

## $T_{cc}$ pole trajectory

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**Protick Mohanta,<sup>a,\*</sup> Srijit Paul<sup>b,c</sup> and Subhasish Basak<sup>d,e</sup>**

<sup>a</sup>*The Institute of Mathematical Sciences, Chennai 600113, India*

<sup>b</sup>*Maryland Center for Theoretical Physics, University of Maryland, College Park, MD 20742 USA*

<sup>c</sup>*Department of Physics, University of Cyprus, Aglantzia 2109, Nicosia, Cyprus*

<sup>d</sup>*School of Physical Sciences, National Institute of Science Education and Research, An OCC of Homi Bhabha National Institute, Jatni 752050, India*

<sup>e</sup>*Homi Bhabha National Institute, Training School Complex, Anushakti Nagar, Mumbai 400094, India*

*E-mail:* [protickm@imsc.res.in](mailto:protickm@imsc.res.in), [spaul137@umd.edu](mailto:spaul137@umd.edu), [sbasak@niser.ac.in](mailto:sbasak@niser.ac.in)

We investigate the spectrum of doubly charmed tetraquark  $T_{cc}$  with quantum number  $I(J^P) = 0(1^+)$  using MILC's  $N_f = 2 + 1 + 1$  HISQ gauge ensembles at two lattice spacings. We have included diquark-antidiquark operator together with molecular and scattering operators in our analysis and varied both the heavy and light quark masses. We employ the anisotropic Clover action for heavy quarks, and  $O(a)$ -improved Wilson–Clover action for the light (up/down) quarks. In order to handle the non-analyticity near the Left Hand Cut we use modified Lüscher's method when we are close to it.

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\*Speaker

## 1. Introduction

The discovery of doubly charmed tetraquark  $T_{cc}$  [1] has led to extensive investigations, both lattice and non-lattice, of doubly charm ( $T_{cc}$ ), doubly bottom ( $T_{bb}$ ) and other possible heavy tetraquark candidates like  $T_{bc}$ ,  $bb\bar{u}\bar{s}$  etc. Though lattice QCD mostly indicates the existence of deeply bound doubly bottom tetraquark  $T_{bb}(bb\bar{u}\bar{d})$  but the theoretical prediction is less certain about the existence of doubly charm tetraquark  $T_{cc}(cc\bar{u}\bar{d})$  [2]. The pole of  $T_{cc}$  lies a mere 0.36 MeV below the  $D^0 D^{*+}$  threshold, whereas the binding energy of  $T_{bb}$  is  $\sim O(100)$  MeV [3–7]. As a first step to address this, we have varied the heavy quark mass and taken three additional mass points between charm and bottom mass. The existence of Left Hand Cut (LHC) complicates the situation as one approaches the physical pion mass limit [8]. Therefore while varying the light quark mass, we used Lüscher method when we are away from LHC and modified Lüscher method [9] when closer to it. In this proceeding we report our preliminary result and the status of our ongoing  $T_{cc}$  pole analysis.

## 2. Operator Basis

Many early studies such as [10] have not included diquark-antidiquark operator in their analysis. Cheung *et al.* [11] found diquark-antidiquark operator not to have significant effects on finite volume spectra and Cheng *et al.* [12] showed such operator results in an unstable  $T_{cc}$ . Most of the past lattice  $T_{cc}$  investigations thus did not consider diquark-antidiquark operator in their simulations. However, studies based on heavy quark symmetries [13, 14] showed the usefulness of diquark operators in doubly heavy tetraquark states. Another striking point is the mass differences in  $J^P = (1/2)^+$  singly heavy baryons  $(\Lambda_b, \Sigma_b)$ ,  $(\Lambda_c, \Sigma_c) \sim 191 - 167$  MeV. The operators

$$\left(O_k^{hl_1l_2}\right)_\alpha = \epsilon_{abc} [l_1^{aT} C \gamma_k l_2^b] Q_\alpha^c \quad \text{and} \quad \left(O_5^{hl_1l_2}\right)_\alpha = \epsilon_{abc} [l_1^{aT} C \gamma_5 l_2^b] Q_\alpha^c \quad (1)$$

