

Exotic vector charmonium and its leptonic decay width (Postprint)

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Abstract

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Full Text

Preamble

Exotic vector charmonium and its leptonic decay width

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We propose a novel type of interpolating field operators that manifest a hybrid-like configuration in which the charm quark-antiquark pair recoils against glu-

onic degrees of freedom. A heavy vector charmonium-like state with a mass of 4.33(2) GeV is disentangled from conventional charmonium states in the quenched approximation. This state shows affinity for the hybrid-like operators but couples less to the relevant quark bilinear operator. We also attempt to extract its leptonic decay constant and provide a tentative upper limit that it is less than one tenth of that of J/ψ , which corresponds to a leptonic decay width of about dozens of eV. The connection of this state with X(4260) is also discussed.

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Introduction

The X(4260) resonance was observed by many experiments as a J/ψ structure in the initial state radiation (ISR) process $e^+e^- \rightarrow \text{ISR}J/\psi$ [1-3]. Its resonance parameters are now determined to be $M_X = 4251(9)$ MeV and $\Gamma_X = 120(12)$ MeV [4]. According to its production mode, X(4260) must have the quantum numbers $J^{PC} = 1^{--}$. In addition, the ratio of $0\ 0J/\psi$ events observed by the CLEO collaboration is consistent with X(4260) being an isoscalar. In other words, X(4260) has the same quantum numbers as vector charmonia such as J/ψ and $\psi(2S)$. However, in contrast to J/ψ states, X(4260) has not been observed directly in e^+e^- annihilation. On the other hand, its mass is well above the D^+D^- threshold, but it has been observed only in the $J/\psi + \pi^+\pi^-$ system instead of D^+D^- channels. These facts may imply that X(4260) has a large branching fraction for the $J/\psi + \pi^+\pi^-$ decay mode. Thus, the small combined width $\Gamma(X(4260) \rightarrow e^+e^-)\text{Br}(X(4260) \rightarrow J/\psi) = 9.2 \pm 1.0$ eV can be understood as X(4260) having a very small e^+e^- width.

These features motivate conjectures that X(4260) might be an exotic state, for example, a hybrid charmonium [5-7]. In any case, more theoretical information about X(4260) is needed to unravel its nature, among which the leptonic decay width of X(4260) is an important quantity.

Regarding hybrid charmonium, extensive lattice QCD studies have been devoted to the $J^{PC} = 1^{--}$ channel. In the constituent quark model picture, this quantum number cannot appear in a $q\bar{q}$ system; therefore, it is usually conjectured that additional degrees of freedom should be involved, with the minimal configuration being $q\bar{q}g$ where g is a constituent gluon. The corresponding $q\bar{q}g$ interpolating fields are used in lattice calculations, which predict the mass of the 1^{--} charmonium-like state to be around 4.3 GeV [8-17]. Similar studies have also been extended to investigate possible hybrids with conventional quantum numbers, but the challenging task is to distinguish these states from conventional mesons. The state-of-the-art approach for this goal is the variational method based on large enough operator sets built through sophisticated methods. In the vector channel, a state has been observed with a mass around 4.4 GeV [14], which couples weakly to the quark bilinear operator but seems intimate with quark-antiquark-gluon operators.

Generally speaking, the appearance of an interpolating field operator does not

necessarily reflect the inner structure of a hadron state. However, for heavy quark systems where the non-relativistic picture may be available to some extent, the coupling of the operator to a specific state may bear useful information about its nature. Taking a mesonic hybrid as an example, even though it is an ambiguous concept from the point of quantum chromodynamics (QCD), it is always thought of as a hadron state made up of a quark-antiquark pair plus a constituent gluonic component in the constituent quark model picture. Of course, one can also relax the definition of a hybrid to an exotic object that has additional degrees of freedom apart from the constituent quarks. This kind of additional degree of freedom can be a fluctuating flux tube in the flux tube model, the color bag of the MIT bag model, etc. In any case, the essence of the exotic nature of a meson state is that the constituent $q\bar{q}$ pair acquires a center-of-mass motion by recoiling against the additional degrees of freedom, which is distinct from conventional hadron states. This is our starting point to build a novel type of hybrid-like operators. We split the charm quark-antiquark pair component and the gluon field in the $q\bar{q}g$ operator into two parts with different spatial separations. In momentum space, this manifests the center-of-mass motion of the $q\bar{q}$ pair in the rest frame of the state. We calculate the correlation functions of these operators, from which we try to extract the possible exotic charmonium state. Since the operators with different spatial separation provide different correlation functions, we fit them simultaneously along with the correlation function involving the electromagnetic current to obtain the decay constants of the states that contribute significantly.

