

Possible hadronic excess in $\psi(2S)$ decay and the $\rho\pi$ puzzle

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We study the so-called $\rho\pi$ puzzle of $\psi(2S)$ decay by incorporating two inputs: the relative phase between the one-photon and the gluonic decay amplitude, and a possible hadronic excess in the inclusive nonelectromagnetic decay rate of $\psi(2S)$. We look into the possibility that the hadronic excess in $\psi(2S)$ originates from a decay process of long-distance origin which is absent from the J/ψ decay. We propose that the amplitude of this additional process happens to nearly cancel the short-distance gluonic amplitude in the exclusive decay $\psi(2S) \rightarrow 1^-0^-$ and turn the sum dominantly real in contrast to the J/ψ decay. We present the general consequences of this mechanism and survey two models which might possibly explain the source of this additional amplitude.

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I. INTRODUCTION

The absence of the $\rho\pi$ decay mode of $\psi(2S)$ has defied a theoretical explanation for more than a decade [1]. The recent measurement by the BES Collaboration [2] has confirmed the absence of $\rho\pi$ with even a higher precision, setting its upper bound at a level of a factor of more than 60 below what one would naively expect from the decay $J/\psi \rightarrow \rho\pi$. The measurement of other decay modes by BES [2,3] seems to rule out all possible resolutions for the $\rho\pi$ puzzle that have so far been proposed by theorists [4]. For instance, the large $\omega\pi$ branching contradicts the helicity suppression [5] with or without large intrinsic charm [6]. A vector glueball near the J/ψ mass, if it should exist, can enhance the $\rho\pi$ branching for J/ψ relative to $\psi(2S)$ [7]. However, the magnitude of $B(J/\psi \rightarrow \rho\pi)$ is in line with expectations when we compare the $B(\rho\pi)/B(\omega\pi)$ with the inclusive ratio $B(J/\psi \rightarrow ggg \rightarrow X)/B(J/\psi \rightarrow \gamma^* \rightarrow X)$. What happens is not an enhancement of $\rho\pi$ in J/ψ but a suppression of $\rho\pi$ in $\psi(2S)$.¹

Meanwhile, the amplitude analysis of the J/ψ decay revealed that the relative phase of the gluonic and the one-photon decay amplitude is close to 90° for all two-body decay channels so far studied: 1^-0^- [8], 0^-0^- [9], 1^-1^- [10], and $N\bar{N}$ [11]. We show in this paper that the recent BES measurement in $J/\psi \rightarrow 1^+0^-$ is also compatible with a large phase.

In contrast, the pattern of a large relative phase does not emerge for $\psi(2S)$. Within experimental uncertainties, the relative phase is consistent with zero in 1^-0^- decay [12,13] and 1^+0^- decay. This marked difference between J/ψ and $\psi(2S)$ is another puzzle if three-gluon decay is equally responsible for the strong decay of J/ψ and $\psi(2S)$.

There is one more piece of experimental information relevant to the issue. That is the hadronic decay rate of $\psi(2S)$ which is normally attributed to $\psi(2S) \rightarrow ggg$. When we compute with the current data the inclusive gluonic decay rate of $\psi(2S)$ by subtracting the cascade and the electromag-

netic decay rate from the total rate, it is 60–70 % larger, within experimental uncertainties, than what we expect from short-distance gluonic decay alone. This excess hadronic branching in $\psi(2S)$ may suggest that something more occurs in the gluonic decay of $\psi(2S)$ than in the J/ψ decay.

In this paper we combine these pieces of information together and search for the origin of the marked difference between J/ψ and $\psi(2S)$. While we should be apprehensive about experimental errors at present, they might give us a clue to a solution of the $\rho\pi$ puzzle. In Sec. II, prompted by the experimental observation in the 0^-0^- , 1^-0^- , 1^-1^- , and $N\bar{N}$ channels, we postulate the universality of the large relative phase between the gluon and the photon decay amplitudes. Specifically, the gluonic decay amplitude acquires a large phase while the photon amplitude is real. We point out that a large phase is consistent with new BES data in $J/\psi \rightarrow 1^+0^-$. Further progress in the BES analysis in this channel will shed more light. We turn to $\psi(2S)$ in Sec. III. The decay branching fractions of $\psi(2S) \rightarrow 1^-0^-$ clearly show a suppression of the gluon amplitude and favor a small relative phase between the gluon and photon amplitudes. We point out that a small phase is more likely in $\psi(2S) \rightarrow 1^+0^-$ too. Taking the possible excess in the inclusive hadronic decay rate of $\psi(2S)$ seriously, we propose that this excess is related to both the suppression and the small relative phase of the 1^-0^- amplitude. Our proposition is that an additional decay process generating the excess should largely cancel the short-distance gluon amplitude in the exclusive decay into 1^-0^- and that the resulting small residual amplitude is not only real but also destructively interferes with the photon amplitude. In Sec. IV, we first present general consequences of the destructive interference. We then examine two scenarios which may possibly generate the excess inclusive gluonic decay. One is the contribution of the virtual $D\bar{D}$ intermediate state. The other is a resonance, a glueball, or a four-quark resonance, near the $\psi(2S)$ mass. Though neither idea is novel nor highly appealing, they seem to be among a very few possibilities that have not yet been ruled out by experiment.

