SU(3) flavour symmetry breaking and charmed states

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By extending the SU(3) flavour symmetry breaking expansion from up, down and strange sea quark masses to partially quenched valence quark masses we propose a method to determine charmed quark hadron masses including possible QCD isospin breaking effects. Initial results for some open charmed pseudoscalar meson states and singly and doubly charmed baryon states are encouraging and demonstrate the potential of the procedure. Essential for the method is the determination of the scale using singlet quantities, and to this end we also give here a preliminary estimation of the recently introduced Wilson flow scales.

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1. Motivation and Strategy

At present there is considerable interest in open charm masses. While much is known about $C = 1$ charmed meson masses and singly charmed baryon masses, the situation is far less clear for doubly charmed quark baryons (i.e. $C = 2$ or $ccq$ with $q = u, d, s$). Here many masses are unknown, but as stable states in QCD they must exist (presently only some candidate states are seen by the SELEX Collaboration [1], but not by BaBar [2], or BELLE [3]). Also there are possible relations to tetraquark states. For example in the $n_c \to \infty$, $m_c \to \infty$ limit we expect the relation

$$M(\Xi_{cc}^+) \approx M_{\Lambda_b^+} + M_{D_s^+} - (M_{D_s^+} + M_{D^0})/4.$$  

The charm quark, $c$, is considerably heavier than the up $u$, down $d$ and strange $s$ quarks, which has hampered its direct simulation using lattice QCD. However as the available lattice spacings become finer, this is becoming less of an obstacle. The sea quarks in present day numerical simulation.

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2. Method

We presently only consider hadrons which lie on the outer ring of their associated multiplet and not the central hadrons. (So we need not consider either any mixing or numerical evaluation of quark–line disconnected correlation functions.) The $SU(3)$ flavour symmetry breaking expansion for the pseudoscalar mesons with valence quarks $a$ and $b$ up to cubic or NNLO terms in the quarks masses is given by

$$M^2(a\bar{b}) = M^2_{0\pi} + \alpha(\delta\mu_a + \delta\mu_b)$$

$$+ \beta_1(\delta\mu_a^2 + \delta\mu_b^2) + \beta_2(\delta\mu_a - \delta\mu_b)^2$$

$$+ \gamma_0\delta m_d\delta m_s + \gamma_1(\delta\mu_a + \delta\mu_b)(\delta m_a^2 + \delta m_b^2 + \delta m_s^2)$$

$$+ \gamma_2(\delta\mu_a + \delta\mu_b)^3 + \gamma_3(\delta\mu_a + \delta\mu_b)(\delta\mu_a - \delta\mu_b)^2,$$  \hspace{1cm} (2.1)
with $\delta \mu_q = \mu_q - m_q, q = a, b, \ldots \in \{u, d, s, c\}$ being valence quarks of arbitrary mass, $\mu_q$ and $\delta m_q = m_q - m, q \in \{u, d, s\}$ being sea quarks. (These have the automatic constraint $\delta m_u + \delta m_d + \delta m_s = 0$.) Note that we have some mixed sea/valence mass terms. The unitary limit occurs when $\delta \mu_q \rightarrow \delta m_q$. The expansion coefficients are $M^2_{\alpha\beta}(\bar{m}), \alpha(\bar{m}), \ldots$ so if $\bar{m}$ is held constant then we have constrained fits to the numerical data. In particular a $n_f = 2 + 1$ flavour simulation, when $\delta m_u = \delta m_d \equiv \delta m_1$ is enough to determine the expansion coefficients. We now use the PQ (and unitary) data to determine the expansion coefficients (i.e. $\alpha$, $\beta$, $\gamma$). This in turn leads to a determination of the ‘physical’ quark masses $\delta m^*_{u}, \delta m^*_{d}, \delta m^*_{s}$ and $\delta \mu^*_c$ by fitting to e.g. $M^\text{exp}_{\pi^+}^{\text{uu}}, M^\text{exp}_{K^+}^{\text{us}}$ and $M^\text{exp}_{\eta^c}^{\text{cc}}$. We can now describe pseudoscalar open charm states with the same wavefunction (and hence expansion) $\mathcal{M} = \bar{q}p c \ (q = u, d, s)$ i.e. $D^0(c\bar{u})$, $D^+(c\bar{d})$ and $D^+_s(c\bar{s})$. Using $\delta m^*_{u}, \delta m^*_{d}, \delta m^*_{s}$ and $\delta \mu^*_c$ gives estimates of their physical masses.