This paper is organized as follows: Section II contains the description of the construction of the new lattice interpolation operators for the hybrid-like vector meson. The lattice parameters and numerical techniques are presented in Sec. III. We discuss our results and their connections with $X(4260)$ in Sec. IV. The conclusions and summary can be found in Sec. V.

II. New Interpolation Field for Exotic Vector Charmonium

In this work, we focus on the exotic vector charmonium ($J^PC = 1^{--}$) by assuming a hybrid-like configuration $c\bar{c}g$. A simple and straightforward local operator possibly reflecting this constituent configuration is $O_i^{\hat{H}}(x) = \bar{c}_a(x) \gamma_5 c_b(x) B_i^{\hat{ab}}(x)$, where a, b are color indices, i is the spatial index, and $B_i^{\hat{ab}}$ is the chromomagnetic field tensor. This kind of operator can be compared with the commonly used quark bilinear operator for the vector meson $O_i^{\hat{M}}(x) = \bar{c}_i c(x)$. In order to find the nonrelativistic form of these interpolation operators, we use the Foldy-Wouthuysen-Tani transformation [18] to decompose the charm quark and antiquark fields (Dirac spinors) in terms of the Pauli spinors χ / χ^\dagger which annihilates/creates a charm quark, and $\bar{\chi}$ / $\bar{\chi}^\dagger$ which creates/annihilates a charm antiquark. The explicit expressions of the operators $O_i^{\hat{H}}(x)$ and $O_i^{\hat{M}}(x)$ to the lowest order of the nonrelativistic approximation can be written as:

$$O_i^{\hat{H}}(x) = \bar{\chi}_a^\dagger \chi_b B_i^{\hat{ab}} + O(1/m_c) \quad O_i^{\hat{M}}(x) = \bar{\chi}_a^\dagger \chi_a +$$

$O(1/m_c)$

where one can see that the block $\bar{c}_a \gamma_5 c_b$ of the $O^{\wedge}(H)$ operator is a spin singlet and color octet, while that of $O^{\wedge}(M)$ is a spin triplet and color singlet. Intuitively, $O^{\wedge}(H)$ couples more to a state of spin singlet charm quark-antiquark pair and less to a state of spin triplet $c\bar{c}$ component owing to the heavy quark mass suppression for spin flipping of a heavy quark, and vice versa for $O^{\wedge}(M)$.

In order to resemble the center-of-mass motion of the $c\bar{c}$ recoiling against an additional degree of freedom, we split the operator $O^{\wedge}(H)$ into two spatial parts, $\bar{c}_a \gamma_5 c_b$ and $B_i^{\wedge}\{ab\}$, separated by an explicit spatial displacement r . In a fixed gauge (the Coulomb gauge in this work), we get a set of spatially extended operators:

$$O_i^{\wedge}(H)(x, t; r) = (\bar{c}_a \gamma_5 c_b)(x, t) B_i^{\wedge}\{ab\}(x + r, t).$$

It is expected that the coupling of this type of operator to conventional charmonia (without the center-of-mass motion of the charm quark-antiquark pair in the nonrelativistic picture) would be suppressed, while the coupling to the exotic state can be enhanced.

III. Numerical Details

We use the tadpole-improved gauge action [19-21] to generate gauge configurations on anisotropic lattices with the temporal lattice spacing much finer than the spatial one. The aspect ratio takes $\beta = a_s/a_t = 5$, where a_s and a_t are the spatial and temporal lattice spacings, respectively. Two lattices $L^3 \times T = 8^3 \times 96$ ($\beta = 2.4$) and $12^3 \times 144$ ($\beta = 2.8$) with different lattice spacings are used to check discretization artifacts, and the relevant input parameters are listed in Table I, where a_s values are determined from $r_0^{-1} = 410(20)$ MeV. We use the tadpole-improved clover action to calculate the quark propagators. The relevant parameters in the fermion action are tuned carefully by requiring that the physical dispersion relations of vector and pseudoscalar mesons are correctly reproduced at each bare quark mass [22, 23].