gives a mass difference of  $\Sigma_b - \Lambda_b \sim 30$  MeV at a pion mass of 490 MeV. When we reduce the pion mass, the above two operators in eqn. (1) produce significantly enhanced mass splitting [15]. And since the diquark-antidiquark operator is related to the  $\Lambda_Q$  operator by heavy quark-diquark symmetry, we expect it to have important contribution to  $T_{cc}$ . We, therefore, use the following operators in our analysis

$$\begin{aligned} \mathcal{D}(x) &= [c(x)^{aT} C \gamma_k c(x)^b] [\bar{u}(x)^a C \gamma_5 \bar{d}(x)^{bT}] \\ \mathcal{M}_1(x) &= [\bar{d}(x)^a \gamma_k c(x)^a] [\bar{u}(x)^b \gamma_5 c(x)^b] \\ \mathcal{M}_2(x) &= \epsilon_{kij} [\bar{d}(x)^a \gamma_i c(x)^a] [\bar{u}(x)^b \gamma_j c(x)^b] \\ \mathcal{S}(t; \mathbf{p}_1, \mathbf{p}_2) &= \sum_{\mathbf{x}} [\bar{d}(x)^a \gamma_k c(x)^a] e^{i\mathbf{p}_1 \cdot \mathbf{x}} \times \sum_{\mathbf{y}} [\bar{u}(y)^b \gamma_5 c(y)^b] e^{i\mathbf{p}_2 \cdot \mathbf{y}} \end{aligned} \quad (2)$$

We perform our analysis in the centre-of-mass frame, hence the  $D^*$  and  $D$  mesons in the scattering operator  $\mathcal{S}$  are given back-to-back momentum  $\mathbf{p}_1 + \mathbf{p}_2 = 0$ . We have generated data up to  $\mathbf{p}^2 = 1$  and performed GEVP analysis of the  $5 \times 5$  correlator matrix. Here we present three diagonal elements

of GEVP matrix,  $C_{\mathcal{D}\mathcal{D}}(t)$ ,  $C_{\mathcal{M}_i\mathcal{M}_i}(t)$  and  $C_{\mathcal{S}\mathcal{S}}(t)$ ,

$$\begin{aligned}
 C_{\mathcal{D}\mathcal{D}}(t) &= \sum_{\vec{x}} \langle \mathcal{D}(x) \mathcal{D}(0)^\dagger \rangle \\
 &= \sum_{\vec{x}} \text{Tr} \left[ \left\{ \mathcal{G}_c(t, \vec{x}; 0) \right\}^{adT} \left( \gamma_k \gamma_4 \gamma_2 \mathcal{G}_c(t, \vec{x}; 0) \gamma_4 \gamma_2 \gamma_k \right)^{bc} \right] \\
 &\quad \times \text{Tr} \left[ \left\{ \gamma_4 \gamma_2 \mathcal{G}_u(t, \vec{x}; 0)^\dagger \gamma_4 \gamma_2 \right\}^{da} \left( \gamma_5 \mathcal{G}_d(t, \vec{x}; 0)^\dagger \gamma_5 \right)^{cbT} \right] \\
 &\quad - \sum_{\vec{x}} \text{Tr} \left[ \left\{ \mathcal{G}_c(t, \vec{x}; 0) \gamma_4 \gamma_2 \gamma_k \right\}^{ac} \left( \gamma_k \gamma_4 \gamma_2 \mathcal{G}_c(t, \vec{x}; 0) \right)^{bdT} \right] \\
 &\quad \times \text{Tr} \left[ \left\{ \gamma_4 \gamma_2 \mathcal{G}_u(t, \vec{x}; 0)^\dagger \gamma_4 \gamma_2 \right\}^{da} \left( \gamma_5 \mathcal{G}_d(t, \vec{x}; 0)^\dagger \gamma_5 \right)^{cbT} \right] \quad (3)
 \end{aligned}$$

