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REVIEW ARTICLE

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Charm physics — A field full of challenges and opportunities

Abstract In this review, we discuss some interesting issues in charm physics, which is full of puzzles and challenges. So far in this field, there exist many problems which have not obtained satisfactory answers yet as more unexpected phenomena Continue to be observed at the current facilities of high energy physics. Charm physics may become an ideal place for searching new resonances and studying non-perturbative QCD effects, moreover it is probably an area for exploring new physics beyond the Standard Model. More data will be available at the BESIII, B-factories, LHC and even the future ILC, which may open a wide window to a better understanding of its nature.

Keywords charm, QCD, exotic state, final state interaction, new physics beyond standard model

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Contents

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1 Introduction

The charm quark was a long-expected member of the quark family. It is noted that if only the u quark is the intermediate fermion at the s-channel, the cross section of the scattering $s + W^+ \to d + W^+$ would increase with the incoming energy, which is unacceptable in physics. It demands the existence of another species of quarks having the same charge as the u quark which serves as an additional intermediate fermion to compensate for the bad high energy behavior of the process, i.e. retaining the unitarity; this new species are the charm quark [1]. Moreover, without the charm quark the anomaly in the electro-weak model cannot be cancelled [2, 3], thus the renormalizability of the whole theory would be spoiled. Later, by studying the K^0-$\bar{K}^0$ mixing, Gaillard and Lee [4] estimated that the mass of the charm quark should be around 1.5 GeV. Thus all the urgency of saving the beautiful theory appealed to discovering this charming "charm" quark. Then the discovery of the J/ψ meson and other members of the ψ family became a milestone in particle physics [5, 6]. The ground states of the family are J/ψ and η_c whose masses are about 3.1 GeV and 2.98 GeV respectively, so roughly it implies that $m_c \sim 1.5$ GeV, which is amazingly consistent with Gaillard and Lee's estimate. However, it is definitely not the end of story for the quark family. The discovery of the bottom quark requires the existence of a partner, which evaded observation for a very long while until it was eventually found at the TEVATRON [7, 8], and it possesses an astonishingly heavy mass of about 176 GeV. Till then, the three generation structure of quarks and leptons seems complete, even though the fourth generation of quarks and leptons is still under discussions.

We know, the u, d and s quarks reside in a triplet of the global $SU(3)$ quark model, which successfully describes relevant phenomenology, but there is, so far, not a special symmetry to associate the rest of the three heavy quarks†. The charm quark may be a special one in the quark family, because it is heavier than the first three light quarks and does not belong to the regular flavor $SU(3)$, but stands in a weak doublet with the light strange quark. The charm is not light at all, but is also not as heavy as the bottom nor even the top quarks. The intermediate mass determines the special characteristics of hadrons which contain charm and/or anti-charm properties. Recently, some researches have noted that the first four quarks are different from the last two heavy ones: the bottom and top quarks, such as the top assisted technicolor model, but this is beyond the scope of our review.

Since the charm quark is indeed sufficiently heavy (as will be discussed below, "heavy" means it is heavier than the binding energy scale Λ_{QCD}), it is natural to use the potential model to evaluate the spectra of J/ψ and other family members: η_c, χ_c and even h_c etc. and their excited states, as well as their fine-structures and decay modes. However, since it is not too heavy, the relativistic correction to the potential which generally consists of a Coulomb piece and a confinement one (for example, the linear potential) [9−11], is more serious than that for the Υ family. Still, by adjusting parameters, great success has been achieved for the heavy charm-quarkonia (charmonia) and the results are satisfactorily consistent with data for ground and lower excited states. Very recently, Voloshin gave an enlightening review on charmonium where the relevant topics were discussed in details [12]. However, the story is far from its end yet. There have been many puzzles in the field, especially as more accurate measurements are being done, they do not disappear. Moreover, several new resonances have been observed and they do not seem easily described by the simplest valence quark structure, i.e. the meson is composed of quark-antiquark and the baryon is composed of three quarks.

Now, let us present a rough list about the puzzles. The first one may be the famous $\rho\pi$ puzzle where the branching ratio of $\psi' \to \rho\pi$ is too small compared with Refs. [13−15]. Then the decay mode $J/\psi \to \rho\pi$ is forbidden by the hadronic helicity conservation, thus its rate must be sufficiently small, but by contrast, it is one of the main decay channels of J/ψ. The sizable D^0-$\bar{D}^0$ mixing which has been recently observed [16, 17], indicates that there must exist a mechanism beyond the Standard Model (SM). There are many newly observed resonances, which may demand interpretation.

On the theoretical aspect, great efforts have been done to look for reasonable explanations. The first step is to make sure whether the puzzles in the new observations are due to the flaws of our theoretical framework or there is new physics beyond the SM. Indeed, to understand the experimental measurements, one must calculate the corresponding quantities, where evaluation of the hadronic matrix elements is the key and the obstacle. Since the hadronic matrix elements are fully governed by the non-perturbative QCD effects for which, at present, there is no reliable way to deal with, all the results we have achieved must possess certain uncertainties. Properly estimating the errors and reliability of the results thus becomes an important issue in theoretical calculations.

Below, we will discuss all the topics in separate sections. We first recall the mechanisms which govern the weak transitions of D mesons, and then focus on discussion of the rare decays, because the regular Cabibbo-favored channels have been thoroughly investigated in both theory and experiment. Then since all the topics are related to theoretical evaluation of the hadronic matrix elements, in the next section we review the current status. In Section 4, we are concerned with the final state interaction, especially focusing on the hadronic triangle calculation and briefly discussing other schemes. In Section 5, we will discuss the hadronic helicity selection rule and its violation. In Section 6 we are concerned with the $\rho\pi$ puzzle, and in Section 7, we show how the QCD multi-expansion theory works very well for the pion radiation from excited states of Υ, but has difficulties for $\psi(nS) \to \psi(mS) + \pi\pi$ $(n > m)$. In Section 8, in some details, we discuss the newly observed resonances at the charm energy region, and some of them seem to be exotic and need a reasonable interpretation. In Section

† There indeed is the so-called Heavy Quark symmetry $SU_f(2) \otimes SU_s(2)$ (which we will discuss later) to connect the b and c quarks, but that symmetry is a symmetry when the quark mass approaches infinity and mainly can simplify the calculation of transition from b to c. It is not like the $SU(3)$ for light flavors and is also not a symmetry in the common sense.

9, we discuss the D^0-$\bar{D}^0$ mixing and related theoretical proposals, while also discussing the possible CP violation observable. In Section 10, we consider the charmed baryons and double heavy baryons Ξ_{cc}, where the difference between the lifetimes of Λ_c and $D^{\pm}$, D^0 is especially focussed on, and we will also briefly discuss the charmed pentaquark. In Section 11, we specially discuss the diquark structure in baryons, especially in the heavy baryons, because it is an important issue when studying baryons. The last section is devoted to a brief discussion.

2 The effective Lagrangian and transition amplitudes

Generally, this is a topic which is familiar to most of theorists working in this field. Thus after a short introduction, we will turn our attention to the application of the theoretical framework for investigating rare decays of D mesons and J/ψ.

2.1 The decay constants of D_s and $D^{0,\pm}$

The most important parameters for the weak decays of D and D_s mesons are their decay constants. Generally, the decay constant of a pseudoscalar meson is measured via its leptonic decays $P \to l\bar{\nu}$ ($l = e,\ \mu,\ \tau$ as long as kinematics allows) and the width is written as [18]

$$\Gamma(P \to l\bar{\nu}) = \frac{G_F^2}{8\pi} f_P^2 m_l^2 M_P \left(1 - \frac{m_l^2}{M_P^2}\right) |V_{qq'}|^2$$

where M_P, m_l are the masses of the pseudoscalar meson and the lepton, $V_{qq'}$ is the corresponding CKM entry and q, q' are the valence quarks in the pseudoscalar meson. There, f_P is the decay constant which we are going to obtain. The CLEO collaboration has achieved

$$f_{D^+} = (222.6 \pm 16.7^{+2.8}_{-3.4})\ \text{MeV}$$

in the decay of $D^+ \to \mu^+\nu$ [19, 20].

Rosner has obtained the average of f_{D_s} as [18]

$$f_{D_s} = (274 \pm 10)\ \text{MeV}$$

There is a discrepancy of about three standard deviations between this result and a recent unquenched lattice calculation [23].