The bare charm quark masses at different β are determined from the physical mass of J/ψ , $m_{J/\psi} = 3.097$ GeV. The spatial extension of both lattices is ~ 1.7 fm, which is tested to be large enough for charmonium states. The ground state masses of 1S and 1P charmonia calculated on these two lattices (see Fig. 2 [Figure 2: see original paper] and Table II of Ref. [24] for details) show that finite a_s effects are small. Since the spatially extended interpolation operators $O^{\wedge}(H)$ discussed above are gauge variant, we carry out the calculation of the quark propagators and correlation functions after transforming each configuration to the Coulomb gauge.

A. Data analysis strategy

Our first task is to verify the existence of the exotic vector charmonium. We use the following source operator to calculate the correlation functions:

$$O_i^{\hat{W}}(y) = \bar{c}_a(y) \gamma_5 B_i^{\{ab\}}(z) c_b(z),$$

where y refers to the source time slice. For the sink operator $O_i^{\hat{H}}(x, t; r)$, the two-point functions we calculate are:

$$C^{\hat{H}}(r, t; i) = \sum_{\{x,y,z,i\}} \text{Tr}[S_F^\dagger(x, t; y) \gamma_5 S_F(x, t; z) B_i^\dagger(z, i)] \times B_i(x+r, t),$$

where $S_F(x, y)$ stands for the charm quark propagator. Accordingly, there are two types of wall-source quark propagators to be calculated. One uses the usual wall source by setting the source element to unity at each spatial site of the source time slice. The other uses the source by multiplying the chromomagnetic field tensor $B_i(z, i)$ to each site of the plain wall source.

In order to increase statistics, for each configuration we calculate T charm quark propagators $S_F(\vec{x}, t; \vec{0}, i)$ by setting the corresponding source vectors on each time slice i . This permits us to average over the temporal direction when calculating the two-point functions. In practice, the two-point functions $C^{\hat{H}}(r, t; i)$ with the same $|r| = r$ are averaged, such that the quantum number is kept to be $J^{PC} = 1^{--}$. After averaging over the time direction, the practical two-point functions we calculate are:

$$C^{\hat{H}}(r, t) = (1/TN_r) \sum_{\{i, |r|=r\}} C^{\hat{H}}(r, t + i; i),$$

where N_r is the degenerate degree of r . At the data analysis stage, we perform a simultaneous multi-exponential fit to $C^{\hat{H}}(r, t)$ using a correlated minimal- χ^2 fit method with the jackknife covariance matrix (we use three mass terms throughout this work).

B. Masses and Bethe-Salpeter amplitudes of conventional vector charmonia

In order to test the reliability of the fitting strategy mentioned above, we first carry out a similar analysis of the correlation functions $C^{\hat{M}}(r, t)$ of the spatially extended version of operator $O^{\hat{M}}$ on the $a = 2.4$ lattice. The procedure is detailed as follows. The spatially extended version of $O^{\hat{M}}$ is defined as:

$$O_i^{\hat{M}}(x, t; r) = \bar{c}(x, i) \gamma_i c(x+r, i),$$

and $C^{\hat{M}}(r, t + i; i)$ is defined similarly as in Eq. (4). After averaging over the temporal direction, we have:

$$C^{\hat{M}}(r, t) = (1/TN_r) \sum_{\{i, |r|=r\}} C^{\hat{M}}(r, t + i; i) = \sum_i \Phi'_i(r) e^{-m_i t},$$

where $\Phi'_i(r)$ is the r -dependent spectral weight of the i -th state. In the fitting procedure, we fix a maximal t (denoted by t_{max}) and vary the lower bound t_{min} of the fit window, then we obtain the masses of the lowest three states which keep constant to some extent for a series of t_{min} , as shown in the upper panel of Figure 1 [Figure 1: see original paper]. We average the masses in this range with each value weighted by its error and get the values $m_1 = 3.097(1)$

GeV, $m_2 = 3.679(19)$ GeV, and $m_3 = 4.007(57)$ GeV, respectively. These three states may correspond to J/ψ , $\psi(3686)$, and $\psi(4040)$, and we can almost reproduce their experimental spectrum.

We also plot the $\Phi'_i(r/a_s)$'s (normalized as $\Phi'_i(0) = 1$) in the right panel of Figure 1, where one can find that there is no radial node for $\Phi'_1(r/a_s)$, one radial node for $\Phi'_2(r/a_s)$, and two radial nodes for $\Phi'_3(r/a_s)$. Given the quark model assignments n^3S_1 state for J/ψ , $\psi(3686)$, and $\psi(4040)$ with $n = 1, 2, 3$, respectively, this is actually not surprising, since $\Phi'_i(r)$ is proportional to the Coulomb Bethe-Salpeter amplitude of the i -th S-wave charmonium, which, at the leading order of the nonrelativistic approximation, corresponds to the radial wave functions in the quark model [25, 26].