Similarly for the baryon octet the same procedure can be applied. We have the $SU(3)$ flavour symmetry breaking expansion

$$M^2(aab) = M^2_{\text{ON}} + A_1(2\delta \mu_a + \delta \mu_b) + A_2(\delta \mu_b - \delta \mu_a) + B_0 \frac{1}{6} (\delta m^2_u + \delta m^2_d + \delta m^2_s) + B_1(2\delta \mu^2_u + \delta \mu^2_b) + B_2(\delta \mu^2_b - \delta \mu^2_d) + B_3(\delta \mu_b - \delta \mu_a)^2 + C_0 \delta m_u \delta m_d \delta m_s + [C_1(2\delta \mu_u + \delta \mu_b) + C_2(\delta \mu_b - \delta \mu_a)](\delta m^2_u + \delta m^2_d + \delta m^2_s) + C_3(\delta \mu_u + \delta \mu_b)^3 + C_4(\delta \mu_u + \delta \mu_b)^2(\delta \mu_u - \delta \mu_b) + C_5(\delta \mu_u + \delta \mu_b)(\delta \mu_u - \delta \mu_b)^2 + C_6(\delta \mu_a - \delta \mu_b)^3, \quad (2.2)$$

(so for example $M_\rho = M(uud)$). Again we use PQ (and unitary) data to first determine the expansion coefficients (i.e. the $A$, $B$, $C$). We can then describe charm states with the same nucleon like wavefunction and hence same expansion. For example for single open charm ($C = 1$) states we have $\Sigma^+_c(uuc), \Xi^0_c(dcc), \Omega^0_c(ssc)$ which all have the wavefunction $\mathcal{B} = \epsilon(q^T C\gamma_5 c)q \ (q = u, d, s)$ while for double open charm ($C = 2$) states we have $\Xi^{++}_c(ccu), \Xi^{++}_c(ccd), \Omega^{++}_c(ccs)$ which all have the wavefunction $\mathcal{B} = \epsilon(q^T C\gamma_5 q)c \ (q = u, d, s)$. In both cases using $\delta m^*_{u}, \delta m^*_{d}, \delta m^*_{s}$ and $\delta \mu^*_c$ gives estimates of these physical masses.

### 3. Lattice details

We use a tree level Symanzik gluon action and an $O(a)$ improved clover fermion action, including mild stout smearing [8]. Thus the quark mass is given by

$$m_q = \frac{1}{2} \left( \frac{1}{\kappa_q} - \frac{1}{\kappa_0} \right), \quad (3.1)$$

where $\kappa_q$ is the hopping parameter, $\kappa_0$ is the hopping parameter along the symmetric line with $\kappa_0c$ being its chiral limit. We shall consider two lattice spacings ($\beta = 5.50$ on $32^3 \times 64$ lattices and preliminary results for $\beta = 5.80$ on $48^3 \times 96$ lattices).