Very recently, the CLEO collaboration reported their new result as [22]

$$f_{D^+} = (205.8 \pm 8.5 \pm 2.5)\,\text{MeV}$$

by assuming $|V_{cd}| = |V_{us}|$. They also obtained

$$\frac{B(D^+ \to \mu^+\nu) - B(D^- \to \mu^-\bar{\nu})}{B(D^+ \to \mu^+\nu) + B(D^- \to \mu^-\bar{\nu})} = 0.08 \pm 0.08$$

which means that no CP violation was observed in the leptonic decay. The Belle group reported [21]

$$f_{D_s} = (275 \pm 16(\text{stat}) \pm 12(\text{syst}))\,\text{MeV}$$

Only when the decay constants are accurately measured can the theoretical predictions on the hadronic transitions be trustworthy, so that more precise experiments are necessary. In fact, measurement on decay constants is not easy, not only because the current available database on D, especially D_s is not large enough to guarantee high statistics, but also because reactions such as $P \to l\bar{\nu} + \gamma$ can influence extraction of f_P. The BESIII will collect the largest database of D mesons, so one may expect to get very accurate decay constants for D and D_s.

2.2 The effective Lagrangian of weak interactions

The charmonia J/ψ, η_c and their excited states mainly decay via strong interactions which are OZI-suppressed (and will be discussed in later sections). Instead, D mesons decay via weak interactions or electromagnetic radiation. In this framework, the semileptonic decays are well understood, so that we concentrate ourselves here on the non-leptonic decays. Let us briefly review the general situation of the D decays. The effective Lagrangian for weak interaction is written as

$$L_{\text{eff}}^{|\Delta c|=1} = \frac{G_F}{\sqrt{2}}\left[V_{cs}^* V_{uq}(C_1 O_1 + C_2 O_2) - V_{cb}^* V_{qb} \sum_{i=3}^{10} C_i O_i\right] + h.c.$$

where q can be either d or s, and the first operator $O_1 = (\bar{s}c)_{V-A}(\bar{u}q)_{V-A}$ with $(\bar{q}q')_{V-A} \equiv \bar{q}\gamma_\mu(1-\gamma_5)q'$ originates from the tree level while others are induced by loops, O_3 to O_6 are the strong penguin operators, and the rest ($i = 7$ to 10) are due to the γ, Z-penguins and box diagram [24].

2.3 Rare decays of D mesons

The weak decays of D mesons can be categorized into Cabibbo favored, Cabibbo suppressed and doubly suppressed modes. The first type includes the modes such as $D^+ \to \bar{K}^0 + \pi^+$ etc. and the second one was discussed by Abbott, Sikivie and Wise [25]. For the Cabibbo favored decay modes, one usually only considers the so called spectator mechanism where the light quark behaves as a spectator when the heavy charm quark transits into s quark plus a $u\bar{d}$ pair. In this picture, annihilation and W-exchange between charm and light anti-quarks can be neglected because of the linear momentum matching (namely, for annihilation and W-exchange, a quark-anti-quark pair must be produced from vacuum, or in other words are produced by soft gluons, so that their linear momenta are small, whereas the quark-antiquark pair occurring directly from the effective vertex possess large linear momenta. As they combine, a quark (anti-quark) emerges from vacuum to constitute a hadron where the two constituents must have close linear momenta; the large momentum difference would greatly suppress the probability). However, for the Cabibbo-suppressed or even doubly-suppressed modes, the annihilation and W-exchange mechanisms as well as the penguin contributions become important. In fact, there are several channels where the spectator mechanisms do not contribute at all, thus these modes would be ideal places to study such small effects which may manifest some unknown mechanisms. Moreover, for the CP violation, the contribution from the penguin and even the electro-weak penguin would be crucially significant. We will discuss these issues in the following sections.

Another interesting rare decay mode are those processes where the light quark (anti-quark) transits while the heavy charm behaves as a spectator which only provides a color source. Such a reaction includes $D^* \to D + \gamma$, $D^* \to D\pi$ etc. where the photon and pion can be emitted from either charm or light flavor. We used to study a special case where in a heavy baryon containing two heavy quarks radiates a photon and transits into a lower states with the same flavor [26] in terms of the Bether-Salpeter (B-S) equation. It was indicated that the branching ratios of such decays are very small and are hard to measure at the present luminosity, however, for BESIII and LHCb, the situation may be greatly improved.

Recently, Li and Yang [27] calculated the branching ratios of $D^+ \to D^0 + e^+ + \nu$, $D_s^+ \to D^0 + e^+ + \nu$, $D_s^+ \to D^+ + e^+ + e^-$ in SM, however, their results indicate that only the branching ratio of $D_s^+ \to D^0 + e^+ + \nu$ could reach 10^{-8}. According to the sensitivity of BESIII, B-factories, Super-B and LHCb, it might be observed at Super-B and LHCb, but not at others. On the other hand, the observation may offer a probe for testing the working mechanisms which govern the behaviors of the light flavors in hadrons which are usually treated as passive spectators in most reactions, as aforementioned.

An interesting discovery has drawn the attention of theorists and it is the observation of baryonic decay $D_s \to p\bar{n}$ [28], which can only occur through W-annihilation topology, since it was supposed to be very suppressed, as aforementioned, for the meson case [29−31]. Chen *et al.* [32] indicated that the short-distance contribution can only make the branching ratio as large as 10^{-6}, which is much smaller than the data. Thus, they suggested that the long-distance contribution via the FSI can enhance this value, so they claimed that it is a dynamical enhancement of the W-annihilation topology in D_s decays. It is worth further study indeed.

2.4 Weak decays of J/ψ

This is another type of rare decays which may provide us with some information about the structure of J/ψ. Generally, J/ψ would decay via strong or electromagnetic interactions. The strong decay is realized via a process where the constituents c and $\bar{c}$ annihilate into three gluons which eventually fragment into hadrons; it is an OZI-suppressed reaction and that is also why J/ψ is a narrow resonance and evaded observation before 1974. It is believed that the decay width is proportional to the wavefunction of J/ψ at the origin, which can be easily obtained by measuring its leptonic decay width. However (see below), a violation of the hadronic helicity selection rule indicates that such a picture may not be completely correct, therefore to investigate the structure of J/ψ (if it has a hybrid component etc.), the study of the weak decay of J/ψ might be very helpful. We calculated the branching ratio of the semi-leptonic decay $J/\psi \to D_s^{(*)} + e^+ + \nu$ [33] in the QCD sum rules and obtained it to be of order of 10^{-10}. Then with the gained parameters, we extended our calculation to the non-leptonic weak decay of J/ψ. The results show that the branching ratio of inclusive weak decays can reach the order the of 10^{-8}, and a special channel $J/\psi \to D_s^{(*)} + \rho$ [34] has a larger branching ratio of about 5.3×10^{-9}, which might be measured by BESIII, B-factories, Super-B and LHCb, as we wish.

In fact, all such rare decays may be important for better understanding of the structure of J/ψ and the governing dynamical mechanisms, even though accurate measurements on them are extremely difficult. We lay our hope on the very large database of the facilities which will be available soon.

3 Hadronic matrix elements

This is probably the most difficult problem in hadron physics which almost covers the whole field of high energy physics, and definitely confronts anybody working in this field. definitely confronted by anybody. The reason is that hadronization occurs at the energy scale below $\Lambda_{\rm QCD}$ where non-perturbative QCD effects dominate, and so far there is no effective way to accurately evaluate the effects yet. Much effort has been made to handle the problem. The simplest way is by using the naive factorization where a hadron, generally a meson, is emitted and can be factorized out from the hadronic transition of one hadron (meson or baryon) to another one. Then the rest transition amplitude can be parametrized by a few form factors which are obtained by fitting data [35]. The transition matrix element can be analytically decomposed into a few terms according to the Lorentz structure and the parity conservation because hadronization is a process where only strong interaction applies. The advantage of this method is that it is simple and since the form factors are fixed by fitting data, they can be extensively applied to study the weak decay rates. It is simple and consistent with data within a rather wide range, however, there are obvious shortcomings. First, the factorization is not always legitimate, and as Buras *et al.* pointed out [36], the matrix element of operator

$$\langle M_1M_2|C_1\bar{q}_{1i}\gamma_\mu(1-\gamma_5)q_{2i}\bar{q}_{3j}\gamma^\mu(1-\gamma_5)q_{4j}$$
$$+C_2\bar{q}_{1i}\gamma_\mu(1-\gamma_5)q_{2j}\bar{q}_{3j}\gamma^\mu(1-\gamma_5)q_{4i}|M\rangle$$

cannot simply be written as the factorized form because a term proportional to $\lambda^a_{ij}\lambda^a_{lm}$ would appear which phenomenologically, causes an effective N_c, and the coefficients are deformed as

$$C_1+C_2/N_c^{\rm eff} \quad \text{or} \quad C_2+C_1/N_c^{\rm eff}$$

where $N_c^{\rm eff}$ is no longer 3 [37, 38].