C. Existence of an exotic vector charmonium and its properties

From the above, one can see that our data analysis strategy is robust for conventional vector charmonium states; therefore, we perform a similar study for the correlation functions described in Eq. (4) and (5). Figure 2 [Figure 2: see original paper] shows the plots of $\Phi_i(r)$ with respect to r (in physical units) through a three-mass-term fit (the upper panel is for $\tau = 2.4$ at $t_{\min} = 12a_t$, and the lower panel for $\tau = 2.8$), whose masses are fitted to be 3.100(7) GeV, 3.58(9) GeV, and 4.6(2) GeV for $\tau = 2.4$, and 3.090(6) GeV, 3.54(5) GeV, and 4.6(1) GeV for $\tau = 2.8$ at $t_{\min} = 16a_t$. $\Phi_1(r)$ and $\Phi_2(r)$ damp more rapidly and are close to zero near $r = 0.3$ fm, while $\Phi_3(r)$ is still relatively large. The lowest two states correspond very possibly to the conventional vector charmonia J/ψ and ψ according to their masses. In contrast, the third state, with a much higher mass, still dominates the two-point functions with r larger than 0.3 fm. This may signal the exotic nature of this state that is reflected by the spatially extended sink operator $\hat{O}(H)$. Of course, higher conventional vector charmonia, such as $\psi(4040)$ and $\psi(4415)$, should also contribute to the two-point function $\hat{C}(H)(r, t)$; however, in our data analysis procedure, $\hat{C}(H)(r, t)$ cannot accommodate more statistically meaningful states. The presence of higher conventional charmonia may result in small shifts of the masses of the fitted states—for example, the mass of the second state deviates from that of the would-be ψ state.

It is seen from Fig. 2 that the r -behaviors of the first (J/ψ) and the second state (ψ) are similar up to an overall factor. If this is the case for all conventional charmonium states, since the r behavior may depict the center-of-mass motion of the $c\bar{c}$ component and conventional charmonia are free of this in the nonrelativistic approximation, we can conjecture that the spectral weights of the i -th conventional charmonia $\Phi_i(r)$ can be factorized into $\Phi(r)W_i$, where $\Phi(r)$ is approximately uniform and insensitive to the different conventional charmonia, such that the two-point functions $\hat{C}(H)(r, t)$ with different r can be linearly combined to eliminate the contribution from conventional charmonium states.

In practice, we combine linearly the correlation functions $\hat{C}(H)(r, t)$ at two

specific r_1 and r_2 as:

$$C(\alpha, t) = C^{\wedge}(\mathbf{H})(r_1, t) - C^{\wedge}(\mathbf{H})(r_2, t),$$

where α is a tunable parameter. For each lattice, an optimal α can be obtained by requiring that the effective mass plateau of $C(\alpha, t)$ is as long as possible. To be specific, for $\beta = 2.4$, we use $r_1 = 0$ and $r_2 = a_s$ and set $\alpha = 2.576$, which is very close to the central value of the ratio $\Phi_1(0)/\Phi_1(a_s) = 2.583$ from the three-mass fit illustrated in Fig. 2. Using this α and the fit results in Fig. 2, the spectral weights of the three states in $C(\alpha, t)$ are roughly 0.002, -0.1, and -1.1, respectively. This implies that the relative contribution of the third state is strongly enhanced by this subtraction scheme. Similarly, for $\beta = 2.8$ we use $r_1 = a_s$ and $r_2 = \sqrt{3}a_s$ and set $\alpha = 3.658$, which is also close to the central value of the ratio $\Phi_1(a_s)/\Phi_1(\sqrt{3}a_s) = 3.681$. The spectral weights of the three states are roughly 0.007, -0.2, and -1.9, respectively.

In Figure 3 [Figure 3: see original paper], we plot the effective mass plateaus of $C(\alpha, t)$ for $\beta = 2.4$ and $\beta = 2.8$. The time and masses are expressed in physical units according to the lattice spacings listed in Table I. One can see that both plateaus are fairly good and lie on top of each other. Since both the sink operator $O^{\wedge}(\mathbf{H})(\mathbf{r})$ and the source operator $O^{\wedge}(\mathbf{W})$ (whose correlation functions are $C^{\wedge}(\mathbf{H})(\mathbf{r}, t)$) are expected to couple strongly to hybrid-like states and the contribution from conventional charmonia is largely subtracted by the above scheme, we take the state reflected by the observed plateau as the exotic vector charmonium and name it X in the rest of this work (we keep the name of the experimental state $X(4260)$). The horizontal line shows the fitted mass $M_X = 4.33(2)$ GeV through a one-exponential fit in the time range from 0.3 to 0.9 fm.

