lagrangian for the light quarks and gluons \([1]\). However, as emphasized in \([13]\), heavy-quark states in Hilbert space should be associated (to some arbitrariness) with a certain velocity \( v \) and residual momentum \( k \). Single-particle (Fock) states are thus obtained by letting the creation operators \( b^+_r(\vec{v}, \vec{k}) \) act on the vacuum. In the infinite mass limit (or in practice for \( m_Q \gg \Lambda_{QCD} \)) any one-particle state with a given value of \( p_Q = m_Q v + k \) could be used to represent the state \( |p> \) of the full theory. In this work, however, we shall consider a superposition of states (wave packet) for describing a heavy quark localized in a definite region of space. In fact, there is a framework to implement in a natural way the Fourier content of the heavy quark effective fields based on the concept of operator-valued distributions.

Since the beginning of quantum field theory, it was pointed out by Heisenberg that the measurement of a field quantity at a space-time point (though itself a useful mathematical entity) must be impossible (see Ref. \([14]\) for an interesting historical account of quantum field theory). Indeed, in experiments the field strength is always measured not at a mathematical point. In this respect Bohr and Rosenfeld \([15]\) early argued that the real meaning of fields was related to average values in space-time regions.

To become properly defined operators in Hilbert space, fields (\( \Phi \)) have to be smoothed by suitable functions of sufficiently regular behaviour on Minkowski space:

\[
\Phi(f) = \int \ d^4x \ \Phi(x) \ f(x)
\]

where \( f \) belongs to a “test” function space, usually taken as the space \( S \) of infinitely often differentiable functions decreasing as well as their derivatives faster than any power as \( x \) moves to infinity \([16]\). The resulting functional \([17]\) is an operator-valued (tempered) distribution, which acting on the vacuum of Hilbert space of states generates normalizable one-particle states or wave packets.

On the other hand, one usually has to deal with several fields, each of which may have tensor or spinor components. Thus Eq. (11) must be generalized as

\[
\Phi(f) = \sum_{r=1}^{n} \int \ \Phi_{r\alpha}(x) \ f_{r\alpha}^\alpha(x) \ d^4x
\]

where the index \( r \) denotes the type of field and \( \alpha \) stands for each component.

In this paper our intention is not to present a rigorous development of HQET based on the framework of the general quantum field theory \([16]\). Rather we shall make use of some ideas underlying the concept of smeared fields to gain insight into several aspects of the effective theory for heavy quarks in hadrons. Indeed, working in the hadron rest frame and due to the projectors \( P_{\pm} \) only the two upper components of the spinor field are not null. Thus we shall dispense with the indices \( r \) and \( \alpha \) of the the smoothing function which becomes just a c-number. (In general, test functions should display a tensor or spinor character in order to fulfill the Wightman axiom relative to the covariance transformation properties of relativistic quantum fields \([16]\).)

Since (heavy) quarks are confined in the hadron “volume”, it seems natural to introduce a Gaussian (peaked at the center of the hadron) as a space averaging function at a given
instant chosen as $t = 0$

$$f(x) = \frac{1}{(2\pi)^{3/4} \sigma_0^{3/2}} \exp \left[ -\frac{x^2}{4\sigma_0^2} \right] \delta(t) \quad (13)$$

where $\sigma_0$ should be interpreted as a measure of the the typical hadronic size, i.e. the hadronic radius $\simeq 1$ fm.

In (residual) momentum space, the Fourier transform reads

$$\tilde{f}(\vec{k}) = \frac{1}{(2\pi)^{3/4} \tilde{\sigma}^{3/2}} \exp \left[ -\frac{\vec{k}^2}{4\tilde{\sigma}^2} \right] \quad (14)$$

where $\sigma_0 = 1/2\tilde{\sigma}$. Therefore, the smeared out annihilation operators are defined as

$$b_r(\tilde{f}) = \int d\mu(k) \ \tilde{f}(\vec{k}) \ \varphi_r(\vec{k}) \ b_r(v, \vec{k}) \quad (15)$$

with $\varphi_r$ standing for the upper components of the spinor. Then it follows that the smeared heavy-quark effective field can be written as

$$h_v(f) = N \sum_{r=-1/2}^{1/2} b_r(\tilde{f}) \quad (16)$$

where $N$ is the appropriate normalization factor in HQET [10].