Moreover, such factorization is based on the spectator mechanism where the transition occurs at the heavier quark leg and where another light component would play the role of a spectator. In this way, the annihilation and W-exchange sub-processes are not properly included. On an other aspect, as we know, the annihilation and W-exchange sub-processes might be important, when dealing with the inclusive processes, especially for evaluating the lifetimes of D mesons [39, 40].

Thanks to the heavy quark effective theory (HQET) where an extra symmetry $SU_f(2)\times SU_s(2)$ is considered, one can reasonably evaluate the transition between two heavy mesons containing b or c quarks. It has already become a criterion for testing the validity of any theoretical calculation where hadronic matrix elements are evaluated, and as when the heavy quark limit is taken, the results must be qualitatively consistent with that obtained by the HQET. Moreover, Georgi generalized this scenario for dealing with transition between two heavy baryons, each of which contains two heavy quarks in terms of the superflavor symmetry [41]. On the other hand, charm is not heavy enough to be treated as a real heavy quark, and the $1/m_c$ corrections may be important, and even for heavier b-quarks, the $1/m_b$ corrections are not negligible in practical computations [42−45]. Therefore, just as the results under large-N_c limit in the $1/N_c$ expansion theory only possess qualitative meaning, the obtained values under the heavy quark limit (i.e. let $m_Q\to\infty$) correspond to the leading order, and corrections must be accounted for while comparing with more accurate experimental data. There are many works to consider when properly evaluating the $1/m_Q$ corrections [42−45].

In the recent years, some effective theories have been developed to justify factorization, especially in B physics. The perturbative QCD (pQCD) is based on the factorization theorem. It states that for a process with large momentum transfer, the physical amplitude can be factorized as a product $\phi\otimes H$ where H is the factor corresponding to the quark-level hard process amplitude which can be calculated in perturbation theory order by order, and where ϕ stands for the soft part. The later one is not calculable in the perturbative way. In general, the product is related to a convolution integration over the wave functions of the initial and final hadrons. In the literature, there are two different versions of pQCD to deal with the factorization, one is the familiar collinear factorization and another is the so-called k_T factorization. The crucial difference between the two schemes may be that k_T factorization is possible in treating the problem of endpoint divergence. The details of the two schemes can be found in relevant literature [46−48]. The soft-collinear effective theory (SCET) is a recently developed effective field theory to simplify the processes containing light energetic hadrons [49−54]. The great development is that the proof of the factorization theorem can be performed at the operator level in the SCET, which is much simpler than the diagrammatic analysis in pQCD. However, on the other hand, except for very few processes, such as the pion transition form factor, which have been proved to be factorizable, most exclusive processes cannot be rigorously proved to be factorizable. Moreover, the factorization proofs are usually limited to the leading order and therefore the application of the factorization theorem is still an assumption. The factorization may be applicable for the decays of bottomed mesons and baryons, but for charmed mesons D and baryons, it is indeed questionable. In D meson decay, the energy of the final light meson is at the order of $\Lambda_{\rm QCD}$, which is not high enough to perform a perturbative analysis. Another question arises from the substantial corrections in power of $\Lambda_{\rm QCD}/m_c$. All these facts make it difficult to apply pQCD or SCET into charm decays. In another approach named as the Transverse Momentum Distribution (TMD), factorization is performed by taking transverse momenta of partons into account. For example in the reaction $\pi^0+\gamma^*\to\gamma$ the form factor can be written as a convolution integral

$$F(Q^2) \sim \phi \otimes S \otimes H(1 + \mathcal{O}(Q^{-2}))$$

where ϕ is the light-cone wavefunction of the pion, S is an additional soft factor and H corresponds to the hard scattering part [55−58].

A very recent work by several authors [59] indicates that for the k_T factorization scheme, at loop-level an extra term which is related to the so-called light-cone divergence appears when a unitary gauge is employed in the calculation and disappears in the Feynman gauge. This explicitly manifests that the the k_T factorization at loop-level is not gauge-independent, so that is violated. This statement is still in dispute. Even though the k_T factorization cannot be a strict theory according to the field theory, it can definitely be treated as a successful phenomenological model and is applied to calculate the transition amplitudes where heavy hadrons are involved. Therefore, generally, we can trust the theoretical results which are obtained based on the k_T factorization. Very, very recently, Li argued that the work [59] might make some calculation mistakes and he presented a result which is free of the light-cone singularity, so the gauge invariance of the k_T factorization is kept [60]. Since it is a serious dispute, we will follow the further development in the interesting regime.

Besides these methods which may stem from the quantum field theory, there are some traditional methods which have been widely applied to calculate the hadronic matrix elements, including for example: the harmonic oscillator model [61], the constituent quark model [62], the constitute quark meson model (CQM) [63−66], the light front quark model [40], the color-singlet model [68, 69], color-octet, and color evaporation [70], and especially the non-relativistic QCD (NRQCD), where an expansion in the powers of velocity of the heavy quark v is naturally made [71]. Besides these phenomenological models whose parameters must be fixed by fitting data, the theoretical framework QCD sum rules [72] is based on quantum field theory where only the perturbative vacuum is replaced by the physical vacuum. Because of the properties of the vacuum, a series of condensates of quark-pair, gluons and quark-gluons etc. are introduced to describe the non-perturbative QCD effects. It has achieved great success in phenomenology. However, on the other side, it is an extrapolation from the region where perturbative QCD works reliably [73]. In the expansion only the operators with lower dimensions are retained and moreover, to extract physical results a reasonable plateau is required, where the threshold values are determined, thus an error of about 15 % is unavoidable.

Quite an amount of phenomenological models has also been employed to calculate the hadronic matrix elements whose energy scale is $\Lambda_{\rm QCD}$. As is well known, the hadronization is fully determined by the non-perturbative QCD effects, and so far there is no reliable way to evaluate them based on quantum field theory or any other first principles. One of the goals of our research in fact, is to determine the mechanism which governs the reaction. Generally, the fundamental theory is the standard model (SM) which at present no one doubts due to its remarkable success, thus we can reliably (if we ignore contributions from new physics beyond the SM) determine the hard factor due to the asymptotic freedom of QCD. To extract important information, such as determining the CKM matrix elements and checking the unitary triangle, by exploring CP violation, one indeed needs to have a more accurate estimation of the hadronic matrix elements, otherwise the physical picture would be contaminated by the inaccuracy. Therefore, to understand the physical world, reliable estimates on the non-perturbative effects are absolutely necessary and all the efforts along the line are worthwhile.

Moreover, one can also use the MIT bag model while taking into account the recoil effects [74], the chiral bag model and even the flux-tube model [75−77], to describe the wavefunctions of the initial and final hadron states when carrying out the calculation of the transition matrix elements.

Indeed, all the models have their own reasonability and advantage, but there are obvious flaws and unreasonability, and because they are not coming from a basic principle, one can never expect that they can be perfect. Therefore on one side, even though the models are not perfect, they have applicability and if they are properly applied, reasonable results should be reached.

As a conclusion, estimation of the hadronic matrix elements is crucially important, but so far, there does not exist a way to fulfil the job yet and one can only apply the available models to estimate them with certain reliability.

4 Final state interaction

Besides the estimate on the hadronic matrix elements which are directly related to the transition, there are secondary reactions, namely the final state interactions (FSIs), which are also very significant for charm-hadron decays. Such processes are due to strong interactions and occur at hadron level, thus also cannot be derived by perturbative QCD. Fortunately, one can use the chiral Lagrangian to evaluate the long-distance effects and we will discuss this issue in this section. Some phenomenological models have been suggested to estimate the FSIs in D meson decays: one-particle-exchange model [78] and the Regge pole model [79].

The final state interaction (FSI) in the charm-tau energy region is very important [80−82]. According to the concept, the FSI can be categorized into the quark-level and hadron level FSI processes. The quark level FSI process refers to the quark interference due to the identical fermion statistics. It was noticed by Stech *et al.* a long time ago while explaining the lifetime difference of D^0 and $D^{\pm}$, and we will come to this subject in later sections. Now let us concentrate on the second category of FSI, i.e., at hadron level. In those processes, the initial hadron first decays into intermediate hadrons (usually two hadrons) and the two hadrons would re-scatter into the final states. Since the re-scattering occurs via strong interaction, the isospin must be conserved.

The re-scattering occurs at hadron level and both of the hadrons are in color-singlets, thus the interaction between the hadrons cannot be described by one or even a few gluon-exchanges and it makes the whole calculation more tricky and uncertain. The responsible effective theory in this field should be the chiral Lagrangian. However, all the coefficients in the Lagrangian cannot be obtained from an underlying theory such as QCD at present, and can only be fixed by fitting data. This brings up very serious problems and uncertainties for the theoretical evaluation.