To this end, we claim that a vector charmonium-like state with a mass of 4.33(2) GeV has been unambiguously singled out, whose exotic nature may be reflected by its distinct coupling to the special interpolation field $O^{\wedge}(\mathbf{H})(\mathbf{r}, t)$ in comparison with conventional vector charmonia. It should be noted that this state has also been observed by previous lattice studies using variational methods based on lattice operator sets [14]; however, the spatially extended operators $O^{\wedge}(\mathbf{H})(\mathbf{r}, t)$ we use provide a clearer picture of its inner structure.

D. Leptonic decay constant of the exotic vector charmonium

Since this hybrid-like charmonium can be disentangled from conventional charmonia with the prescription above, its leptonic decay constant can be investigated accordingly. The leptonic decay constant f_V of a vector meson state V is defined by:

$$\langle 0 | J_{\mu}^{\wedge}(\text{em}) | V(\sim \mathbf{p}, r) \rangle = m_V f_V \epsilon_{\mu}(\sim \mathbf{p}, r),$$

where $J_{\mu}^{\wedge}(\text{em})(0)$ is the electromagnetic current and $\epsilon_{\mu}(\sim \mathbf{p}, r)$ is the polarization vector of V at momentum $\sim \mathbf{p}$. For vector charmonium states, $J_{\mu}^{\wedge}(\text{em})(0)$ can be approximated by $\bar{c} \gamma_{\mu} c(0)$ if the contribution from other quark flavors through annihilation diagrams is neglected. Since the vector current $J_{\mu}^{\wedge}(\text{em})(x)$

defined in the continuum limit is no longer conserved on the lattice, we perform a nonperturbative renormalization procedure [27] to extract the multiplicative renormalization constant Z_V of the current. The renormalization constant of the spatial components of \hat{J}_{em} is determined to be $Z_V(s) = 1.39(2)$ for $\beta = 2.4$ and $Z_V(s) = 1.11(1)$ for $\beta = 2.8$ [28].

Since the spatial components of $\hat{J}_{em}(x)$ are exactly the normal quark bilinear operator $\hat{O}(M)$ for vector mesons, the matrix elements in Eq. (10) can be derived from the corresponding correlation functions involving the operator $\hat{O}(M)$ along with the vector current renormalization constant $Z_V(s)$. In order to obtain these matrix elements, we also calculate two other categories of correlation functions in addition to $\hat{C}(H)(r, t)$:

$$\hat{C}(J)(t) = \langle \hat{J}_i(x, t) \hat{O}^\dagger(W) \rangle, \quad \hat{C}(W)(t) = \langle \hat{O}^{(W)(t)} \hat{O}^\dagger(W) \rangle,$$

where averaging over the temporal direction is also taken implicitly in the above expressions. After inserting intermediate states into $\hat{C}(J)(t)$, $\hat{C}(H)(r, t)$, and $\hat{C}(W)(t)$, we have:

$$\hat{C}(J)(t) = \sum_n Z_n^{(J)Z_n(W)*} e^{-m_n t}, \quad \hat{C}(H)(r, t) = \sum_n Z_n^{(H)(r)Z_n(W)*} e^{-m_n t}, \quad \hat{C}(W)(t) = \sum_n |Z_n^{(W)}|^2 e^{-m_n t},$$

where m_n is the mass of the n -th state and the parameter $Z_n(K)$ with K referring to H or W is defined as:

$$Z_n(K) = \langle 0 | \hat{O}(K) | V_n(p, r) \rangle / (2m_n V_3).$$

Accordingly, the leptonic decay constant f_{Vn} can be derived from $Z_n(J)$ from the definition Eq. (10) as:

$$f_{Vn} = C Z_V(s) Z_n(J) / m_n,$$

where C is an overall constant prefactor owing to the redefinition of our quark fields and the anisotropic lattices we are using.