II. PHASES OF J/ψ DECAY AMPLITUDES

The relative phase between the gluon and photon amplitudes in the decay $J/\psi \rightarrow 1^-0^-$ has been analyzed with bro-

¹More comparisons of experiment with models are found in Ref. [2].

ken flavor SU(3) symmetry [8] including the ρ - ω mixing. All analyses clearly show that the relative phase should be very large and not far from 90° with experimental uncertainties. The SU(3) analysis was made also for the 0^-0^- modes [9] and the 1^-1^- modes [10] for which the leading gluon amplitude is SU(3) violating. The relative phases were found to be equally large for these modes. Furthermore, comparison of the electromagnetic form factors in the timelike region with the J/ψ decay branching fractions revealed that the relative phase is very close to 90° in the $N\bar{N}$ decay channels too [11]. A question arises as to whether this large relative phase is universal to all decays of J/ψ or not. There is no persuasive theoretical answer to it at present.²

In addition to those two-body channels already analyzed, the recent BES measurement [3] on $J/\psi \rightarrow 1^+0^-$,

$$B(J/\psi \rightarrow K_1^\pm(1400)K^\mp) = (3.8 \pm 0.8 \pm 1.2) \times 10^{-3},$$

$$B(J/\psi \rightarrow K_1^\pm(1270)K^\mp) < 3.0 \times 10^{-3}, \quad 90\% \text{ C.L.}, \quad (1)$$

is relevant to this issue. We examine here these branching fractions together with the $b^\pm \pi^\mp$ branching fraction $B(J/\psi \rightarrow b^\pm \pi^\mp) = (3.0 \pm 0.5) \times 10^{-3}$ [15].³

Since $K_1(1270)$ and $K_1(1400)$ are superpositions of K_A and K_B of the 1_A^{++} and the 1_B^{+-} octet, respectively,

$$K_1^+(1400) = K_A^+ \cos \theta + K_B^+ \sin \theta,$$

$$K_1^+(1270) = -K_A^+ \sin \theta + K_B^+ \cos \theta, \quad (2)$$

with $\theta \approx 45^\circ$ [17], we can parametrize the three decay amplitudes in terms of the gluon amplitude a_1 of the 1_B^{+-} octet and the photon amplitudes $a_{\gamma A/B}$ of 1_A^{++} and 1_B^{+-} :

$$A(b_1^+ \pi^-) = a_1 + \sqrt{1/5} a_{\gamma B},$$

$$A(K_1^+(1270)K^-) = (a_1 + \sqrt{1/5} a_{\gamma B}) \cos \theta - a_{\gamma A} \sin \theta,$$

$$A(K_1^+(1400)K^-) = (a_1 + \sqrt{1/5} a_{\gamma B}) \sin \theta + a_{\gamma A} \cos \theta. \quad (3)$$

Since there are two independent helicity amplitudes (or s and d waves) for 1^+0^- , we should use this parametrization separately for the s -wave and the d -wave amplitudes. The three branching fractions can be fitted with

$$|a_1| > |a_{\gamma A}| \approx |a_{\gamma B}|,$$

$$\arg(a_1^* a_{\gamma A}) \approx \arg(a_1^* a_{\gamma B}) \approx 90^\circ. \quad (4)$$

²Some attempt was recently made to argue in favor of universal large phases [14].

³A previous analysis [16] of these J/ψ decay modes assumed a zero relative phase and used a preliminary value of the upper bound on $B(J/\psi \rightarrow K_1^\pm(1270)K^\mp)$. Therefore the J/ψ analysis of Ref. [16] should be disregarded. However, the analysis of the $\psi(2S) \rightarrow 1^+0^-$ of Ref. [16] remains valid.

If the ratio $\Gamma(J/\psi \rightarrow \gamma \rightarrow 1^+0^-)/\Gamma(J/\psi \rightarrow ggg \rightarrow 1^+0^-)$ is comparable in magnitude to $\Gamma(J/\psi \rightarrow \gamma \rightarrow X)/\Gamma(J/\psi \rightarrow ggg \rightarrow X) \approx 1/5$, we should expect that $|a_{\gamma A/B}| \approx 0.7|a_1|$ [16]. For $a_{\gamma A} = a_{\gamma B} = \pm 0.7i a_1$ and $\theta \approx 45^\circ$, the ratios of the branching fractions prior to the phase space corrections, denoted by B_0 , take values as

$$B_0(b^\pm \pi^\mp) : B_0(K_1^\pm(1270)K^\mp) : B_0(K_1^\pm(1400)K^\mp) \\ \approx 1 : 0.5 : 0.9. \quad (5)$$

While the inequality $B(K_1^\pm(1270)K^\mp) < B(K_1^\pm(1400)K^\mp)$ can easily be realized by a wide range of parameter values, the other inequality $B(b^\pm \pi^\mp) < B(K_1^\pm(1400)K^\mp)$ is a little tight. If we allow $|a_{\gamma A/B}|$ larger than $0.7|a_1|$ and/or increase the value of θ , however, the current central values of the branching fractions are consistent with the large phase hypothesis. We should also point out that if the SU(3) breaking correction is made by the meson wave functions $(f_P f_A)^2$, it is likely to enhance $B(K_1^\pm(1400)K^\mp)$ over $B(b^\pm \pi^\mp)$ to the direction in favor of the large phase fit.

If we leave $a_{\gamma A/B}$ unrestricted in magnitude and phase, a triangular relation holds for the amplitudes as

$$A(K_1^\pm(1270)K^\mp) \cos \theta + A(K_1^\pm(1400)K^\mp) \sin \theta = A(b^\pm \pi^\mp). \quad (6)$$

Determination of the s -to- d wave ratio of the amplitudes and further study of $B(J/\psi \rightarrow K_1^\pm(1400)K^\mp)$ will eventually resolve the composition of amplitudes and test the large phase hypothesis in 1^+0^- . As it was pointed out previously [16], it is also important to resolve the discrepancy between $B(b^\pm \pi^\mp)$ and $2B(b^0 \pi^0)$ [15], which theory predicts to be equal.

To summarize the experimental situation of the two-body J/ψ decay amplitudes, the existing data strongly favor large relative phases close to 90° between the gluon and the photon decay amplitudes for 1^-0^- , 0^-0^- , 1^-1^- , and $N\bar{N}$, and are consistent with a large phase for 1^+0^- .