To evaluate the re-scattering effects, some authors suggested calculating the absorptive part of the triangle diagrams whose internal legs are intermediate hadrons and external lines corresponding to the initial hadron and the final daughter hadrons [83−87]. In fact, the absorptive part of the triangle corresponds to the real Final State rescattering because the two intermediate hadrons which are directly coming from the initial decaying hadron are on their mass shells. The hadron(s) exchanged between the two intermediate hadrons at the t-channel not only needs to possess proper quantum numbers, but also has to conserve energy-momentum, so that it is obviously off-shell. As at the triangle apexes where we apply the effective interaction vertices which are extracted from the chiral Lagrangian, one needs to introduce phenomenological form factors to compensate the off-shell effects. Usually there are various types of the form factors which are widely adopted in literature, and the simplest one is the pole form††,

$$\frac{\Lambda^2 - m^2}{\Lambda^2 - q^2}$$

where q and m are the momentum and mass of the t-channel exchanged hadron, respectively, and Λ is a phenomenological parameter whose value is believed to be close to 1 GeV. It is noted that the form factor can also paly the role of the cut-off in the Pauli-Villas renormalization scheme, so that in the calculations, no ultraviolet and infrared divergences bother us. On an other aspect, this also leads to parameter-dependence and makes theoretical predictions uncertain. Therefore, this calculation can only tell us the order of magnitude for the concerned reaction unless we can use some data as inputs to fix the model parameters.

The importance of FSI can be understood as one studies the decays of $D^0 \to K^0\bar{K}^0$ and $D^0 \to K^+K^-$ [79]. The former process can only occur via a W-exchange diagram which is very suppressed according to the general analysis, moreover, due to a CKM cancelation, this reaction should be proportional to an $SU(3)$ violation, so that should be very small. In comparison, the later one is a Cabibbo favored external emission process and should be overwhelmingly larger than the former one. However the data show that $B(D^0 \to K^+K^-) = (3.84\pm0.3)\times10^{-3}$ and $B(D^0 \to 2K^0_S) \sim (3.7\pm0.7)\times10^{-4}$, which implies that $B(D^0 \to K^0\bar{K}^0)$ is comparable with $B(D^0 \to K^+K^-)$. This can be easily realized via a rescattering of $K^+K^- \to K^0\bar{K}^0$. In our work [79], we showed that it can be realized with the data measured in experiments on KK scattering as inputs. This simple example confirms the importance of FSI in charm physics.

Moreover, the final state interaction provides a strong phase which may lead to CP violation. As is well known, the direct CP violation in decays is

$$\Gamma(A \to B) - \Gamma(\bar{A} \to B) \sim \sin(\alpha_1 - \alpha_2)\sin(\phi_1 - \phi_2)$$

where B may be a CP eigenstate and α, ϕ are the strong phase and weak phase, respectively. It is obvious that there at least exist two independent channels which have different strong and weak phases, otherwise the direct CP asymmetry is zero. Since the final products may originate from another weak process and therefore the reaction can have different weak phases from the direct transition and the FSI can provide a strong phase, i.e. the phase shift in the language of scattering, their interference can result in a non-zero CP asymmetry [82].

To calculate effects of the final state interaction, there are two possible methods that can be adopted, i.e. the Regge-pole model and the hadronic loop. It was discussed that the Regge model may apply in higher energy

†† The usually adopted form of the form factor is $\left(\frac{\Lambda^2 - m^2}{\Lambda^2 - q^2}\right)^n$ where $n = 1$ is the monopole form, $n = 2$ is the dipole form and since $n > 2$, it is a multi-pole form which is seldom selected in literature. Besides, there are exponential and other forms.

Table 1 The first two modes are well measured and the theoretical model parameters are obtained by fitting them [89].

Decay mode	$\rho^0\pi^0$	$K^{*+}K^- + c.c.$	$\phi\eta$	$\phi\eta'$	$\omega\eta$	$\omega\eta'$
BR$\times10^{-3}$(Experiment)[95]	4.2 ± 0.5	5.0 ± 0.4	0.65 ± 0.07	0.33 ± 0.04	1.58 ± 0.16	0.167 ± 0.025
$\mathcal{G}^{PV}(10^{-3}\mathrm{GeV}^{-1})$	2.08 ± 0.25	1.65 ± 0.26	0.89 ± 0.096	0.71 ± 0.086	1.27 ± 0.13	0.46 ± 0.069
$\mathcal{G}_H^{PV}(10^{-3}\mathrm{GeV}^{-1})$(Theory)	6.44	6.01	3.47	5.01	5.33	3.97
$\mathcal{G}^{PV}(10^{-3}\mathrm{GeV}^{-1})$(Theory)	2.08(fitting)	1.65(fitting)	0.93	0.61	0.93	0.43
BR$\times10^{-3}$(Theory)	4.2(fitting)	5.0(fitting)	0.71	0.25	0.84	0.15

regions whereas the hadronic loop method is more suitable in lower energy regions. To a certain accuracy they are consistent.

Generally, as many authors discussed, the two intermediate hadrons are real on-shell particles, therefore, one only needs to calculate the absorptive (i.e., imaginary) part of the triangle diagrams. In fact, it is the real final state interaction by common sense. However, as Suzuki [88] pointed out, the dispersive part of the triangle can also paly a role in influencing the transition amplitude. By calculating the dispersive part of the triangle, we evaluate the branching ratio of $J/\psi \to PV$, where P and V stand for pseudoscalar and vector mesons [89]. This is related to the famous $\rho\pi$ puzzle which will be discussed in later sections. Our strategy is that we calculate the dispersive part of the triangle by keeping the parameter Λ as a free parameter to be determined, and by assuming the $SU(3)$ symmetry, we set the direct transition rate as another free parameter, then by fitting two special channels ($J/\psi \to \rho^0\pi^0$, $K^{+*}\bar{K}^-$ which are more accurately measured) we obtain the two parameters. With these we calculate the branching ratios of other channels and find that the results, which are listed in the table below, are well consistent with data.

It is interesting to note that usually only at lower energy regions, the FSI effects are more significant, but sometimes, the small effects may be also non-negligible. This is the case for studying CP violation. As indicated, the direct CP violation is induced by an interference between at least two channels which have different weak and strong phases. Even though the FSI in certain cases is much smaller than the main contribution which usually comes from the tree level and does not possess a strong phase, FSI then provides a non-zero strong phase, and its weak phase may be completely different from the tree level one, thus an interference between it and the tree amplitude would result in a non-zero CP violation. We calculated such a possibility for the B_c decay [90] and conclude that in this case the pQCD calculation is still valid and the FSI effect only raises a minor contribution to the total decay amplitude, but its strong phase is not zero, so is may help to build up an observable CP violation effect as it interferes with the tree contribution which is calculated in pQCD.

This indicates that the FSI is important, but as aforementioned, the large uncertainties in the whole scenario and the parameters involved in the calculations make the theoretical prediction in accurate, and all these are worth further study. To place the theory in a better shape, we need more information from experiments.

5 Hadronic Helicity suppression in J/ψ decays

It is generally believed that the narrowness of J/ψ is due to the Okubo-Zweig-Iizuka (OZI) rule [91−94]. Violation of the OZI rule would give rise to a non-zero transition rate. The main two-body decays of J/ψ can be categorized as $J/\psi \to PP'$, $J/\psi \to PV$ and $J/\psi \to VV'$ where P, V stand for pseudoscalar and vector mesons. Further, there are definitely other modes, for example, the scalar and axial vector final states. As well, there is also the possibility of decaying into two baryons such as $p\bar{p}$, $\Lambda\bar{\Lambda}$ etc. but the corresponding branching ratios of such channels are rather small [95] and the result is well understood in our theory. To get a better understanding of the OZI rule, people turned to study the radiative decays of orthoquarkonia [96, 97] where only a hadronic transition matrix element is needed. With certain approximations they obtained numerical results which were roughly consistent with data. The process involved a five-point Green's function; the Feynman integration is complicated. Thanks to the developments of calculating techniques, we carried out a full calculation [98] and the results are qualitatively consistent with that obtained by the authors of Refs. [96, 97] and the data available at that moment. The success indicates that our knowledge on the OZI rule might be correct, with a tolerable error. It is also noticed that in the radiative decays, the hadronic effects of J/ψ are included in its wavefunction at the origin, which is obtained by fitting data of the leptonic decays of J/ψ and it must be accurate enough.