The time dependence of the correlation functions is usually observed from their effective mass plots. The effective masses of different correlation functions $\hat{C}(X)(t)$ are defined as:

$$M_{\text{eff}}(t)_a = \log[\hat{C}^{(K)(t)/C}(K)(t+1)],$$

where K stands for J , H , and W , and are illustrated in Fig. 4 [Figure 4: see original paper], where the mass of the ground state ($J/$) is also plotted as a horizontal dashed line to guide the eye. The left panel of Fig. 4 is for the $\beta = 2.4$ lattice, and the right panel is for $\beta = 2.8$. In the plots, we present in the third row the effective mass of $\hat{C}(H)(r, t)$ at a specific r ($r = \sqrt{6}a_s$ for $\beta = 2.4$ and $r = 3a_s$ for $\beta = 2.8$, respectively), where $\hat{C}(H)(r, t)$ is dominated by the third state in the short time range, changes sign, and is finally saturated by the ground state when t increases. This is manifested in the effective mass plot by the phenomenon that $M_{\text{eff}}(t)$ shows a meta-stable plateau (roughly 4.4 GeV) higher than the ground state before the discontinuity time, and then converges

to the mass of the ground state. This phenomenon also implies that the spectral weight of the third state is much larger than those of the lower states and the signs of the two spectral weights are different.

In the fitting procedure, we fix the maximal time t_{\max} of the fitting window ($t_{\max} = 30a_t$ for $\beta = 2.4$ and $t_{\max} = 40a_t$ for $\beta = 2.8$), and let the minimal time vary in a range. The fitted spectral weights of $C^{\hat{J}}(t)$ and $C^{\hat{W}}(t)$ are listed in Table II (raw data), from which the decay constants can be derived. The fitted masses and decay constants are converted to values in physical units (GeV) and are presented in Table III. The spectral weights of $C^{\hat{H}}(r, t)$ are less relevant and omitted here to save space (one can refer to Fig. 2 to see their relative magnitudes for the three states). All errors are statistical and obtained through a jackknife analysis. To illustrate the fit quality, we also show the effective masses (in blue bands) of the functions in Eq. (12) using the fitted parameters at $t_{\min} = 12a_t$ for $\beta = 2.4$ and $t_{\min} = 16a_t$ for $\beta = 2.8$. It is seen that the fit functions describe the measured data very well. Actually, for all the fits, the $\chi^2/N_{\text{d.o.f}}$ values are around 1 even though the number of degrees of freedom $N_{\text{d.o.f}}$ is always several hundred.

As shown in Table II and Table III, the fitted parameters are almost stable and insensitive to t_{\min} . The masses of the first and second states are consistent with those of J/ψ and ψ' , while the mass of the third state is a little higher than the hybrid-like state we obtained before. This can be attributed to contamination from even higher states to some extent. In the right part of Table II, the spectral weights $|Z_3^{(W)}|^2$ are an order of magnitude larger than those of the lowest two states (corresponding mostly to J/ψ and ψ'). This is not strange since $Z_n^{\hat{W}}$ represents the coupling of the hybrid-like wall-source operator $O^{\hat{W}}$ to the n -th state and is expected to be enhanced when coupling to a hybrid-like state. In contrast, the spectral weights $Z_n^{(J)Z_n^{\hat{W}}}$ (in the left part of Table II) of the lowest two states, even suppressed by $Z_n^{\hat{W}}$, are much larger than that of the third state, which are close to zero within errors. This may imply that the decay constant of the third state is very small, since $Z_n^{\hat{J}}$ is proportional to the decay constant of the n -th state.

Using the fitted spectral weights $Z_n^{(J)Z_n^{\hat{W}}}$ and $|Z_n^{(W)}|^2$, we can obtain the concrete values of $Z_n^{\hat{J}}$, from which the decay constant of the n -th state can be derived using Eq. (14). The derived decay constants of the three states are also listed in Table III, where the last row lists the experimental values for comparison. For J/ψ , we obtain its decay constant to be roughly 0.46(2) GeV at $\beta = 2.4$ and 0.43(2) GeV at $\beta = 2.8$, which are close to the experimental value although 5%-10% larger. The deviation can be attributed to artifacts of finite lattice spacing (and also uncertainty due to the quenched approximation). The derived decay constant of ψ' , $f_{\psi'}$, seems compatible with the experimental value but with huge errors that come mainly from the uncertainty of $Z_n^{\hat{W}}$.