What does theory say about these relative phases? In the perturbative picture, the gluonic decay of J/ψ proceeds as depicted in Fig. 1(a). The inclusive decay rate is computed with the gluons placed on mass shell. In contrast, the photon, being far off shell, no corresponding on-shell intermediate state appears in the perturbative diagrams of the photon amplitude [Fig. 1(b)]. Although perturbative QCD (p QCD) is a good description of inclusive charmonium decays, it is questionable whether it works for two-body decay channels of charmonia. To be specific, the p QCD prediction of the asymptotic pion form factor [18],

$$F_\pi(q^2) \approx 16\pi\alpha_s(q^2)f_\pi^2/q^2, \quad (7)$$

has not been reached at the J/ψ mass [19]. Furthermore, the helicity suppression argument of p QCD fails for the $\omega\pi$ decay channel. Though it is tempting, therefore, we cannot argue for the large relative phases on the basis of perturbative diagrams. Whether or not these large relative phases are universal to all two-body decay modes of J/ψ must be determined by experiment. Despite lack of a good theoretical

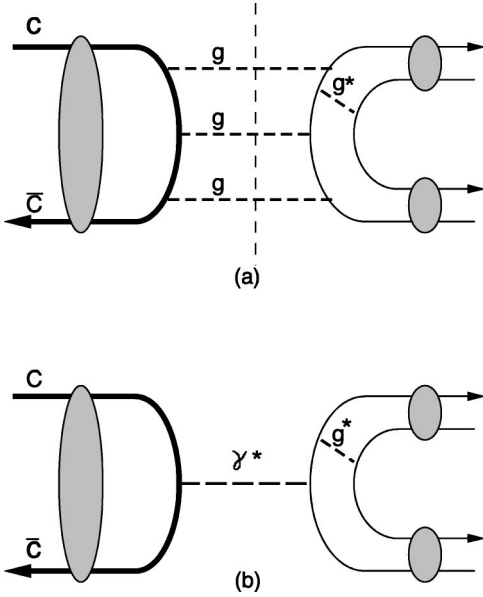


FIG. 1. Decays of charmonium into two mesons (a) through ggg and (b) through one photon. The vertical dashed line in diagram (a) indicates that the gluons are placed on the mass shell when the inclusive decay rate is computed with ggg . If perturbative QCD dominated, the one-gluon-exchange diagram depicted in (b) would dominate in the final state of the one-photon process.

argument at present, we suspect nonetheless that the universal large phases so far found are not an accident.

If the relative phases are close to 90° , it is more likely that the photon decay amplitudes are real and consequently the gluon decay amplitudes are imaginary. The reason is as follows: In order for the photon decay amplitude to have a substantial phase, the final $q\bar{q}$ created by the virtual γ should have a large absorptive part of a long-distance origin. This can happen if there should be a relatively sharp $q\bar{q}$ resonance just around J/ψ . More likely is that many resonances exist below J/ψ as in the vector-meson-dominance scenario or that as in the dual resonance model many increasingly broader resonances appear all the way to high energies [20]. In the former case, only the tails of low-lying resonances contribute to the real part. In the latter case, a few nearby broad resonances can contribute to the imaginary part, but

they are outnumbered by many more resonances below and above J/ψ that contribute to the real part. In comparison, we have less insight in hadronization dynamics of the gluon decay.

Motivated by the results in the amplitude analyses of the two-body of J/ψ decays, we make two postulates.

(1) The relative phases between the gluon and the photon decay amplitudes are universally large for all two-body decays of J/ψ . The photon decay amplitudes are predominantly real and consequently the gluon decay amplitudes are imaginary.

(2) The same pattern holds for $\psi(2S)$ decay as well.

These are the starting assumptions of our analysis that follows.

III. $\psi(2S)$ DECAY

A. Relative phase from experiment

The only large energy scale involved in the three-gluon decay of charmonia is the charm quark mass m_c . Whether one accepts the argument of the universal large phase in exclusive channels or not, therefore, one would naively expect that the corresponding phases should not be much different between the J/ψ decay and the $\psi(2S)$ decay. However, experimental data so far available show that the phases are small at least in some two-body decay modes of $\psi(2S)$. The strongest evidence is in the decay $\psi(2S) \rightarrow 1^- 0^-$, which includes the puzzling $\psi(2S) \rightarrow \rho\pi$. In the case that the final 0^- meson is an octet, we can parametrize the $\psi(2S) \rightarrow 1^- 0^-$ decay amplitudes with the SU(3) singlet amplitude a_1 , the SU(3) breaking correction ϵ due to $m_s - m_{ud}$, and the photon amplitude a_γ . The corresponding amplitudes b_1 and b_γ are introduced for the 0^- singlet. For ϕ - ω mixing, we assume the nonet scheme. The parametrization of the amplitudes [21] is listed in Table I for the decay modes so far studied in experiment [2]. Since the analyzed channels are limited and the uncertainties in their branching fractions are still large, we are unable to perform a meaningful χ^2 fit at present. Therefore, we present only fits to the central values by referring to Table I.

First of all, if we ignored the photon amplitude a_γ , we would obtain $B(K^{*0}\bar{K}^0 + \text{c.c.}) = B(K^{*+}K^-)$, which contradicts experiment, $(0.81 \pm 0.24 \pm 0.16) \times 10^{-4}$ vs < 0.30

TABLE I. Parametrization of the $\psi(2S) \rightarrow 1^- 0^-$ amplitudes and values of branching fractions. The amplitude ϵ represents the T_{33} breaking of $m_s - m_{ud}$ instead of λ_8 breaking. The η - η' mixing angle is chosen to be $\theta_P = -20^\circ$. The fits to the central values are shown for the minimum and the maximum relative phases, $\delta = -\arg(a_1^* a_\gamma) = 0^\circ$ and $\pm 90^\circ$. The ranges of values for $\rho\pi$ and $\omega\eta$ are given for $|\epsilon/a_1| < 1/3$.