To further validate the OZI rule, we calculate the process of $J/\psi \to \pi\pi$, which is an isospin violating one. Usually it is supposed that this process is induced by the electromagnetic interaction, i.e. via $J/\psi \to \gamma^\star \to \pi\pi$. However, as is well known, there are two sources for isospin violation, one is the electromagnetic interaction and another is the mass splitting of u and d quarks. Considering the mass difference of u and d quarks, we calculate the OZI process $J/\psi \to ggg \to \pi\pi$ [99] and get its amplitude which is comparable to the contribution from electromagnetic interaction. Concretely, we find the OZI amplitude is proportional to $(m_u - m_d)/M_{J/\psi}$, which clearly manifests the isospin violation. We should carry out a similar calculation of $J/\psi \to \gamma^* \to \pi\pi$ in the same scenario and compare it with the OZI process in our later work. Then we continue to calculate the branching ratio of $J/\psi \to \rho\pi$, which is supposed to be fully dominated by the OZI process $J/\psi \to ggg \to \rho\pi$. We obtain a branching ratio which is one order smaller than the data. More discussion will be presented in the next section).

In fact, it is not a surprise, because a long time ago, Brodsky and Lepage [100, 101] indicated that as the vector-gluon coupling conserves the quark helicities, the total hadronic helicity is conserved and can only be violated at orders of m/Q or higher, where m is the light quark mass in the final hadrons and Q is the transferred momentum scale. Our numerical results are consistent with this rule, i.e. the process is suppressed by the hadronic helicity conservation and we explicitly show that the amplitude of the transition is proportional to $(m_u + m_d)/M_{J/\psi}$, which confirms the observation of Lepage *et al.* However, the allegation sharply contradicts the data where the branching ratio of $J/\psi \to \rho\pi$ is $(1.69\pm0.15)\%$ [95] and is one of the dominant hadronic modes. To reach a compromise with this obvious discrepancy, Branbilla [100, 101] suggested that either J/ψ contains other constituents, for example it is a hybrid, or there is some unknown mechanism which violates the hadronic helicity conservation and results in a larger transition amplitude. However, both interpretations would receive very crucial challenges because for a long time, people believed that J/ψ is composed of charm and anti-charm quarks and most overwhelming works are based on this picture. The second one also does not seem to work, because if it were true, the mechanism would exist in other channels and influence all theoretical predictions.

Associating with the phenomenon where $\psi' \to \rho\pi$ is peculiarly suppressed, we are inclined to the first interpretation, namely J/ψ is not a pure c-$\bar{c}$ bound state, but ψ' is. We will give more discussions in the later section about the $\rho\pi$ puzzle. On an other aspect, we used to follow Suzuki and consider the long-distance contribution to $J/\psi \to \rho\pi$ and our results show that it is possible to explain data, but the answer is still not satisfactory yet (see later section for more discussions).

6 The $\rho\pi$ puzzle decays

As has been widely discussed, the puzzle has been raised for a long while. In the regular theoretical framework, there should be a relation

$$R = \frac{BR(\psi' \to ggg)}{BR(J/\psi \to ggg)} = \frac{\Gamma(\psi' \to e^+e^-)}{\Gamma(J/\psi \to e^+e^-)} \cdot \frac{\Gamma_t(J/\psi)}{\Gamma_t(\psi')}$$

where Γ_t is the total width. This ratio comes from the fact that if both J/ψ and ψ' are $c - \bar{c}$ bound states, in the hadronic decays, c and $\bar{c}$ annihilate into three gluons which then convert into hadrons, whereas in the leptonic decays, c and $\bar{c}$ annihilate into a virtual photon which turns into a lepton-pair. In this picture, the amplitudes of the hadronic decay which occurs via a three-gluon intermediate state, and the leptonic decay which occurs via a virtual photon intermediate state, are proportional to the wavefunction at origin $\psi(0)$. If everything work well, the ratio should be close to 12 %$\sim$ 14 %, which is called as the 14 % rule (now, it is sometimes called the 12 % rule; anyhow it is a sizable number in contrast to the data.). However, the data tell us that this ratio is much smaller than this value.

Some theoretical interpretations have been proposed. Rosner *et al.* [102] suggested that the quantum number of the observed ψ' may be not a pure 2S state, which is the first radial excited state of the $c\bar{c}$ system, but a mixture of 2S and 1D states. The amplitudes are instead

$$\begin{aligned}\langle\rho\pi|\psi'\rangle &= \langle\rho\pi|2^3S_1\rangle\cos\phi - \langle\rho\pi|1^3D_1\rangle\sin\phi \sim 0\\ \langle\rho\pi|\psi''\rangle &= \langle\rho\pi|2^3S_1\rangle\sin\phi + \langle\rho\pi|1^3D_1\rangle\cos\phi\\ &\sim \langle\rho\pi|2^3S_1\rangle/\sin\phi\end{aligned}$$

where ϕ is fixed as $-27°$ or $12°$ by fitting data. By the destructive interference between the contributions of the two components to the amplitude of $\psi' \to \rho\pi$, the smallness is explained. Suzuki [88] alternatively suggested that the relative phase between the one-photon and gluonic decay amplitudes or hadronic excess in the ψ' decay may result in the small branching ratio. The final state interactions may also give a reasonable explanation [103]. The first proposal can be tested in the decays of $\psi'' \to \rho\pi$ which has not been well measured yet. In Ref. [88], the author suggested that the one-photon amplitude is sizable and it can be tested in some other modes, for example $\psi \to \pi\pi$ if the process is dominated by the electromagnetic interaction. The hadronic excess can also receive tests in the decays of other higher excited states of the ψ- family and even the Υ- family. The final state interaction may play an important role in D and B decays, and also in decays of ψ mesons as suggested in literature. The difficulties are in how to properly evaluate such effects. The final state interaction process is induced by strong interaction at lower energy regions, thus it is governed by the non-perturbative QCD, which is not fully understood in the present theoretical framework yet. People need to invoke some phenomenological models to carry out the calculations. We will give a more detailed discussion on the estimation of the final state interaction, here we only use our results to discuss the puzzle.

In our work [89], we simultaneously considered the FSI and the direct decay of J/ψ into a vector and a pseudoscalar meson and concludeed that both of them contribute to the widths, and their interference should be destructive to explain data. This observation indicates that even though the OZI-forbidden process is sizable, it cannot be consistent with the data. The result implies that the hadronic helicity conservation indeed greatly suppresses the process of $J/\psi \to \rho\pi$ and as the data demand an explanation, one should consider what the origin of the problem is.

In a straightforward calculation based on the SM, we estimate the decay width of the OZI forbidden process $J/\psi \to \rho\pi$ [99], and find that the width is indeed proportional to $(m_q/m_{J/\psi})^2$ which comes from the hadronic helicity suppression factor. Numerically, the branching ratio of $J/\psi \to \rho\pi$ should be smaller than 0.1 %. The same situation appears for $\psi' \to \rho\pi$. It was qualitatively discussed by Brodsky *et al.* [100, 101]. As aforementioned in the last section, to prove the calculation, we recalculate the subprocess $J/\psi \to 3g \to \pi\pi$, which is an isospin violating reaction and is usually supposed to be dominated by the subprocess $J/\psi \to \gamma\pi\pi$ because the EM interaction violates isospin, as is well known. Our result indicates that in the OZI forbidden subprocess the transition amplitude is proportional to $(m_u - m_d)/m_{J/\psi}$, i.e. the mass difference results in the isospin violation instead. Our numerical result is of the same order as the contribution from $J/\psi \to \gamma\pi\pi$. All the results are consistent with our physics picture and are qualitatively reasonable. Therefore, we can trust our calculations for the process $J/\psi \to \rho\pi$. Our numerical results are listed in Table 2.

As indicated in Refs. [100, 101], the structure of J/ψ may be not a pure $c\bar{c}$ charmonium, but consists of other components, such as hybrids $c\bar{c}g$, $c\bar{c}q\bar{q}$ and etc.

To understand the smallness of the ratio R, one can expect that either there is a problem with ψ' as Rosner *et al.* did, or there is something obscure in the J/ψ structure, as Brodsky and many others indicated. Our above numerical results show that even though the hadronic selection rule works in the cases of $J/\psi \to \rho\pi$ and $\psi' \to \rho\pi$, the suppression is not too serious and the theoretical prediction is only one order smaller than the data.

Therefore, a tentative conclusion may be drawn that the $\rho\pi$ puzzle may not be due to the mixing structures of ψ' and ψ'', and neither to an anormal structure of J/ψ itself. It seems that both proposals cannot independently explain the "puzzle"' more complicated mechanisms may be needed.