$$\begin{aligned}
 C_{\mathcal{M}_1\mathcal{M}_1}(t) &= \sum_{\vec{x}} \langle \mathcal{M}_1(x) \mathcal{M}_1(0)^\dagger \rangle \\
 &= \sum_{\vec{x}} \text{Tr} \left[ \left\{ \gamma_5 \mathcal{G}_d(t, \vec{x}; 0)^\dagger \gamma_5 \right\} \left( \gamma_k \mathcal{G}_c(t, \vec{x}; 0) \gamma_k \right) \text{Tr} \left[ \left\{ \gamma_5 \mathcal{G}_u(t, \vec{x}; 0)^\dagger \gamma_5 \right\} \left( \gamma_5 \mathcal{G}_c(t, \vec{x}; 0) \gamma_5 \right) \right] \right] \\
 &\quad - \sum_{\vec{x}} \text{Tr} \left[ \left\{ \gamma_5 \mathcal{G}_d(t, \vec{x}; 0)^\dagger \gamma_5 \right\} \left( \gamma_k \mathcal{G}_c(t, \vec{x}; 0) \gamma_5 \right) \left\{ \gamma_5 \mathcal{G}_u(t, \vec{x}; 0)^\dagger \gamma_5 \right\} \left( \gamma_5 \mathcal{G}_c(t, \vec{x}; 0) \gamma_k \right) \right] \quad (4)
 \end{aligned}$$

$$\begin{aligned}
 C_{\mathcal{S}\mathcal{S}}(t; \mathbf{p}_1, \mathbf{p}_2, \mathbf{p}_4) &= \langle \mathcal{S}(t; \mathbf{p}_1, \mathbf{p}_2) \mathcal{S}(0; \mathbf{p}_3, \mathbf{p}_4)^\dagger \rangle \\
 &= \sum_{\mathbf{x}, \mathbf{y}, \mathbf{z}} e^{i(\mathbf{p}_1 \cdot \mathbf{x} + \mathbf{p}_2 \cdot \mathbf{y} - \mathbf{p}_4 \cdot \mathbf{z})} \text{Tr} \left[ \left\{ \gamma_5 \mathcal{G}_d(t, \mathbf{x}; 0)^\dagger \gamma_5 \right\} \left( \gamma_k \mathcal{G}_c(t, \mathbf{x}; 0) \gamma_k \right) \right] \\
 &\quad \times \text{Tr} \left[ \mathcal{G}_c(t, \mathbf{y}; 0, \mathbf{z}) \left\{ \gamma_5 \mathcal{G}_u(0, \mathbf{z}; t, \mathbf{y}) \gamma_5 \right\} \right] \\
 &\quad - \sum_{\mathbf{x}, \mathbf{y}, \mathbf{z}} e^{i(\mathbf{p}_1 \cdot \mathbf{x} + \mathbf{p}_2 \cdot \mathbf{y} - \mathbf{p}_4 \cdot \mathbf{z})} \text{Tr} \left[ \left\{ \gamma_k \gamma_5 \mathcal{G}_d(t, \mathbf{x}; 0)^\dagger \gamma_5 \gamma_k \right\} \left( \mathcal{G}_c(t, \mathbf{x}; 0, \mathbf{z}) \right. \right. \\
 &\quad \left. \left. \gamma_5 \mathcal{G}_u(0, \mathbf{z}; t, \mathbf{y}) \gamma_5 \mathcal{G}_c(t, \mathbf{y}; 0) \right) \right] \quad (5)
 \end{aligned}$$

Except for the correlator  $C_{\mathcal{S}\mathcal{S}}(t; \mathbf{p}_1, \mathbf{p}_2, \mathbf{p}_4) = \langle \mathcal{S}(t; \mathbf{p}_1, \mathbf{p}_2) \mathcal{S}(0; \mathbf{p}_3, \mathbf{p}_4)^\dagger \rangle$ , all other require computation of only point-all propagators. For the  $C_{\mathcal{S}\mathcal{S}}$  we made use of the one end trick as suggested in [16]. In this case, we make use of translational invariance to remove the phase related to  $\mathbf{p}_3$ . We implemented complex  $\mathbb{Z}(2) \times \mathbb{Z}(2)$  random numbers at time slice  $t = 0$  for inversion of the Dirac operators corresponding to charm and up/down quarks. For instance, the first term of the eqn. (5) can be written as