Modes	Amplitudes	Branchings (in 10^{-4})	Fits	
			$\delta=0$	$\delta=\pm 90^\circ$
$\rho^+ \pi^- (= \rho^0 \pi^0)$	$a_1 + a_\gamma$	< 0.09	0–0.04	0.21–0.70
$K^{*+} K^-$	$a_1 + \epsilon + a_\gamma$	< 0.15	0.00	0.30
$K^{*0} \bar{K}^0$	$a_1 + \epsilon - 2a_\gamma$	$0.41 \pm 0.12 \pm 0.08$	0.41	0.41
$\omega \pi^0$	$3a_\gamma$	$0.38 \pm 0.17 \pm 0.11$	0.38	0.38
$\omega \eta$	$\sqrt{1/3}(a_1 + a_\gamma)(\omega \eta_8)$	< 0.33	0.06–0.22	0.05–0.31
$\omega \eta'$	$\sqrt{2/3}(b_1 + b_\gamma)(\omega \eta_1)$	$0.76 \pm 0.44 \pm 0.18$	0.76	0.76

$\times 10^{-4}$. The large splitting between these branching fractions requires that a_γ be comparable to a_1 . If we set a_1 and ϵ to zero, we would obtain up to phase space corrections

$$B_0(\omega\pi)/B_0(K^{*0}\bar{K}^0 + \text{c.c.}) = 9/8 \quad (8)$$

in contradiction with the measurement, $(0.38 \pm 0.17 \pm 0.11)/(0.81 \pm 0.24 \pm 0.16)$.⁴ In order to come closer to this ratio of the measured values, a large constructive interference should occur in $K^{*0}\bar{K}^0$; that is, the relative phase must be small between $a_1 + \epsilon$ and $-2a_\gamma$. Then, assuming that a_1 and ϵ have a common phase, we have a large destructive interference between a_1 and a_γ for both $\rho\pi$ and $K^{*\pm}K^\mp$ in agreement with experiment.⁵ This solves the $\rho\pi$ puzzle and explains also the missing of the $K^{*\pm}K^\mp$ mode in experiment.

With these qualitative observations in mind, we have fitted to the central values of the observed branching fractions and then have computed with those parameter values the branching fractions of the modes for which only the upper bounds have been determined. In Table I we have listed the fit with $\delta \equiv -\arg(a_1^* a_\gamma) = 0^\circ$ and the large phase fit with $\delta = \pm 90^\circ$ for comparison. When the central values of $B(\omega\pi)$ and $B(K^{*0}\bar{K}^0 + \text{c.c.})$ are fitted with $\delta = 0^\circ$, the ratio of the photon and the gluon amplitude turns out to be

$$a_\gamma/(a_1 + \epsilon) \simeq -0.76. \quad (9)$$

For comparison, $a_\gamma/(a_1 + \epsilon) \simeq 0.14$ in the case of J/ψ . We would expect $|a_\gamma/a_1| \approx 0.22$ if the ratio of $\Gamma(\psi(2S) \rightarrow \gamma^* \rightarrow 1^- 0^-)$ to $\Gamma(\psi(2S) \rightarrow ggg \rightarrow 1^- 0^-)$ is roughly equal to $\Gamma(\psi(2S) \rightarrow \gamma^* \rightarrow X)/\Gamma(\psi(2S) \rightarrow ggg \rightarrow X)$. Since experiment shows that $\Gamma(\omega\pi)/\Gamma(1^+ 1^-)$ ($\propto |a_\gamma|^2$) is about the same for J/ψ and $\psi(2S)$, the large number in Eq. (9) results from a strong suppression of the total gluonic amplitude $a_1 + \epsilon$ in $\psi(2S)$. As the value of ϵ is varied in the range of $|\epsilon/a_1| < 1/3$, the value of $B(\rho^\pm \pi^\mp)$ varies between 0 and 0.04×10^{-4} . The values for $B(\rho^\pm \pi^\mp)$ and $B(K^{*\pm} K^\mp)$ can be increased if we stretch $B(\omega\pi)$ and $B(K^{*0}\bar{K}^0 + \text{c.c.})$ within the experimental uncertainties. In contrast, the fit with $\delta = \pm 90^\circ$ overshoots the upper bound on $B(K^{*\pm} K^\mp)$ and, if $|\epsilon| < \frac{1}{3}|a_1|$, the upper bound on $B(\rho^\pm \pi^\mp)$ very badly. A fit with $\delta = \pm 90^\circ$ is virtually impossible even with experimental uncertainties unless $|\epsilon| \gg |a_1|$. We thus conclude that the relative phase between a_1 and $-a_\gamma$ should be small in $\psi(2S) \rightarrow 1^- 0^-$ contrary to the J/ψ decay.

Though it is less conclusive, a small phase seems to be favored in the $1^+ 0^-$ decay of $\psi(2S)$ too. It is conspicuous in experiment [3] that the $K_1^\pm(1400)K^\mp$ mode is strongly suppressed relative to the $K_1^\pm(1270)K^\mp$:

⁴It is possible that the SU(3) breaking in the strange and non-strange meson wave functions $(f_{\pi f_\rho}/f_{K f_{K^*}})^2$ may be responsible for part of the discrepancy.