The most striking observation is that the decay constant f_3 of the third state is consistent with zero within error. Superficially, it seems that the exotic vector charmonium has a nearly zero decay constant. However, there is a possibility

that this is a mixing effect of two nearby states (for example, the would-be exotic state and (4415)), whose contributions to $C^{\wedge}(J)(t)$ cancel to some extent because we perform the simultaneous fit using only three mass terms. We have established that an exotic vector charmonium state does exist and contributes substantially to the correlation functions $C^{\wedge}(H)(r, t)$ in the previous context. This is also the case for $C^{\wedge}(W)(t)$ since $|Z_{-3}^{(W)}|^2$ is one magnitude or even larger than $|Z_{-1,2}^{(W)}|^2$. Based on these facts and considering the possible admixture of a conventional charmonium state to the exotic state, we try to estimate the upper limit of the decay constant of the exotic vector charmonium.

If the third state actually receives contributions from the would-be exotic state X and the adjacent vector charmonium state (4415), then the spectral weight $Z_{-3}^{(J)Z_{-3}}(W)^*$ can be expressed as:

$$Z_{-3}^{(J)Z_{-3}}(W)^* = Z_{-X}^{(J)Z_{-X}}(W)^* + Z_{-(4415)}^{(J)Z_{-(4415)}}(W)^* - Z_{-X}^{(J)Z_{-X}}(W)^*,$$

since $|Z_{-(4415)}^{\wedge}(W)| = |Z_{-X}^{\wedge}(W)|$. Thus we have:

$$|Z_{-X}^{\wedge}(J)| = |Z_{-3}^{(J)Z_{-3}}(W)^*| / |Z_{-X}^{\wedge}(W)|.$$

According to Eq. (14), this is equivalent to:

$$f_{-X} = (C Z_{-V}^{\wedge}(s) |Z_{-3}^{(J)Z_{-3}}(W)^*|) / (m_{-X} |Z_{-X}^{\wedge}(W)|).$$

If X(4260) is tentatively assigned to the X state investigated in this study, combining the above value with the leptonic decay width of X, we can give an estimate of the branching ratio of X(4260) decaying into $J/\psi + \pi$.

Since $|Z_{-X}^{\wedge}(W)| = |Z_{-(4415)}^{\wedge}(W)|$, we can take the approximation $|Z_{-X}^{\wedge}(W)| = |Z_{-3}^{\wedge}(W)|$. Furthermore, if we assume $|Z_{-(4415)}^{\wedge}(J)| = |Z_{-X}^{\wedge}(J)|$, from Table II we can take $|Z_{-(4415)}^{\wedge}(W)| = |Z_{-X}^{\wedge}(W)|$. Experimentally, the leptonic decay width of (4415) is measured to be $\Gamma((4415) \rightarrow e+e-) = 0.58(7)$ keV, and $f_{-(4415)}$ is extracted to be 157 MeV using the relation:

$$\Gamma(V_{-c}^{-}c \rightarrow e+e-) = (4 \alpha_{\text{QED}}^2 f_{-V}^2) / (3m_{-V}),$$

where we take $\alpha_{\text{QED}} = 1/134$ at the charm quark mass scale. Therefore, f_{-X} can be roughly estimated to be:

$$f_{-X} = (|Z_{-3}^{(J)Z_{-3}}(W)| / |Z_{-(4415)}^{\wedge}(W)|) \times (m_{-} / m_{-X}) \times f_{-(4415)} \approx 30 \text{ MeV}.$$

So we can assign a safer upper limit of f_{-X} as:

$$f_{-X} < 40 \text{ MeV},$$

which gives the upper limit of the leptonic decay width of the exotic vector charmonium:

$$\Gamma(X \rightarrow e+e-) < 40 \text{ eV}.$$

IV. Discussion

The surprisingly small e^+e^- decay width of the exotic vector charmonium X is in sharp contrast to that of conventional vector charmonia, which are usually of the order of keV. In other words, if this hybrid-like vector charmonium does exist in the real world, its contribution to the inclusive cross section of e^+e^- annihilation is rather small. Actually, the R-value scan versus the invariant energy \sqrt{s} of e^+e^- collision does not show any indication of an extra vector charmonium-like state around $\sqrt{s} = 4.3$ GeV (there is, however, a small dip in this energy range). The BEPCII/BESIII in Beijing is now accumulating data from e^+e^- collisions in this energy range and will hopefully give a more precise line shape of the R-value here to clarify the situation.