$$\begin{aligned}
 &\frac{1}{N} \sum_n \sum_{\mathbf{x}} e^{i(\mathbf{p}_1 \cdot \mathbf{x})} \left\{ \gamma_k \gamma_5 \mathcal{G}_d(t, \mathbf{x}; 0)^\dagger \gamma_5 \gamma_k \right\}_{s_1 s_2}^{c_1 c_2} \left( \mathcal{G}_c(t, \mathbf{x}; 0) \right)_{s_2 s_1}^{c_2 c_1} \\
 &\quad \times \sum_{\mathbf{y}} e^{i(\mathbf{p}_2 \cdot \mathbf{y})} \left( \phi_c^n(\mathbf{y}, t) \right)_{s_3}^{c_3} \left( \phi_u^n(\mathbf{y}, t)^\dagger \right)_{s_3}^{c_3}, \quad (6)
 \end{aligned}$$

where,

$$\begin{aligned}
 \left( D_c(r, x) \right)_{s_1 s_2}^{c_1 c_2} \left( \phi_c^n(x) \right)_{s_2}^{c_2} &= \delta_{r_0, 0} \left( \Xi(\mathbf{r})[n] \right)_{s_1}^{c_1} \\
 \left( D_u(r, x) \right)_{s_1 s_2}^{c_1 c_2} \left( \phi_u^n(x) \right)_{s_2}^{c_2} &= \delta_{r_0, 0} \left( \Xi(\mathbf{r})[n] \right)_{s_1}^{c_1} e^{i(\mathbf{p}_4 \cdot \mathbf{r})} \quad (7)
 \end{aligned}$$

### 3. Quark action and tuning

#### 3.1 Clover action for light up/down quark

To simulate light up/down quarks, we employ the  $O(a)$  improved Wilson–Clover fermion action, which includes a clover leaf term that systematically removes leading discretization error at  $O(a)$ . The corresponding fermion action is expressed as

$$\begin{aligned} S_{\text{clover}} = & \sum_n \bar{\psi}(n)\psi(n) - \kappa_s \left[ \sum_{n,\mu} \bar{\psi}(n)(1 - \gamma_\mu)U_\mu(n)\psi(n + \hat{\mu}) \right. \\ & \left. + \sum_{n,\mu} \bar{\psi}(n)(1 + \gamma_\mu)U_\mu^\dagger(n - \hat{\mu})\psi(n - \hat{\mu}) \right] - \kappa_s c_{\text{SW}} \sum_{n,\mu < \nu} \bar{\psi}(n)\sigma_{\mu\nu}F_{\mu\nu}(n)\psi(n) \end{aligned} \quad (8)$$

where  $\kappa_s = \frac{1}{2(m+4)}$  and  $c_{\text{SW}}$  is the clover coefficient. We vary the pion mass from  $m_{\eta_s} = 688.5$  MeV [19] down to 400 MeV, where  $m_{\eta_s}$  denotes the mass of the fictitious  $\eta_s$  meson [18]. The gauge links are HYP smeared and  $c_{\text{SW}}$  is obtained from [22]

$$c_{\text{SW}} = \frac{1}{u_0^3}, \quad (9)$$

where  $u_0$  is the tadpole improvement factor defined by the fourth root of the average plaquette. We tabulate the pion masses used in our analysis in Table 2.

#### 3.2 Anisotropic clover action for charm quark

In this work, we use the anisotropic clover-improved Wilson action *i.e.* Relativistic Heavy Quark (RHQ) action [20, 21] for the charm quark. The RHQ action modifies the standard Wilson–Clover action by introducing an anisotropy factor  $\nu$ . This enhances temporal resolution which is crucial for heavy quark systems on lattices with shorter time direction sizes. The RHQ action  $S_{\text{RHQ}}$  is given by

$$\begin{aligned} \frac{S_{\text{RHQ}}}{m_0 + 1 + 3\nu} = & \sum_n \bar{\psi}(n)\psi(n) - \kappa \left[ \sum_n \bar{\psi}(n)(1 - \gamma_0)U_0(n)\psi(n + \hat{0}) \right. \\ & + \sum_n \bar{\psi}(n)(1 + \gamma_0)U_0^\dagger(n - \hat{0})\psi(n - \hat{0}) + \nu \sum_{n,i} \bar{\psi}(n)(1 - \gamma_i)U_i(n)\psi(n + \hat{i}) \\ & \left. + \nu \sum_{n,i} \bar{\psi}(n)(1 + \gamma_i)U_i^\dagger(n - \hat{i})\psi(n - \hat{i}) \right] - \kappa c_P \sum_{n,\mu < \nu} \bar{\psi}(n)\sigma_{\mu\nu}F_{\mu\nu}(n)\psi(n), \end{aligned} \quad (10)$$