3 Heavy-Quark States in Hilbert Space

When acting on the vacuum, $\overline{h}_v(f)$ generates a normalizable state as an ensemble of plane-wave states with momenta about $m_Qv$, corresponding to a minimum wave packet with the choice (14). Let us remark that the velocity superselection rule of HQET [5] - which forbids combinations from different sectors of Hilbert space related to different $v$'s - does not actually apply in this case since the residual momenta correspond to quantum fluctuations around $m_Qv$.

In order to study the evolution of the wave packet in space-time according to the Schrödinger picture, let us write at $t = 0$

$$< x|\overline{h}_v(f)|0 > = N \sum_{r=-1/2}^{1/2} \int d\mu(k) \ \tilde{f}(\vec{k}) \ \varphi_r(\vec{k}) \ e^{-i\vec{k}\vec{x}} \quad (17)$$

where the Gaussian function $\tilde{f}$ actually restricts the integral to small values of $\vec{k}$.

The $\tilde{\sigma}$ parameter in Eq. (14) can be interpreted as the root-mean-square residual momentum (for each spatial component) of the heavy quark inside the hadron,

$$< \vec{k}^2 > = 3 \ \tilde{\sigma}^2 \quad (18)$$

which, in turn, can be related to the average non-relativistic kinetic energy $K$

$$K = \frac{< k^2 >}{2m_Q} = \frac{3 \ \tilde{\sigma}^2}{2m_Q} \quad (19)$$
On the other hand, the expectation value of the \( 1/m_Q \) kinetic term of the HQET Lagrangian in (9) is usually written as

\[
K^{HQET} = -\frac{< H(v) | \overline{D}_\perp (D_\perp)^2 v | H(v) >}{2m_Q} = \frac{-\lambda_1}{2m_Q} \tag{20}
\]

with a mass-independent normalization of states. The \( \lambda_1 \) parameter together with \( \Lambda \) and \( \lambda_2 \) (related to the light degrees of freedom and the chromomagnetic interaction respectively) are basic quantities of HQET, playing a fundamental role in many of its phenomenological applications.

Although the HQET kinetic operator is not multiplicatively renormalized, its mixing with the identity operator under ultraviolet renormalization should lead to an additive quadratically divergent contribution to \( \lambda_1 \) which has to be non-perturbatively subtracted \[18\] in the quest for a "physical" quantity \( \lambda_1^{\text{phys}} \) \[3, 19\]. Although the use of a hard ultraviolet cut-off delivers matrix elements from renormalon ambiguities, unfortunately, the arbitrariness in the choice of the subtraction scheme reintroduces the ambiguity. Such an uncertainty, however, is expected to be small if a Lorentz-invariant regularization is employed since then the mixing occurs at next-to-leading order (i.e. two-loop order), as explicitly shown in \[3\]. On the other hand, determinations of \( \lambda_1 \) using dimensional regularization are contaminated by the renormalon ambiguity problem.

Hence, in our characterization of heavy-quark states by wave packets we shall assume

\[< \vec{k}^2 > = -\lambda_1^{\text{phys}}, \]

where current numerical values of the last parameter in literature may be possibly affected by ambiguities of order \( \Lambda_{QCD}^2 \).

As already mentioned, the action when the field operator \( \overline{b}_v(f) \) acts on the vacuum is to create a wave packet of typical width \( \sigma_0 \) in ordinary space at \( t = 0 \). However, as time elapses the wave packet representing the “temporarily free” heavy-quark spreads out, its spatial width becoming increasingly larger according to \[20\]

\[
\sigma(t) = \sigma_0 \left( 1 + \frac{4\bar{\sigma}^2 t^2}{m_Q^2} \right)^{1/2} \tag{21}
\]

whereas the momentum width \( \bar{\sigma} \) remains the same. In order for the wave packet description of localized (within the hadron volume) massive quarks to remain meaningful, one must require that

\[
\bar{\sigma}^2 \ll \frac{m_Q}{2t} \tag{22}
\]