⁵This main feature of the fit to the $\psi(2S) \rightarrow 1^- 0^-$ amplitudes is found in the earlier paper by Chen and Braaten [12] and, in particular, in the paper by Tuan [13].

$$B(\psi(2S) \rightarrow K_1^\pm(1270)K^\mp) = (10.0 \pm 1.8 \pm 2.1) \times 10^{-4},$$

$$B(\psi(2S) \rightarrow K_1^\pm(1400)K^\mp) < 3.1 \times 10^{-4}. \quad (10)$$

$B(\psi(2S) \rightarrow b^\pm \pi^\mp)$ is halfway between them [3]:

$$B(\psi(2S) \rightarrow b^\pm \pi^\mp) = (5.2 \pm 0.8 \pm 1.0) \times 10^{-4}. \quad (11)$$

We can use Eq. (4) as the parametrization of $\psi(2S) \rightarrow 1^+ 0^-$. First of all, if a_1 dominated over $a_{\gamma A/B}$, we would have $B(b^\pm \pi^\mp) \simeq 2B(K_1^\pm(1270)K^\mp) \simeq 2B(K_1^\pm(1400)K^\mp)$ for $\theta \simeq 45^\circ$ in disagreement with experiment. Just as in $\psi(2S) \rightarrow 1^- 0^-$, $|a_1|$ is comparable to $|a_{\gamma A/B}|$. Next, the strong suppression of $K_1^\pm(1400)K^\mp$ relative to $K_1^\pm(1270)K^\mp$ can be realized only when $a_1 + \sqrt{1/5}a_{\gamma B}$ interferes destructively with $a_{\gamma A}$. Therefore, the relative phase between a_1 and $a_{\gamma A/B}$ must be small modulo π . To obtain $B(K_1^\pm(1270)K^\mp) \simeq 2B(b^\pm \pi^\mp)$, we need $a_1 + \sqrt{1/5}a_{\gamma B} \simeq -a_{\gamma A}$. The allowed range of the amplitude ratios was plotted in Ref. [16] by choosing all amplitudes as relatively real and assuming tentatively the s -wave decay for phase-space corrections. Though it is not impossible to fit the three branching fractions with $\theta \simeq 90^\circ$, we must have $a_1 \simeq 0$ and $\sqrt{1/5}a_{\gamma B} \simeq -a_{\gamma A}$ in that case.

To summarize for $\psi(2S)$, the data on $1^- 0^-$ virtually exclude the possibility of a small phase between a_1 and a_γ . A fit to $1^+ 0^-$ has more room when the relative phase is small. There is no evidence for that the relative phase must be large in $\psi(2S)$. For some reason, a large relative phase does not seem to occur in the $\psi(2S)$ decay. We ask what causes this marked difference between J/ψ and $\psi(2S)$ when we postulate the universal large phase for $\psi(2S)$ as well as for J/ψ .

B. Excess hadronic rate in inclusive hadronic decay

It has been noticed that when one computes the inclusive hadronic decay rate of $\psi(2S)$ through ggg by subtracting the rates of the cascade and electromagnetic decays from the total decay rate, it is substantially larger than what we expect from an extrapolation of J/ψ . The number with a conservative error estimate based on the listings of *Reviews of Particle Physics* [15] is

$$\frac{B(\psi(2S) \rightarrow ggg + gg\gamma)}{B(J/\psi \rightarrow ggg + gg\gamma)} = 0.23 \pm 0.07, \quad (12)$$

which should be compared with

$$\left(\frac{\alpha_s(\psi(2S))}{\alpha_s(J/\psi)} \right)^3 \frac{B(\psi(2S) \rightarrow 1^+ 1^-)}{B(J/\psi \rightarrow 1^+ 1^-)} = 0.134 \pm 0.034. \quad (13)$$

Smaller errors (0.226 ± 0.052 vs 0.141 ± 0.012) have been attached in a recent paper [22] with a different error estimate. We would expect that the two numbers should be equal to each other since the wave functions at origin appear in common in Eqs. (12) and (13). The discrepancy of 60–70%

TABLE II. Branching fractions of J/ψ and $\psi(2S)$, and the ratio $B(\psi(2S)\rightarrow h)/B(J/\psi\rightarrow h)$. The $\pi^+\pi^-\pi^0$ mode is not included here since $J/\psi\rightarrow h$ is entirely $J/\psi\rightarrow\rho\pi$ within experimental uncertainties.

Modes	Branchings of J/ψ	Branchings of $\psi(2S)$	Ratio
$2(\pi^+\pi^-\pi^0)$	$3.37\pm 0.26\times 10^{-2}$	$3.0\pm 0.8\times 10^{-3}$	$8.9\pm 2.5\%$
$3(\pi^+\pi^-\pi^0)$	$2.9\pm 0.6\times 10^{-2}$	$3.5\pm 1.6\times 10^{-3}$	$12.1\pm 6.1\%$
K^+K^-	$2.37\pm 0.31\times 10^{-4}$	$1.0\pm 0.7\times 10^{-4}$	$42\pm 30\%$
$\pi^+\pi^-K^+K^-$	$7.2\pm 2.3\times 10^{-3}$	$1.6\pm 0.4\times 10^{-3}$	$22.2\pm 8.8\%$
$p\bar{p}$	$2.12\pm 0.10\times 10^{-3}$	$1.9\pm 0.5\times 10^{-4}$	$8.9\pm 2.4\%$
$p\bar{p}\pi^0$	$1.09\pm 0.09\times 10^{-3}$	$1.4\pm 0.5\times 10^{-4}$	$12.8\pm 4.7\%$
$p\bar{p}\pi^+\pi^-$	$6.0\pm 0.5\times 10^{-3}$	$8\pm 2\times 10^{-4}$	$13.3\pm 2.1\%$

between them alarms us particularly because all numbers involved have been repeatedly measured over many years.⁶