This is a great challenge to our understanding because the $c\bar{c}$ structure of J/ψ has been recognized almost from the very beginning of its discovery. If it is not a pure $c\bar{c}$, all the previous works in terms of the potential models where many parameters are fixed by fitting data should be re-considered. Or there may be some other mechanisms which were not taken into account, or there may exist contributions from new physics beyond the standard model (SM). However, the latter does not seem to be very promising because the concerned energy range is rather low and SM works perfectly well in explaining the data for most states and processes. Thus, it is obviously inclined to the first proposal that J/ψ is not a pure S-wave bound state of $c\bar{c}$.

It is noted that the calculated results, unless for the distribution ϕ_3, are one order smaller than the data. The same situation happens to ψ', but it is still hard to draw a conclusion that the 14 % rule is due to the existence of higher Fock states in J/ψ or other mechanisms which further suppress the reaction of $\psi' \to \rho\pi$, maybe both. It forms an intriguing challenge to our understanding of the hadron structures. This whole picture also applies to the Υ family, therefore the future experiments would provide hints to finally solve the puzzle.

There have been some theoretical explanations besides that we discussed above. In the work by Mo, Yuan and

Table 2 Decay widths (Γ) of $J/\psi \to \pi^+\rho^- + \pi^-\rho^+$ based on the three distribution functions, ϕ_1, ϕ_2 and ϕ_3, respectively [99].

m_u/MeV	m_d/MeV	$\Gamma(\phi_1)$/MeV	$\Gamma(\phi_2)$/MeV	$\Gamma(\phi_3)$/MeV	exp/MeV
2	2	1.04×10^{-4}	7.21×10^{-5}	5.11×10^{-4}	
3	3	2.36×10^{-4}	1.6×10^{-4}	1.17×10^{-3}	
4	4	4.12×10^{-4}	2.9×10^{-4}	2.08×10^{-3}	$(1.06 \pm 0.08) \times 10^{-3}$
5	5	6.69×10^{-4}	4.54×10^{-4}	3.38×10^{-3}	
6	6	9.75×10^{-4}	6.68×10^{-4}	4.88×10^{-3}	

Wang [104], the authors described the recent status of theoretical research as well as the experimental measurements on the interesting subject.

Interestingly, there is also an alternative opinion towards the subject. Suzuki [88], Zhao [105−107] deny it as a "puzzle", because they consider that the electromagnetic interaction may play an important role in J/ψ decay where $c\bar{c}$ annihilates into a virtual photon which later fragments into hadrons. In the picture, it is supposed that a destructive interference between the contribution of three-gluon and single-photon processes would suppress $\psi' \to \rho\pi$. Since its amplitude can be roughly estimated in terms of the measured rate of J/ψ leptonic decay, there should be a strong constraint on the proposal. Moreover, if it is true, the interference would also appear in other decay modes and the picture should be further investigated and tested by more accurate experimental data available in the future, especially from the BESIII.

7 QCD multi-expansion and strong radiations

In an enlightening paper by Kuang [108], in detail, the author explained the work of Yan and Kuang [109, 110] where they successfully initiated and developed a complete theoretical framework, the multi-expansion method in QCD. The theory properly deals with the emission of light hadrons during heavy quarkonia transitions. Concretely, one mainly studies the processes such as $\Upsilon(nS) \to \Upsilon(mS)+\pi+\pi$ or $\psi(nS) \to \psi(mS)+\pi+\pi$ with $n > m$. In Ref. [108], the author investigated the emission of h_c which was found by CLEOc [111] and invited great interest from theorists [112, 113] as well as experimentalists [114]. The decay width of such transitions can be written as

$$\Gamma(n_I{}^3S_1 \to n_F{}^3S_1) = |C_1|^2 G |f^{l,P_I,P_F}_{n_I,l_I,n_F,l_F}|^2$$

where $|C_1|^2$ is a constant to be determined and it comes from the hadronization of gluons into pions, G is the phase space factor, $f^{l,P_I,P_F}_{n_I,l_I,n_F,l_F}$ is an overlapping integral over the concerned hadronic wave functions, and their concrete forms were given in Refs. [109, 110] as

$$f^{l,P_I,P_F}_{n_I,l_I,n_F,l_F} = \sum_K \frac{\int R_F(r) r^{P_F} R^*_{Kl}(r) r^2 \mathrm{d}r \int R^*_{Kl}(r') r'^{P_I} R_I(r') r'^2 \mathrm{d}r'}{M_I - E_{Kl}}$$

where n_I, n_F are the principal quantum numbers of the initial and final states, l_I, l_F are the angular momenta of the initial and final states, l is the angular momentum of the color-octet $q\bar{q}$ in the intermediate state, P_I, P_F are the indices related to the multipole radiation, and for the E1 radiation P_I, P_F=1 and $l = 1$. R_I, R_F and R_{Kl} are the radial wave functions of the initial and final states, M_I is the mass of the initial quarkonium and E_{Kl} is the energy eigenvalue of the intermediate hybrid state.

The framework provides a more elegant way to deal with the long-distance QCD effects even though it only concerns transitions between states of heavy quarkonia. Moreover, between two gluon emissions the intermediate state is a hybrid state which definitely is a subject to draw the interest of all high energy physicists. Just as aforementioned, the QCD theory predicts the existence of exotic states, such as glueball, hybrid, tetraquark and pentaquark etc., and at least does not exclude their existence. However, so far none of such exotic states have been experimentally identified yet, so that any direct or indirect information about the exotic states must be valuable. Since the subject is heavy quarkonia, the potential model is reasonable for describing their structures. Recently, many works have been devoted to studing the potential which can describe the hybrids, which are composed of heavy quark, anti-quark and a gluon. In this case the quark and heavy quark reside in a color octet to keep the hybrid meson in a color singlet, so that the Coulomb potential between them is repulsive. The recent literature suggests a form by Swanson and Szczepaniak [115]

$$V(r) = br + \frac{\pi}{r}\left(1 - \mathrm{e}^{-fb^{1/2}r}\right)$$

where the concerned parameters were given in Ref. [115]. Alternatively, Allen *et al.* proposed another potential form which includes a repulsive Coulomb piece as [116]

$$V(r) = \frac{\kappa}{8} + \sqrt{(br)^2 + 2\pi b} + V_0$$

where V_0 is the zero-point energy and other parameters were also depicted in Ref. [116].

When the works of Refs. [108−110] were done, there were not many data about the transitions available, so that the authors assumed that $\psi(4.03)$ is the ground state of hybrid $|c\bar{c}g\rangle$ and accordingly obtained the concerned parameters in the potentials listed above. Recently, thanks to the great work of Belle, Barbar and CLEOc, much more data have been collected and they enable us to take an inverse strategy to study the problem.

In our strategy, we treat the parameters in the hybrid potential as free parameters to be determined. Using all the available data on $\Upsilon(nS) \to \Upsilon(mS) + \pi + \pi$ and $\psi(nS) \to \psi(mS) + \pi + \pi$ as inputs, we apply the χ^2 analysis with the form of $\bar{\chi}^2$ defined in Ref. [117] as

$$\bar{\chi}^2 = \sum_i \frac{(W_i^{\rm th} - W_i^{\rm exp})^2}{(\Delta W_i^{\rm exp})^2}$$

where i represents the i-th channel, $W_i^{\rm th}$ is the theoretical prediction on the width of channel i, $W_i^{\rm exp}$ is the corresponding experimentally measured value, and $\Delta W_i^{\rm exp}$ is the experimental error.

Carrying out all the procedures, we have obtained that the masses of ground states of the charmonium family and Υ family are 4.23 GeV (for charmonium) and 10.79 GeV (for bottonium). It is noted that they are not the physical states which are experimentally observed. This is quite understandable because the present knowledge may suggest that the glueballs and hybrids may not exist as real resonances with fixed masses and widths, but mix with hadrons with regular valence-quark compositions [118−121], and the physical states are the eigenstates of the mass matrices.

However, from other aspects, these results are not accurate due to large uncertainties of the experimental data. Especially, when we reached the results, $\Upsilon(5S)$ was not measured yet, and we did not include its transition to lower the Υ members.

Last year, the Belle Collaboration reported their measurements on $\Upsilon(5S) \to \Upsilon(1S)\pi^+\pi^-$ and $\Upsilon(5S) \to \Upsilon(2S)\pi^+\pi^-$ with decay widths of 0.59 ± 0.04(stat.) $\pm$ 0.09(syst.) MeV and 0.85 ± 0.07(stat.)±0.16(syst.) MeV. These values are about two orders larger than the previously measured partial widths for dipion transitions between lower Υ resonances [122].