The characteristic time \( t \) of the above expression can be identified with \( \Delta t \) given by Eq. (3) \[3\]. Therefore setting \( t \simeq (3/2) \Lambda_{QCD}^{-1} \) in Eq. (22) we get the condition:

\[
\bar{\sigma}^2 \ll \frac{1}{3} m_Q \Lambda_{QCD} \tag{23}
\]

and with the aid of (18), one can deduce that \[\Box\]

\[
-\lambda_1^{\text{phys}} \ll m_Q \Lambda_{QCD} \tag{24}
\]

\footnote{Tentatively identifying \( t \) with the typical lifetime of B mesons, of the order of the psec \( 10^{-12} \)s makes no sense at all since in order to satisfy the inequality (22) the initial wave packet should be too narrow in momentum space, or equivalently exceedingly large in ordinary space in contradiction with the typical size of a hadron where the heavy quark is to be confined. Therefore it seems natural to consider the hadron lifetime divided in tiny time intervals \( \simeq \Lambda_{QCD}^{-1} \) in the hadron rest frame during which the free wave packet description is meaningful.}

\footnote{Any other physical scale of the order \( \Lambda_{QCD} \) such as \( \Lambda \) could be used instead.}
in spite of ambiguities of order $\Lambda_{QCD}^2$. As a matter of fact, what the above strong inequality actually means is that, for consistency, $-\lambda_1^{phys}$ must be much smaller than the scale of the light degrees of freedom times the heavy quark mass, both sides being ambiguous to a similar extent which should not invalidate it though.

On the other hand, the uncertainty principle also implies a lower bound to the physical value of the heavy quark residual square momentum. Hence we conclude that the following double inequality must be satisfied for a particle-like description of heavy-quarks:

$$\Lambda_{QCD}^2 \leq -\lambda_1^{phys} \ll m_Q \Lambda_{QCD}$$  \hspace{1cm} (25)

Observe that the rhs of expression (25) crudely states that the kinetic energy of the heavy quark must be much smaller than $\Lambda_{QCD}$, i.e. much smaller than the energy exchange with the gluonic cloud. (In the infinite mass limit the massive quark would become a pure static source of colour.) In particular, setting standard numerical values for the $b$-quark mass and $\Lambda_{QCD}$ we determine the following range (in GeV$^2$ units):

$$0.06 \leq -\lambda_1^{phys} \ll 1$$  \hspace{1cm} (26)

Consequently note that the last expression favours relatively small values for the $-\lambda_1^{phys}$ parameter, as for example the result quoted by Neubert in [21] making use of the field-theory analog of the virial theorem for heavy-light hadrons [22] (see also Ref. [23] for phenomenological determinations from $B$ and $D$ meson decays). Nevertheless, since our expression (26) is in reality an order-of-magnitude estimate, somewhat larger values for the average square momentum [2, 24] do not necessarily lead to contradiction with a particle-like behaviour of the heavy quark.

What about charm quarks in hadrons?

Setting $m_c = 1.5$ GeV and $\Lambda_{QCD} = 0.25$ GeV the double inequality (in GeV$^2$ units)

$$0.06 \leq -\lambda_1^{phys} \ll 0.4$$  \hspace{1cm} (27)

now leaves little room for the meaning of a localized $c$-quark in a hadron. Needless to say, if light quarks are considered $(m_q \simeq \Lambda_{QCD})$ it becomes apparent that the inequalities (25) can not be simultaneously satisfied any more and the formalism is not appropriate at all. In the language of ordinary Quantum Mechanics light quarks would display a wave-like character.

\section{4 Conclusions}

In this work we have averaged effective fields $h_v(x)$ defined in HQET for almost on-shell heavy quarks by means of smearing functions (Gaussians spatially peaked at the center of the hadron) yielding operator-valued distributions in Hilbert space. Thus, smoothed field operators acting on the vacuum create normalized wave packets representing heavy quarks as superposition of plane-wave states with residual momenta around $m_Q v$. Their subsequent time evolution in ordinary space leads, however, to increasingly larger spatial spreads, becoming incompatible with a particle-like interpretation for a typical $B$ meson lifetime. In fact, the validity of this description has to be limited to the characteristic
hadronic time scale $\Lambda_{QCD}^{-1}$. Thereby, we conclude that the double inequality (25) heuristically suggests relatively small absolute values of the HQET parameter $\lambda_1^{phys}$ for a $b$-quark in a hadron, either meson or baryon. Such a picture becomes doubtful for $c$-quarks and clearly meaningless for light quarks.

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