In comparison we find no similar excess in $Y(2S)$ though experimental uncertainties are large. In terms of the ratio of branching ratios, $\bar{B}(ggg+gg\gamma)\equiv B(Y\rightarrow ggg+gg\gamma)/\alpha_s(Y)^3B(Y\rightarrow\mu^+\mu^-)$, three Y 's are more in line:

$$\bar{B}(ggg+gg\gamma)=\begin{cases} (4.5\pm 0.2)\times 10^3, & Y(1S), \\ (4.9\pm 0.9)\times 10^3, & Y(2S), \\ (4.0\pm 0.4)\times 10^3, & Y(3S), \end{cases} \quad (14)$$

where the total leptonic branching⁷ for $Y(3S)$ has been substituted with three times $B(\mu^+\mu^-)$, the only quoted leptonic branching. It appears that the excess in $\bar{B}(ggg+gg\gamma)$ is unique to $\psi(2S)$. However, this excess in the inclusive rate has not shown up in the rates of the exclusive channels so far measured. In fact, the ratio $B(\psi(2S)\rightarrow h)/B(J/\psi\rightarrow h)$ scatters around the expected value [≈ 13 – 14% of Eq. (13)], which was often called the *14% rule*. Some remarks should be in order on it.

First of all, the 14% rule is largely violated in many of two-body and quasi-two-body channels, as we recently learned in the BES data [2]. The $\rho\pi$ channel is an extreme case. For multihadron channels, there are actually not so many modes that are available for testing the 14% rule.⁸ In Table II, we have tabulated the ratios for the modes not listed in [2] but available for comparison. We see that the ratios scatter rather widely above and below 14% with some tendency of being smaller than 14%, but with fairly large ex-

perimental uncertainties. It is important to notice that the branching fractions of all modes in Table II add up to no more than 15% of the total gluonic decay branching of J/ψ . Indeed, only one charge state has been available for comparison from each of 5π , 7π , $2\pi K\bar{K}$, and $N\bar{N}n\pi$. We have not yet seen a comparison of the rest. The so-called 14% rule is based on a very limited number of decay modes. It is premature to preclude the hadronic excess with the data of multihadron exclusive channels.

If future experiments show that the hadronic excess in $\psi(2S)$ is real, it may have something to do with the $\rho\pi$ puzzle [22] and with the abrupt change of the relative phase of amplitudes from J/ψ to $\psi(2S)$. A process responsible for the excess inclusive hadron rate can interfere with the short-distance gluon decay in exclusive modes. If its amplitude makes a large destructive interference with the three-gluon amplitude in $\psi(2S)\rightarrow 1^-0^-$ and if the sum is nearly real and comparable to the photon amplitude in magnitude, our puzzle can be solved. We shall look into possible sources of this rate excess in the following.

IV. ADDITIONAL HADRONIC AMPLITUDE IN $\psi(2S)$

A. General consequences

If the origin of the problem is in the interference of an unknown additional process with the short-distance gluon decay of $\psi(2S)$, we expect a general pattern of correlation between the decay angular distribution and suppression or enhancement.

The decay angular distribution for two final hadrons is generally of the form

$$d\Gamma/d\Omega \propto 1 + a \cos^2 \theta, \quad (|a| \leq 1), \quad (15)$$

where θ is the polar angle measured from the e^+e^- beam direction. For 1^-0^- and 0^-0^- decays, the value of a is constrained kinematically to $+1$ and -1 , respectively, while it is determined dynamically by the helicity content, ± 1 or 0 , of the final state in other decays. In 1^-0^- and 0^-0^- , therefore, any additional amplitude has the same angular dependence as the three-gluon and the photon amplitude irrespective of its origin. Consequently a high degree of destructive or constructive interference with an additional amplitude is possible in these decays. Observation of the strongest suppression in the 1^-0^- mode is consistent with this pattern [2].

⁶The author learned that the BES Collaboration is considering a different determination of the cascade decay branchings [23].

⁷The branching to $Y(2S)\gamma\gamma$ quoted in [15] is a sum of the cascade decay branchings $\chi_{bJ}(2P)\gamma\rightarrow Y(2S)\gamma\gamma$ in view of the $\gamma\gamma$ invariant mass spectrum [24] and also of its magnitude [$>B(Y(2S)\pi^0\pi^0)$]. It is counted in the radiative decay branchings separately listed in [15].

⁸Gu and Li [22] lumped multihadron modes together and compared between J/ψ and $\psi(2S)$. Then the number is dominated by three modes $\pi^+\pi^-\pi^0$, $2(\pi^+\pi^-\pi^0)$, and $3(\pi^+\pi^-\pi^0)$, which happen to be below 14%. The mode $\pi^+\pi^-\pi^0$ is actually $\rho\pi$ and its nonresonant content is consistent with zero [15]. More revealing are the ratios of the individual modes.

We expect that the decay rates of $\psi(2S) \rightarrow 0^- 0^-$ may also be quite different from those of $J/\psi \rightarrow 0^- 0^-$. In terms of $\bar{B}(0^- 0^-) \equiv B(0^- 0^-) / \alpha_s^3 B(\mu^+ \mu^-)$, the current data [15] give

$$\bar{B}(\pi^+ \pi^-) = \begin{cases} 0.15 \pm 0.02 & \text{for } J/\psi, \\ 0.8 \pm 0.5 & \text{for } \psi(2S), \end{cases}$$

$$\bar{B}(K^+ K^-) = \begin{cases} 0.24 \pm 0.03 & \text{for } J/\psi, \\ 0.94 \pm 0.66 & \text{for } \psi(2S). \end{cases} \quad (16)$$

Within the large experimental uncertainties we see a hint of large constructive interference in $\psi(2S) \rightarrow 0^- 0^-$. In contrast, in other processes an additional amplitude and the three-gluon amplitude have different angular distributions in general. A large interference can occur only when the dynamical mechanisms of the two processes are similar. Otherwise it should be a result of a high degree of accident. When a large disparity is observed between the corresponding two-meson decay rates of J/ψ and $\psi(2S)$, therefore, the decay angular distribution of this channel will also be very different between J/ψ and $\psi(2S)$. This will give a good test of the idea of interference with an additional amplitude.