Meng and Zhao suggested that the anomalous enhancement is due to the final state interaction [123, 124]. Namely, because $\Upsilon(5S)$ and $\Upsilon(4S)$ are heavy and above the production threshold of $B\bar{B}$, they may therefore first decay into a $B\bar{B}$ pair and then by a re-scattering, $B-\bar{B}$ would turn into $\Upsilon(mS)+\pi\pi$ with $m \leqslant 3$. By this picture and fixing the concerned parameters within reasonable ranges, the enhancement may be understood. If it is true, one cannot further use the data of $\Upsilon(5S) \to \Upsilon(mS)+\pi\pi$ in our above calculations because the re-scattering contribution contaminates the whole picture and one is no longer enable to gain direct information about the hybrid intermediate state at all.

In the calculation, we have found a strange phenomenon that for $\Upsilon(nS) \to \Upsilon(mS)+\pi+\pi$, the results are pretty stable, however for $\psi(nS) \to \psi(mS)+\pi+\pi$ there exists a cancelation among large numbers with smaller numbers remaining. This is due to the closeness of the charmonia masses and the hybrids, therefore the results on the charmonia transitions are not very reliable.

On an other aspect, such unstableness may be the result of a mistreatment of the charmonia transition. If the final state interaction is important as the authors of Refs. [123, 124] suggested, for bottonia transitions, it would also apply to the charmonia transition when masses of $\psi(nS)$ are above the production threshold of $D\bar{D}$. Taking into account such FSI effects, one may re-extract information about the direct transitions. As one did for the potential model, one may expect that the parameters are universal for cases b and c and then we can reduce theoretical uncertainties in the calculation.

Indeed, such information is very necessary for determining the model parameters and even judge the whole scenario of hybrids. Therefore, we are expecting more data in the charmonia energy regions to be collected in Babar, Belle and even the LHCb, as well as improvements in the theory.

8 The X, Y and Z resonances

The QCD theory and quark model have been proven to be very successful, however, it is by no means the end of the story. Both QCD and the quark model have soft bellies where many problems are not answered yet. Interestingly, the two aspects in the quark model and QCD are connected to each other. In the QCD theory, thanks to the asymptotic freedom, the perturbation can be believably applied to evaluate any high energy processes, however, on other side, the low energy processes are governed by the non-perturbative QCD effects for which so far there is a lack of any reliable way to deal with. We have already briefly discussed this issue in previous sections. Unfortunately, many real physics quantities are related to the low processes, such as the fragmentation in high energy collisions and hadronic transition form factors in hadron decays. On an other aspect, the quark model demands that mesons are composed of a pair of quark and anti-quark, baryons consist of three quarks (or anti-quarks), and the QCD interaction (i.e. strong interaction) binds all the constituents into hadrons. Both of the theories do not prohibit the existence of exotic states, such as glueballs, hybrids and multi-quark states (tetraquark, pentaquark etc.) or even favor their existence. The pentaquark was a hot topic for a while as several groups claimed that pentaquarks containing an anti-strange quark or anti-charm quark were observed, then new data from most of the major labs gave completely negative results. Do they really exist or mix with regular hadrons as suggested in literature [125, 126]? That is still an unsolved question.

Recently, the Babar, belle, CLEO and BES reported many newly observed resonances which were randomly named as X, Y and Z particles. In a review paper, Godfrey and Olsen discussed this issue in some detail [127]. There are many theoretical works devoted to the exciting field.

$D^*_{sJ}(2317)$, $D_{sJ}(2460)$, $D_{sJ}(2860)$ and $D_{sJ}(2715)$

The discoveries of mesons $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ [128−138] whose spin-parity structure are respectively 0^+ and 1^+, have attracted great interest from both theorists and experimentalists of high energy physics, because they seem to be exotic. Bardeen *et al.* supposed that $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ are $(0^+, 1^+)$ chiral partners of D_s and D^*_s [139−141], i.e. p-wave excited states of D_s and D^*_s [142]. By studying the mass spectra, Beveren and Rupp suggested that $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ are made of c and $\bar{s}$ [143, 144]. With the QCD spectral sum rules, Narison calculated the masses of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ by assuming them to be quark-antiquark bound states and obtained results consistent with the experimental data within a wide error range [145]. Very recently, considering the contribution of the DK continuum in the QCD sum rules, Dai *et al.* obtained the mass of $D^*_{sJ}(2317)$, which is consistent with experiments [146]. Meanwhile, some authors suggested that $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ may be of four-quark structure [147−152].

Thus to clarify the mist of the structures of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$, serious theoretical works are needed. The studies of the productions and decays of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ are very interesting topics. Several groups have calculated the strong and radiative decay rates of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ in different theoretical approaches: the Light Cone QCD Sum Rules, the Constituent Quark model, Vector Meson Dominant (VMD) ansatz, the constituent quark meson model, etc. [153−161]. The authors of Refs. [150, 162] also calculated the rates based on the assumption that $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ are in non-$c\bar{s}$ structures. Their predictions on the $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ decay rates are obviously a few orders larger than that obtained by assuming the two-quark structure. Recently, Faessler *et al.* studied the same subject, assuming $D^*_{sJ}(2317)$ as a DK molecule state using an effective Lagrangian approach [163].

The semileptonic decay of B_s is one of the ideal platforms to study the productions of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$. Especially, the Large Hadron Collider (LHC) will be running in 2008, which can produce large amounts of B_s. Thus, the measurements on $B_s \to D_{sJ}(2317, 2460)l\bar{\nu}$ would be realistic. In Ref. [164], the author calculated rate of $B_s \to D_{sJ}(2317, 2460)l\bar{\nu}$ in terms of the QCD sum rules and HQET. Recently, authors of Refs. [165, 166] completed the calculations of $B_s \to D_{sJ}(2317, 2460)l\bar{\nu}$ semileptonic decays in the QCD sum rules and obtained large branching ratios. However, the results obtained by authors of Refs. [165, 166] are one order smaller than those given in Ref. [164]. In Ref. [167], authors studied the same topic in terms of the Constituent Quark Meson (CQM) model. The branching ratios of $B_s \to D_{sJ}(2317, 2460)l\bar{\nu}$ estimated by the authors of Ref. [167] and that obtained in terms of the QCD sum rules [165, 166] are of the same order of magnitude.

In Ref. [168], the authors used the heavy quark effective theory (HQET) and a non-relativistic model to evaluate the production rate of $D^*_{sJ}(2317)$ in the decay of $\psi(4415)$, and found that it is sizable and may be observable at BES III and CLEO, if it is a p-wave excited state of $D_s(1968)$.

Because $D_{sJ}(2632)$ was only observed by the SELEX collaboration [169], but not by Babar [170], Belle [171] and FOCUS [172], its existence is still in dispute, so we do not intend to include $D_{sJ}(2632)$ in this review.

In the summer of 2006, the Babar collaboration observed a new $c\bar{s}$ state $D_{sJ}(2860)$ with a mass $2856.6 \pm 1.5 \pm 5.0$ MeV and width $\varGamma = (48 \pm 7 \pm 10)$ MeV. Babar observed it only in the $D^0K^+, D^+K^0_S$ channel and found no evidence of $D^{*0}K^+$ and $D^{*+}K^0_S$. Thus its quantum numbers should correspondingly be $J^P = 0^+, 1^-, 2^+, 3^-, \cdots$ [173]. At the same time, the Belle collaboration reported a broader $c\bar{s}$ state $D_{sJ}(2715)$ with $J^P = 1^-$ in $B^+ \to \bar{D}^0D^0K^+$ decay [174, 175]. Its mass is $2715 \pm 11^{+11}_{-14}$ MeV and width $\varGamma = (115 \pm 20^{+36}_{-32})$ MeV.

According to the heavy quark effective field theory, heavy mesons form doublets. For example, we have one s-wave $c\bar{s}$ doublet $(0^-, 1^-) = (D_s(1965), D^*_s(2115))$ and two p-wave doublets $(0^+, 1^+) = (D^*_{sJ}(2317),\ D_{sJ}(2460))$ and $(1^+, 2^+) = (D_{s1}(2536), D_{s2}(2573))$ [95]. The two d-wave $c\bar{s}$ doublets $(1^-, 2^-)$ and $(2^-, 3^-)$ have not been observed yet. The possible quantum numbers of $D_{sJ}(2860)$ may be $0^+(2^3\mathrm{P}_0)$, $1^-(1^3\mathrm{D}_1)$, $1^-(2^3\mathrm{S}_1)$, $2^+(2^3\mathrm{P}_2)$, $2^+(1^3\mathrm{F}_2)$ and $3^-(1^3\mathrm{D}_3)$. The $2^3\mathrm{P}_2$ $c\bar{s}$ state is expected to lie around $(2.95 - 3.0)$ GeV while the mass of the $1^3\mathrm{F}_2$ state will be much higher than 2.86 GeV.