where  $\kappa = 1/[2(m_0 + 1 + 3\nu)]$ ,  $c_P$  is the clover coefficient,  $\sigma_{\mu\nu} = \frac{i}{2}[\gamma_\mu, \gamma_\nu]$  is the antisymmetric tensor and  $F_{\mu\nu}(n)$  represents the discretized gluon field strength. First we tuned the strange propagator using Wilson–Clover action (8) and subsequently tuned the RHQ action parameters  $\{m_0, c_P, \nu\}$ . The  $m_0$  is tuned by matching spin-average  $aM_{D_s}$  and  $aM_{D_s^*}$  obtained from lattice with the corresponding experimental value  $\overline{M}_{D_s}$

$$\overline{M}_{D_s} = \frac{1}{4}M_{D_s} + \frac{3}{4}M_{D_s^*} \approx 2.076 \text{ GeV} \quad (11)$$

The clover coefficient  $c_P$  is adjusted so that the hyperfine splitting matches its PDG value,

$$\Delta M_{D_s} = M_{D_s^*} - M_{D_s} \approx 143.8 \text{ MeV}. \quad (12)$$

The anisotropy coefficient  $\nu$  is tuned to satisfy relativistic dispersion relation

$$E^2(\vec{p}) = M^2 + \vec{p}^2, \quad (13)$$

ensuring the speed of light is normalized on the lattice to  $c^2 = 1$ . We used  $|\vec{p}| = \frac{2\pi}{L}$  in tuning the RHQ action parameters.

#### 4. Simulation Details

The simulations have been performed using two ensembles of MILC  $N_f = 2 + 1 + 1$  HISQ lattices [24], the details of which are given in Table 1. The parameter tuning and production run on  $24^3 \times 64$  lattices are ongoing, here we present a status update of the project. All measurements are carried out using point source.

$\beta = 10/g^2$	$m_l$	$m_s$	$m_c$	$L^3 \times T$	$a$ (fm)	$N_{cfg}$
5.80	0.013	0.065	0.838	$16^3 \times 48$	0.15	700
6.00	0.0102	0.0509	0.635	$24^3 \times 64$	0.12	700

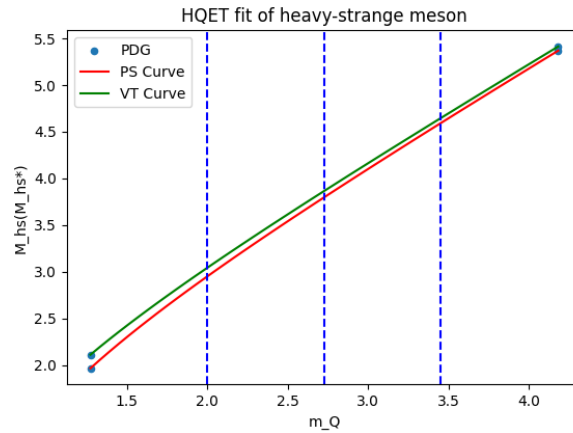
**Table 1:** MILC configurations used in this work and  $N_{cfg}$  is the number of configurations used.

##### 4.1 Heavy-light meson mass

With the varying of heavy quark mass, we need both the heavy-strange pseudoscalar and vector meson masses for the tuning of RHQ action parameters. HQET provides the relevant expression for the masses of heavy-light mesons [23]

$$M_{hl} = m_h + \bar{\Lambda} - \frac{\lambda_1}{2m_h} - \frac{3\lambda_2}{2m_h} \quad \text{and} \quad M_{hl}^* = m_h + \bar{\Lambda} - \frac{\lambda_1}{2m_h} + \frac{\lambda_2}{2m_h} \quad (14)$$

In figure 1 we plot our heavy-strange pseudoscalar  $M_{hl}$  and vector meson  $M_{hl}^*$  mass points, starting from  $D_s, D_s^*$  (the two leftmost points) all the way to  $B_s, B_s^*$  (the two rightmost points).