The other consequence is in multibody final states. Since the additional process enhances the inclusive rate, a large number of exclusive decay channels should receive enhancement rather than suppression. When there are many hadrons in the final state, the chance of interference between the amplitudes of different decay mechanisms is much smaller because of the difference in the subenergy dependence and event topology. Therefore the enhancement will not be dramatic. While we have not yet seen such an enhancement in Table II, we expect that the branching fraction tends to be enhanced in many nonresonant multibody channels of $\psi(2S)$ relative to J/ψ .

Where does the additional amplitude possibly come from? There are few options left in modifying charmonium physics radically. Since the aspects of perturbative QCD have been well understood, we are bound to look for the origin of the problem in long-distance physics of one kind or another.

B. $\psi(2S) \rightarrow D\bar{D} \rightarrow$ hadrons

One unique feature of $\psi(2S)$ is the close proximity of its mass to the $D\bar{D}$ threshold. The $\psi(2S)$ mass is only 43 MeV (53 MeV) below $D^0\bar{D}^0$ (D^+D^-), while $Y(3S)$ is 200 MeV away from the $B\bar{B}$ threshold. Can the small energy difference⁹ between $\psi(2S)$ and $D\bar{D}$ have anything to do with the excess? It may happen that $\psi(2S)$ picks up a light quark pair through soft gluons and dissociates virtually into $D\bar{D}$, which in turn annihilate into light hadrons. (See Fig. 2.) The dominant process of the $D\bar{D}$ annihilation is through $c\bar{c}$ annihilation through a single hard gluon. The small energy denominator enhances the creation of virtual $D\bar{D}$ while

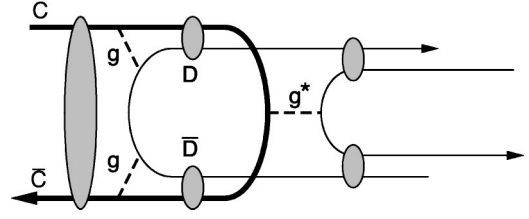
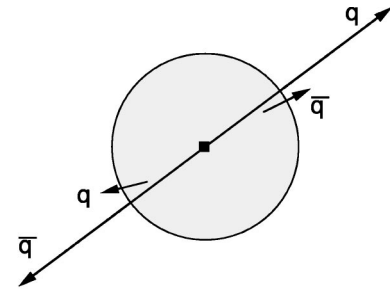


FIG. 2. The decay $\psi(2S) \rightarrow D\bar{D}$ (off shell) \rightarrow meson + meson.

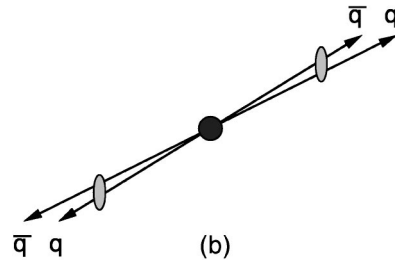
p -wave creation compensates the enhancement. It is difficult, actually nearly impossible, to give a reliable computation of this sequence. Deferring an estimate of the rate to the future, we shall comment here only on whether the $D\bar{D}$ contribution can have a final-state interaction phase large enough to cancel the perturbative gluon amplitude or not.

Since the $D\bar{D}$ intermediate state is above the $\psi(2S)$ mass, a phase of amplitude must come from the subsequent annihilation of $c\bar{c}$ and thereafter. After an energetic light quark pair $q\bar{q}$ is created from $c\bar{c}$, each q and \bar{q} picks up a soft light quark from the light quark cloud of $D\bar{D}$ to form mesons. [See Fig. 3(a)]. Kinematically, this step of the hadron formation process is quite different from that of the timelike electromagnetic form factor of a meson in which energetic light quarks pick up collinear quarks created by a hard gluon. [See Fig. 3(b)]. In our case the color-dipole moment is large for all pairs of quarks [26]. Furthermore, the c.m. energy of a hard quark in one meson and a soft quark in the other meson is in the low-energy resonance region:

$$\sqrt{s} = O(\sqrt{2\Lambda_{QCD}m_c}) < 1 \text{ GeV}. \quad (17)$$



(a)



(b)

FIG. 3. (a) Formation of a light meson pair by energetic $q\bar{q}$ from $c\bar{c}$ and wee $q\bar{q}$ in $D\bar{D}$ annihilation. The invariant mass is small for qq and for $\bar{q}\bar{q}$. The arrows denote directions and magnitudes of momenta. (b) Light meson pair formation in one-photon annihilation where all quarks are hard.

⁹This was brought to the author's attention by Rosner [25].

Therefore, one cannot argue that final-state interactions should be small between final mesons. We expect that there is a good chance for the amplitude of $\psi(2S) \rightarrow D\bar{D} \rightarrow$ mesons to acquire a substantial final-state interaction phase. The weakness in this argument is that the phase can be large but need not be large.

The idea of virtual $D\bar{D}$ dissociation actually has some common features with that of the ‘‘higher Fock component’’ of charmonia [27,28]. The $D\bar{D}$ state can be viewed as part of the four-quark Fock space of $\psi(2S)$. The higher Fock component was proposed as an additional contribution to $J/\psi \rightarrow 1^-0^-$ to solve the $\rho\pi$ puzzle [28]. It was argued that it is more significant in J/ψ than in $\psi(2S)$. As we have emphasized, however, there is nothing anomalous about $J/\psi \rightarrow 1^-0^-$.