We shall now briefly mention here our progress in defining and determining the scale using singlet quantities, collectively denoted here by $X_0$. There are many possibilities such as pure gluon quantities like the $r_0$ Sommer scale: $X_0 = 1/r_0$, or the $\sqrt{t_0}$ [4], $\sqrt{t_0}$ [4] scales based on the Wilson gauge action flow: $X_0 = 1/\sqrt{t_0}, X_{w_0} = 1/w_0$, or quantities constructed using fermions. One simple possibility in this case is to take the ‘centre of mass’ of the hadron octet. In all these cases it can
easily be shown that linear terms in $\delta m_q$ are absent, [3]. We then have from eqs. (2.1), (2.2) in the unitary limit (with $\overline{\delta m^2} = (\delta m_0^2 + \delta m_s^2 + \delta m_2^2)/3$)

$$
X_0^2 = \frac{1}{6} \left( M_{K^0}^2 + M_{K^\pm}^2 + M_{\pi^0}^2 + M_{\pi^\pm}^2 + M_{K^-}^2 + M_{K^0}^2 \right) = M_{0\pi}^2 + \left( \frac{1}{2} \beta_0 + 2 \beta_1 + 3 \beta_2 \right) \overline{\delta m^2} + \ldots , \quad (3.2)
$$

and

$$
X_N^2 = \frac{1}{6} \left( M_N^2 + M_{\pi^0}^2 + M_{\pi^\pm}^2 + M_{\pi^0}^2 + M_{K^-}^2 + M_{K^0}^2 \right) = M_{0N}^2 + \frac{1}{2} (B_0 + B_1 + B_3) \overline{\delta m^2} + \ldots . \quad (3.3)
$$

In the left panel of Fig. 1 we plot various singlet quantities $(a_S X_S)^2$ for $S = t_0$, $w_0$, $\pi$, $\rho$, $N$. It is apparent that the constancy $(a_S X_S)^2$ holds over the complete range from the symmetric point down to the physical point. Using this enables us to use $X_S^{\exp}$ to determine the lattice spacing by

$$
da_S^2 = \frac{(a_S X_S)^2}{X_S^{\exp 2}} . \quad (3.4)
$$

We shall define our lattice spacing here using $S = N$: $a_N$. Of course depending on how well we have chosen our initial $\kappa_0$ point, using a different singlet quantity (i.e. $S \neq N$) will give a slightly different lattice spacing. More ambitiously we can vary $\kappa_0$ to try to find a point where we have a common scale. We have initiated a programme to investigate this. In the right panel of Fig. 1 we plot $a_S^2$ again for $S = \pi$, $N$, $\rho$, $w_0$, $t_0$ against various $\kappa_0$. The crossing of the $a_S^2$ for $S = \pi$, $N$ and $\rho$ give an estimation of the common scale as $\alpha \approx 0.074 (2)$ fm. We can now adjust $X_{0\pi}^{\exp}$, $X_{0\rho}^{\exp}$ to also cross at this point to find a preliminary estimate for these ‘intermediate scales’ of $\sqrt{X}^{\exp} \approx 0.153 (7)$ fm, $w_0 \approx 0.179 (6)$ fm.

Practically it is numerically advantageous to form dimensionless ratios (within a multiplet): $\tilde{M}^2 \equiv M^2 / X_S^2$ and re-write eqs. (2.1), (2.2) in terms of $\tilde{\alpha} \equiv \alpha / M_{0\pi}$, $\ldots$ and $\tilde{A_i} \equiv A_i / M_{0N}$, $\ldots$ in the expansions. About $\sim O(80)$ PQ and unitary masses are used to determine these expansion coefficients and hence the ‘physical’ quark masses $\delta m_u$, $\delta m_d$, $\delta m_s$ and $\delta \mu$ as described in section 2.

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4. Results and Conclusions

We now discuss our results. In Fig. 2 left panel, we show the diagonal pseudoscalar mesons 

\[ M_2^2(a^a) = M_2^2(a^a')/X^2 \] versus \( \delta \mu_a \) for \( \beta = 5.50 \), together with the fit from eq. (2.1). The vertical dashed line represents the symmetric point, while the horizontal dashed line is the physical value of \( \tilde{M}_{\eta_c}^2 \). Right panel: ‘fan’ plot for the baryon octet, from the symmetric point \( \delta m_l = 0 \) to the physical point \( \delta m_l = (\delta m_u + \delta m_d)/2 \) (vertical dashed line and stars) for \( \beta = 5.50 \). The filled triangles are from \( 32^3 \times 64 \) sized lattices, while the open triangles are from \( 24^3 \times 48 \) sized lattices (not used in the fits). The fits are again given from eq. (2.1).