On the other hand, as mentioned above, the vector charmonium state $X(4260)$ was observed by several experiments in the initial state radiation of e^+e^- annihilation into $J/\psi + \pi$. The combined decay width of $X(4260)$ is:

$$\Gamma(X(4260) \rightarrow e^+e^-)\text{Br}(X(4260) \rightarrow J/\psi) = 9.2 \pm 1.0 \text{ eV}.$$

If $X(4260)$ is tentatively assigned to the X state investigated in this study, combining the above value with the leptonic decay width of X , we can give an estimate of the branching ratio of $X(4260)$ decaying into $J/\psi + \pi$:

$$\text{Br}(X(4260) \rightarrow J/\psi) > 20\%,$$

which means that J/ψ is one of the most important decay modes of $X(4260)$. This can naturally explain why $X(4260)$ has been observed only in this channel till now. Furthermore, given the likely hybrid nature of the X state, it can be expected that the spin singlet $c\bar{c}$ component of X prefers a hadronic transition into spin singlet charmonium, such as h_c . So $X(4260) \rightarrow h_c$ can also be an important decay mode of $X(4260)$. Recently, the BESIII Collaboration studied the $e^+e^- \rightarrow \pi^+ - h_c$ process at center-of-mass energies from 3.90 GeV to 4.42 GeV. They found that the cross sections are of the same order of magnitude as, but have different line shapes from, those of $e^+e^- \rightarrow \pi^+ - J/\psi$ [29]. Anyway, it is highly desirable to investigate whether they are from the same resonance structure or not.

The reason for the large branching ratio of $X(4260) \rightarrow J/\psi$ can be depicted as follows. In e^+e^- annihilation, the charm quark-antiquark pair $c\bar{c}$ is produced in the short range through the virtual photon. During the hadronization procedure, the charm quark and antiquark emit soft gluons continuously, which form a colored gluon halo around the gradually localized color octet $c\bar{c}$ (in a relative sense). Finally, a meta-stable state is formed as the $X(4260)$ particle. Obviously, the color octet $c\bar{c}$ kernel can be readily neutralized into a color singlet charmonium by absorbing (emitting) soft gluons from (to) the halo, and the gluon halo thereby becomes color neutral and is emitted as light hadrons—for example, the π pair. If this is actually the case, the color flux between the charm quark and antiquark has less chance to be excited to a high enough energy to break, and thus the possibility of $D\bar{D}$ decay modes is suppressed.

There is a little similarity between this “halo charmonium” picture and the so-called “hadro-charmonium” picture, where the relatively localized color neutral $c\bar{c}$ kernel is surrounded by a light hadron cloud [30]. However, the advantage of the former resides in the fact that the direct color interaction between the halo and the kernel provides an obvious binding mechanism, while in the “hadro-charmonium” picture, more phenomenological assumptions are needed to describe the interaction between the meson cloud and the charmonium kernel.

V. Conclusion

To summarize, we use a new type of spatially extended hybrid-like operator to investigate the possible existence of exotic vector charmonia. In the non-relativistic approximation of these operators, the localized color octet charm quark-antiquark component is in the spin singlet state and separates from the chromomagnetic field strength with a spatial distance. These operators couple preferably to a higher vector state X with a mass of $4.33(2)$ GeV when the distance increases. This observation indicates that the charm quark-antiquark pair of X may acquire a center-of-mass motion by recoiling against additional degrees of freedom depicted by the chromomagnetic field strength operator, which are necessarily gluonic in the quenched approximation. In this sense, the state X can be taken as a hybrid-like vector charmonium.

In addition, through a simultaneous multi-state fit to different correlation functions built from the vector current operator and the hybrid operator mentioned above, the leptonic decay constant of X is tentatively determined to be roughly one order of magnitude smaller than $f_{J/\psi}$, say, $f_X < 40$ MeV, which gives a very small leptonic decay width $\Gamma(X \rightarrow e^+e^-) < 40$ eV. This is a very important characteristic parameter for X to be identified in experiments. Obviously, the mass and leptonic decay width of X are consistent with the production and decay properties of $X(4260)$, which has so far escaped direct measurement in e^+e^- annihilation. Based on the combined width $\Gamma_{ee} \Gamma_{J/\psi} + \dots / \Gamma_{\text{tot}} = 9.2$ eV of $X(4260)$, if it can be assigned to the X state in this study, its decay branching fraction for the J/ψ mode can be larger than 20%, which also naturally explains why $X(4260)$ is dominantly observed in J/ψ . By virtue of the inner structure of X , $X(4260)$ should also be observed in the h_c channel.

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