**Figure 1:** Heavy-light meson points

### 4.2 Light quark tuning

For the strange quark, the  $\kappa_s$  has been tuned separately on each ensemble to arrive at ChiPT  $\bar{s}\gamma_5 s$  state  $\eta_s = 688.5$  MeV mass as discussed before. The rest of the  $\kappa$  values are chosen to get the pion masses approximately equidistant between successive points. To avoid uncontrolled fluctuations, we limited ourselves to  $m_\pi \geq 400$  MeV. We wish to point out that without the HYP-smearing the GEVP matrix loses definite positivity at around time slice  $t = 4$ .

$16^3 \times 48$	$\kappa$	0.12566	0.1260	0.1263	0.1266	0.1269	0.1272	0.1275
	$m_\pi$	688.8(3)	644.7(3)	603.3(3)	559.3(3)	511.7(3)	459.1(3)	400.4(5)
$24^3 \times 64$	$\kappa$	0.1256	0.125845	0.126062	0.12628	0.126498	0.126718	0.126937
	$m_\pi$	688.7(2)	644.7(2)	603.2(2)	559.2(2)	511.7(2)	459.2(2)	400.7(2)

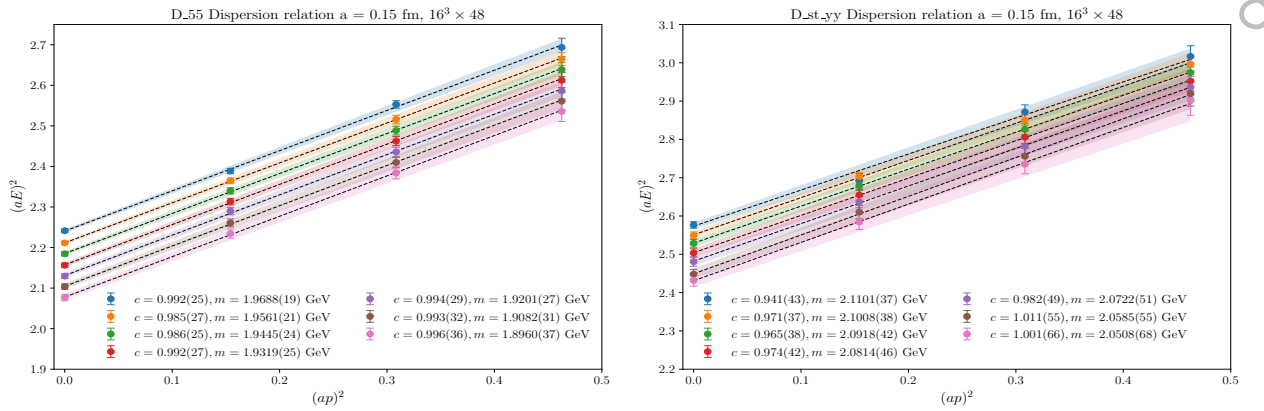
**Table 2:** The light  $\kappa$  used in our analysis. The  $m_\pi$  are in MeV.

### 4.3 RHQ parameters tuning

The step to tune charm quark involves determining  $m_0$  along with the other two RHQ action parameters  $\{c_P, \nu\}$  appearing in (10). As discussed in subsection 3.2, tuning involve matching the spin-averaged mass, the hyperfine splitting and the velocity of light  $c^2 = 1$  for the  $m_Q$  points obtained from eqn. (14) and the plot Fig. 1. The tuned RHQ parameters for  $16 \times 48$  are given in Table 3. In figure 2 we present the dispersion relation of  $D$  and  $D^*$  for  $\kappa = 0.12566$ .