C. $\psi(2S) \rightarrow$ resonance \rightarrow hadrons

The second idea is a twist of an old one: A noncharm resonance may exist near the $\psi(2S)$ mass and give an extra contribution to the hadronic decay rate. A glueball was proposed earlier at the J/ψ mass to boost the $\rho\pi$ decay rate of J/ψ [7]. However, we now want it near $\psi(2S)$, not near J/ψ . We look into the possibility that some resonance around the $\psi(2S)$ mass, a glueball, or a four-quark resonance, destructively interferes with the perturbative $\psi(2S) \rightarrow ggg \rightarrow 1^-0^-$ decay. Admittedly, the idea is *ad hoc* and there is some difficulty aside from unnaturalness.

Light quark resonances of high mass (~ 3.7 GeV) and low spin are normally too broad to be even recognized as resonances. Four-quark resonances may be an alternative if they exist at all. The mass of 3.7 GeV is normally considered as too high for the lowest vector glueball. An excited glueball state of $J^{PC} = 1^{--}$ serves our purpose. Whatever its origin is, let us introduce here such a resonance, call it R , and see its consequences.

In order for $\psi(2S)$ to decay through R as strongly as through three gluons, the coupling f of R to $\psi(2S)$ defined by $-fm_R\psi_\mu R^\mu$ must be large enough. The $\psi(2S)$ - R mixing at the $\psi(2S)$ mass is given by

$$\varepsilon \approx \frac{f}{\Delta m - i\Gamma_R}, \quad (18)$$

where $\Delta m = m_R - m(\psi(2S))$ and Γ_R is the total width of R . It leads to $\Gamma(\psi(2S) \rightarrow R \rightarrow \text{hadrons}) \approx |f|^2/\Gamma_R$ when $|\Delta m| < O(\Gamma_R)$. To obtain $\Gamma(\psi(2S) \rightarrow R \rightarrow \text{hadrons}) \approx \Gamma(\psi(2S) \rightarrow ggg)$, we need therefore

$$|f|^2 \approx \Gamma_R \Gamma(\psi(2S) \rightarrow ggg). \quad (19)$$

If R is a light quark resonance $q\bar{q}$, $|f|$ would be much too small for the following reason: While $|f|^2$ is of the order of $\Gamma(\psi(2S) \rightarrow ggg)\Gamma(R \rightarrow ggg)$ for $q\bar{q}$, we expect $\Gamma(R \rightarrow ggg) \ll \Gamma_R$ because of the α_s^3 suppression of $q\bar{q} \rightarrow ggg$. Therefore there is no chance to satisfy Eq. (19). It is likely that the same argument applies to four-quark resonances.

For glueballs, we simply do not have enough quantitative understanding to rule out a large enough coupling to $\psi(2S)$. Hou and Soni [7] proposed a glueball near the J/ψ mass in order to enhance $J/\psi \rightarrow \rho\pi$ (and its symmetry-related modes) but not other decay modes. To accomplish it, this glueball must have very special, if not unnatural, properties [29]: It is nearly degenerate with J/ψ with a quite narrow width for an object of mass ~ 3 GeV and decays predominantly into 1^-0^- . Later Hou [30] relaxed the constraint on the $\rho\pi$ branching to argue that such a glueball was not yet ruled out by the search of the BES Collaboration [31].

In our case, since a glueball is introduced to account for the hadronic excess, it should couple not primarily to the 1^-0^- channels, but to many other channels. What we need is a generic vector glueball with no special or unusual properties. If Γ_R is as narrow as 100 MeV, for instance, the mixing $|\varepsilon| = O(10^{-2})$ would be able to account for the excess in the inclusive hadron decay of $\psi(2S)$. The width can be wider. In that case the pole transition strength f should be stronger according to Eq. (19). Since the glueball R couples to a photon only indirectly through its mixing to a quark pair, it is hard to detect R in the hadronic cross section of e^+e^- annihilation near the $\psi(2S)$ mass. Searching by hadronic reactions such as $p\bar{p}$ annihilation is a daunting task. From a purely experimental viewpoint, such a resonance has not been ruled out [30].

Though the resonance scenario is admittedly a longshot, it is one of very few options left to us. One reason to pursue this somewhat unnatural scenario is that the amplitude for $\psi(2S) \rightarrow R \rightarrow X$ has automatically a large phase when $\Delta m < O(\Gamma_R)$ since the coupling f is likely real, that is, dominated by the dispersive part. If that is the case, the resonant amplitude can interfere strongly with the three-gluon amplitude in two-meson decays.

V. CONCLUSION

We have searched for a clue to solve the $\rho\pi$ puzzle in this paper. Our purpose is to locate the source of the problem rather than to offer a final solution of the problem. Two threads have been exposed which may eventually lead us to a solution of the $\rho\pi$ puzzle. They are the phases of the decay amplitudes and a possible excess in the inclusive hadronic decay rate of $\psi(2S)$. An experimental confirmation of the excess will be the most useful in directing theorists. If it is confirmed, it will be quite an important experimental discovery by itself. One crucial piece of experimental information will be the angular distributions of J/ψ and $\psi(2S)$ into channels other than 1^-0^- and 0^-0^- . The difference in the angular distributions should have a direct correlation with the enhancement and suppression in general. As for the source of an additional process, the virtual $D\bar{D}$ pair and the vector glueball are two options that cannot be ruled out. To be frank, however, we admit that both ideas have unnaturalness. More of an attractive alternative is highly desired. It is possible that the large relative phases so far observed in J/ψ decay are an accident and that the $\rho\pi$ puzzle is a problem of incalculable long-distance complications. However, our hope is that there might be something novel, simple, or fundamental hidden beneath the issue.

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