In Fig. 3 we show \( D_0(c^d) \), \( D_+^+(c^d) \) and \( D_+^+(c^s) \) against our \( a_N^2 \) lattice spacings (left panel) and their mass differences (right panel). These mass differences in particular are sensitive to unknown QED effects (the present computation is for pure QCD only). As we currently have only two lattice spacings (and are also increasing their statistics) the results are to be regarded as preliminary and we do not presently attempt a continuum extrapolation. However there do not seem to be strong scaling violations present.

In Figs. 4 and 5 we show the \( C = 1 \) and \( C = 2 \) charmed baryons (from the spin 1/2 20-plet). Again while we do not see significant lattice effects in either case, we gain an impression that those present are a little larger for the doubly charmed baryons than for the singly charmed mesons.

In conclusion we note that we have developed a method to determine some open charm states using a precise SU(3) flavour symmetry breaking expansion enabling \( u, d, s \) quarks to approach the physical point while the \( c \) quark is treated as PQ. The expansion appears to be highly convergent. The method can be extended to other states. In a \( 2 + 1 \) world there is no \( \Sigma^0 - \Lambda^0 \) mixing, but the determined coefficients can be used to compute \( \Sigma^0(uds) - \Lambda^0(uds) \) mixing, [12]. Therefore
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\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{Figure3}
\caption{Left panel: $D^0(c\bar{u})$, $D^+(c\bar{d})$ and $D^+_s(c\bar{s})$. Right panel: $D^+(c\bar{d}) - D^0(c\bar{u})$, $D^+_s(c\bar{s}) - D^0(c\bar{u})$ and $D^+_s(c\bar{s}) - D^+(c\bar{d})$ mass splittings. (All values in MeV.) The experimental values are given as red stars. To guide the eye, we extend these values as horizontal dashed lines.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{Figure4}
\caption{Left panel: $\Sigma^{++}_c(uuc)$, $\Sigma^0_c(ddc)$, $\Omega^0_c(ssc)$, $\Omega^{+}_c(ssc) - \Sigma^0_c(ddc)$, $\Omega^0_c(ssc) - \Sigma^0_c(uuc)$, $\Sigma^0_c(uuc) - \Sigma^0_c(uuc)$, $\Sigma^0_c(ddc) - \Sigma^0_c(uuc)$, $\Omega^0_c(ssc) - \Sigma^0_c(ssc)$. Right panel: $\Sigma^+_c(uuc)$, $\Sigma^0_c(ddc)$, $\Omega^0_c(ssc)$, $\Omega^0_c(ssc) - \Sigma^0_c(ddc)$, $\Omega^0_c(ssc) - \Sigma^0_c(uuc)$, $\Sigma^0_c(uuc) - \Sigma^0_c(uuc)$, $\Sigma^0_c(ddc) - \Sigma^0_c(uuc)$, $\Omega^0_c(ssc) - \Sigma^0_c(ssc)$. Computing e.g. $\Sigma^+_c - \Lambda^+_c$, $\Xi^0_c - \Xi^0_c$ mixing is possible. Furthermore the method can be applied to the baryon decuplet and QED effects can be introduced, [13].}
\end{figure}

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Figure 5: Left panel: $\Xi_{cc}^+(ccd) - \Xi_{cc}^+(ccu)$. The SELEX Collaboration result, [1], for $\Xi_{cc}^+$ is shown as a red star. Right panel: $\Omega_{cc}^+(ccs) - \Xi_{cc}^+(ccu)$, $\Omega_{cc}^+(ccs) - \Xi_{cc}^+(ccd)$ mass splittings.

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