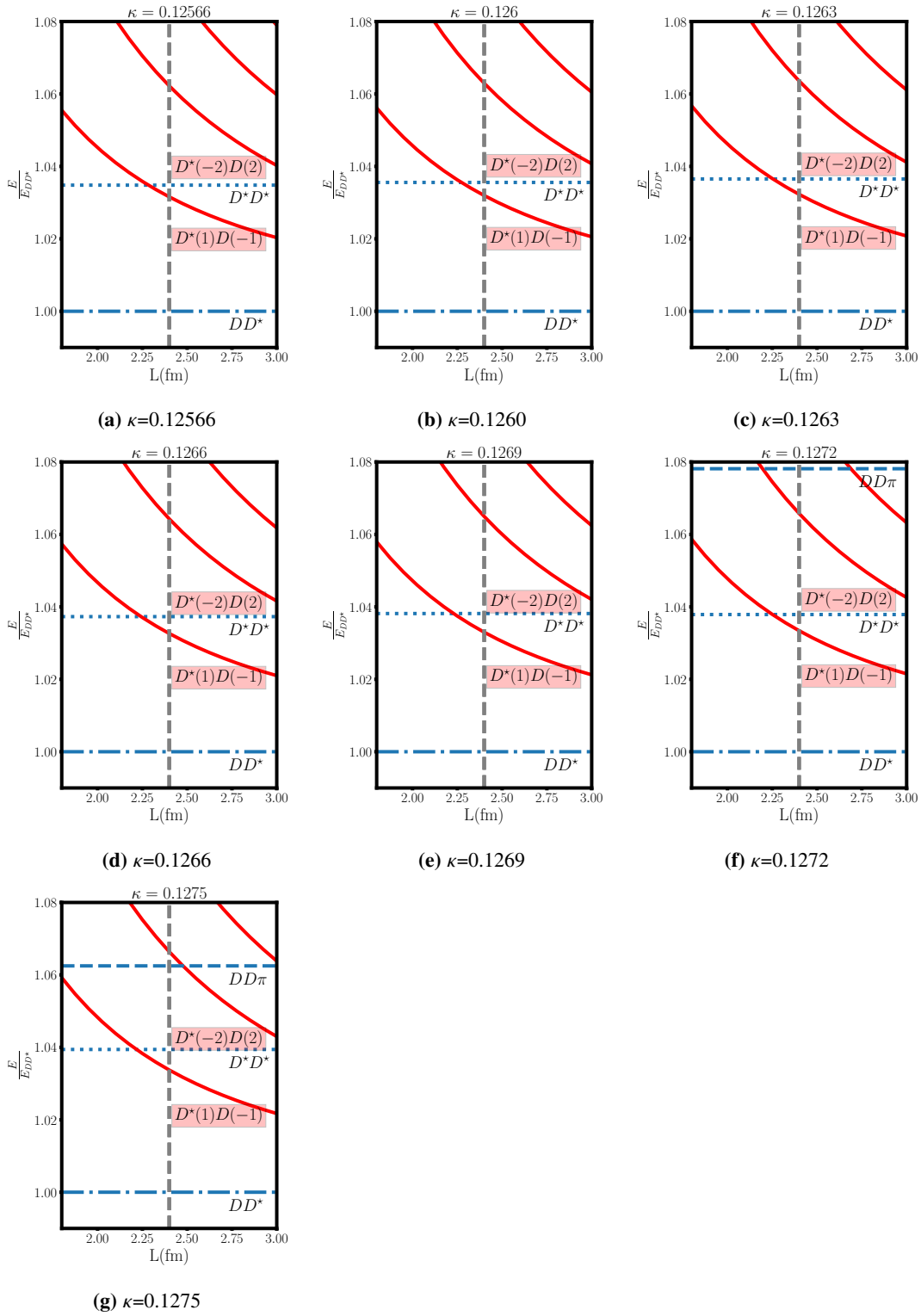
$m_h$	$m_0$	$\nu$	$C_p$	$\overline{M}_{hs}$	$\Delta M_{hs}$	$c^{\text{latt}}$
1.273	1.462	1.351	1.976	2076.8(9)	144.6(9)	0.992(25)
2.00	4.009	2.036	2.710	3020.5 (9)	93.5 (9)	1.021 (23)
2.73	8.499	2.951	3.922	3853.0 (11)	69.8 (11)	1.004 (27)
3.45	16.244	4.461	5.953	4631.4 (13)	57.5 (13)	1.001 (29)
4.183	30.061	7.041	9.014	5403.0(14)	48.3(14)	1.006 (30)

**Table 3:** RHQ action parameter tuned on  $16^3 \times 48$  ensemble. The  $\overline{M}_{hs}$  and  $\Delta M_{hs}$  are in MeV.



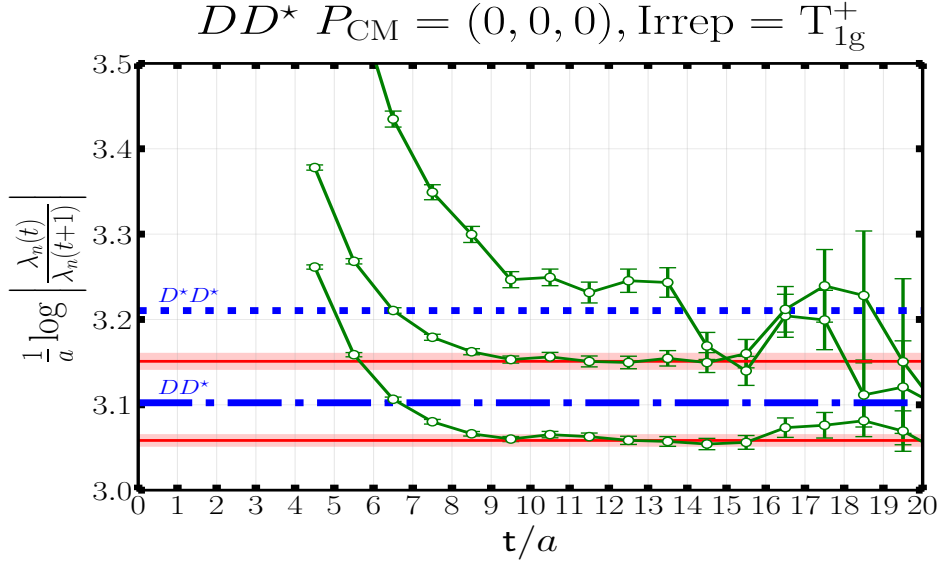
**Figure 2:** Dispersion relation of  $D_s$  and  $D_s^*$  meson obtained on  $16^3 \times 48$  ensemble

### 5. Results



**Figure 3:** Variation of thresholds with changing light quark  $\kappa$  from  $m_\pi = 688.8 (= m_{\eta_s})$  to  $m_\pi = 400$  MeV.

In figure 3 we show the trajectory of thresholds for different  $\kappa$ 's starting from the onset of elastic  $D^*D$  threshold up until the first inelastic  $D^*D^*$  threshold. This is the relevant domain of the Lüscher formalism where it is valid. We also depict the back-to-back non-interacting  $D^*D$  energy level which is in the relevant domain. The vertical line denotes the  $16^3 \times 48$  lattice volume we use in our calculation. When we perform our GEVP analysis we expect to find the energy levels close to the intersection points of the vertical line. The presence of valid non-interacting level in the relevant domain confirms that we have the necessary finite volume setup to reliably perform the  $T_{cc}$  analysis and extract its pole trajectory from Lüscher formalism.



**Figure 4:** Energy levels of the GEVP matrix for light quark  $\kappa = 0.12566$ .

In Fig. 4 we present our preliminary GEVP results on the  $16^3 \times 48$  ensemble with light quark  $\kappa = 0.12566$ . The correlator matrix has been constructed after projecting the correlators into the  $T_1^+$  irrep. We use time reversal symmetry to average the forward and backward running principal correlators. The principal correlators are fitted with single exponential ansatz to obtain the finite volume spectrum. We plot the 3 lowest lying levels below the inelastic  $D^*D$  threshold. The fit results are consistent with Fig. 3 in Ref. [17]. Following the Ref. [17], we can discard the third energy level close to the  $D^*D^*$  threshold. The energy shift from the non-interacting energy levels are a clear indication of the presence of a interaction between the  $D$  and  $D^*$ .

## 6. Summary and Outlook

In this work we study the spectrum of doubly charm tetraquark  $T_{cc}$  using MILC  $N_f = 2 + 1 + 1$  HISQ configurations. We include three set of operators, namely diquark-antidiquark, molecular and scattering and perform GEVP analysis on  $5 \times 5$  correlation matrix. For actions, we use Wilson-Clover for up/down and strange and anisotropic clover RHQ for charm. We find that for  $\kappa = 0.12566$  ( $m_\pi = 688.8$  MeV) the ground state of  $T_{cc}$  lies below the  $D^*D$  threshold, the first and the second excited states appear to lie below and above the  $D^*D^*$  thresholds. The ground state

likely indicates an attractive potential. It would be interesting to see what happens at lower pion mass  $\kappa = 0.1275$  ( $m_\pi = 400.4$  MeV) and at heavier quark mass points.

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