$D_{sJ}(2860)$ was proposed as the first radial excitation of $D^*_{sJ}(2317)$ [176, 177], or as a $J^P = 3^-$ $c\bar{s}$ state [178] or as $c\bar{s}(2P)$ state [179]. By the potential model, one can see that $D_{sJ}(2715)$ sits exactly at the position predicted by the quark model, at 2715 MeV if it is a $2^3\mathrm{S}_1$ $c\bar{s}$ state [180]. The 1^- state should lie around 2721 MeV if a $(1^+, 1^-)$ $c\bar{s}$ chiral doublet is formed [181].

In Ref. [182], authors investigated the strong decays of the excited $c\bar{s}$ states using the $^3\mathrm{P}_0$ model. After comparing the theoretical decay widths and decay patterns with the available experimental data, they tend to conclude: (1) $D_{sJ}(2715)$ is probably the $1^-(1^3\mathrm{D}_1)$ $c\bar{s}$ state although the $1^-(2^3\mathrm{S}_1)$ assignment is not completely excluded; (2) $D_{sJ}(2860)$ seems unlikely to be the $1^-(2^3\mathrm{S}_1)$ and $1^-(1^3\mathrm{D}_1)$ candidate; (3) $D_{sJ}(2860)$ as either a $0^+(2^3\mathrm{P}_0)$ or $3^-(1^3\mathrm{D}_3)$ $c\bar{s}$ state is consistent with the experimental data; (4) experimental search of $D_{sJ}(2860)$ in the channels $D_s\eta$, DK^*, D^*K and $D_s^*\eta$ will be crucial to distinguish the above two possibilities.

In Ref. [183], the strong decay of the D wave $c\bar{s}$ meson to the light pseudoscalar meson is studied in the framework of the light cone QCD sum rule (LCQSR).

Recently, Dubynskiy and Voloshin proposed an interesting picture to explain the newly observed rich family of X, Y and Z. They suggested that the charmonium states, such as J/ψ, $\psi(2S)$, η_c, can be bound inside light hadronic matter, especially inside higher resonances made from light quarks and gluons, and they named such states as hadro-charmonium [184, 185]. Definitely, this picture should undergo some more serious theoretical and experimental tests.

Indeed, this is a wide world to be explored, which may help to validate the quark model and the low-energy QCD principles as well as all the working phenomenological models.

9 D^0-$\bar{D}^0$ mixing

This is an extremely interesting subject since a sizable mixing of D^0-$\bar{D}^0$ generally implies existence of new physics beyond the standard model.

In the SM, mixing of particle and anti-particle, such as K^0-$\bar{K}^0$, $D^0\bar{D}^0$ and $B^0_{(s)}$-$\bar{B}^0_{(s)}$, occurs via the box-diagrams [186, 187]. The calculation is standard based on the the GIM mechanism [188]. The contribution of the box diagram is proportional to m_i^2/m_W^2, where m_i is the mass of the exchanged quarks in the box and the CKM matrix elements [189]. Thus for the B^0-$\bar{B}^0$ and B_s-$\bar{B}_s$ mixing, the exchanged quark is the top quark and the factor $m_t^2/m_W^2 > 1$ becomes an enhancement, so that the mixing is large; such mixing has been reported to have been observed a long time ago and in history it played an important role in hinting that the top quark might be heavier than the W-boson. However, for D^0-$\bar{D}^0$, the exchanged quark can only be a b-quark, which is much lighter than the top quark, so that the resultant mixing must be very small. The D^0-$\bar{D}^0$ mixing has indeed been measured by the Babar [190, 191] and Belle [192, 193] collaborations. Therefore, it implies the possible existence of new physics beyond SM. There have been some theoretical suggestions about the mechanisms which may cause a sizable D^0-$\bar{D}^0$ mixing, for example, via a FCNC process in the non-universal Z' model [194] or the unparticle model [195, 196].

The eigenstates of the mass matrix are [197, 198]

$$|D_H\rangle = p|D^0\rangle + q|\bar{D}^0\rangle$$
$$|D_L\rangle = p|D^0\rangle - q|\bar{D}^0\rangle$$

with $|p|^2 + |q|^2 = 1$, and the corresponding eigenvalues are

$$\begin{aligned}&(m_H - m_L) - \mathrm{i}(\Gamma_H - \Gamma_L)/2\\ &= 2\sqrt{(M_{12} - \mathrm{i}\Gamma_{12}/2)(M_{12}^* - \mathrm{i}\Gamma_{12}^*/2)}\end{aligned}$$

where M_{12} and Γ_{12} are the off-diagonal matrix elements and are obtained by calculating the box diagrams; it is noted that both of them are complex. In SM they are small, thus a sizable non-zero mass and life difference of the two physical states $|D_S\rangle$ and $|D_L\rangle$ must be caused by new physics, as aforementioned.

Since the mass and life differences of the two eigenstates are not very large, it is hard to measure it as one did for the K^0-$\bar{K}^0$ system, and we can investigate the evolution process of the D^0-$\bar{D}^0$ system which, in general, is produced in B-decays or at higher excited states of the ψ family. One has [197]

$$|D_p^0(t)\rangle = g_+(t)|D^0\rangle + \frac{q}{p}g_-(t)|\bar{D}^0\rangle$$
$$|\bar{D}_p^0(t)\rangle = \frac{p}{q}g_-(t)|D^0\rangle + g_+(t)|\bar{D}^0\rangle$$

and

$$g_\pm = \frac{1}{2}\mathrm{e}^{-imt-\frac{\gamma}{2}t}\left(\mathrm{e}^{\mathrm{i}\frac{\Delta m}{2}t-\frac{\Delta\gamma}{4}t} \pm \mathrm{e}^{-\mathrm{i}\frac{\Delta m}{2}t+\frac{\Delta\gamma}{4}t}\right)$$

and $\Delta m = m_H - m_L$, $\Delta\gamma = \gamma_H - \gamma_L$. The important parameters are

$$x = \frac{\Delta m}{\gamma}, \quad y = \frac{\Delta\gamma}{\gamma}$$

where γ is the average lifetime of D_H and D_L.

In analogy to the K-system, there exists the indirect CP violation which can be observed in the time evolution of the system. The direct CP violation will be discussed later.

The data of the Babar collaboration about the D^0-$\bar{D}^0$ mixing [190, 191] are $x'^2 = [-0.22 \pm 0.30(\mathrm{stat.}) \pm 0.21(\mathrm{syst.})]\times 10^{-3}$ and $y' = [9.7\pm 4.4(\mathrm{stat.})\pm 3.1(\mathrm{syst.})]\times 10^{-3}$, where $x' = x\cos\delta_{K\pi} + y\sin\delta_{K\pi}$; $y' = -x\sin\delta_{K\pi} + y\cos\delta_{K\pi}$ and $\delta_{K\pi}$ is the strong phase between the Cabibbo favored (CF) and doubly Cabibbo suppressed (DCS) amplitudes. $x'$is consistent with zero and y' obviously deviates from zero. The data given by the Belle collaboration [192, 193] are $y_{CP} = [1.31 \pm 0.32(\mathrm{stat.}) \pm$

0.25(syst.)]%, where y_{CP} is defined as $y_{CP} = y\cos\phi - \frac{1}{2}A_M \sin\phi$. When CP is conserved, $A_M = \phi = 0$, and the results are consistent with that obtained by the Babar collaboration.

The unparticle model was first proposed by Georgi [201, 202]. Georgi argued that operators O_{BZ} made of BZ fields in the scale invariant sector may interact with operators $O_{\rm SM}$ of dimension $d_{\rm SM}$ made of Standard Model (SM) fields at some high energy scale by the exchange of particles of large masses, $M_{\mathcal{U}}$, with the generic form $O_{\rm SM}O_{BZ}/M_{\mathcal{U}}^k$. At another scale $\Lambda_{\mathcal{U}}$, the BZ sector induces dimensional transmutation; below that scale the BZ operator O_{BZ} matches onto unparticle operator $O_{\mathcal{U}}$ with dimension d_U and the unparticle interaction with SM particles at low energy has the form

$$\lambda \Lambda_{\mathcal{U}}^{4-d_{\rm SM}-d_U} O_{\rm SM} O_{\mathcal{U}}$$

In the SM, the weak phase comes from the CP phase in the CKM matrix. Since the SM contribution to M_{12} and Γ_{12} can be neglected, the weak phase must be induced in the new physics. We would like to point out some salient features of the unparticle contribution to M_{12}^U and Γ_{12}^U due to an extra phase factor $\mathrm{e}^{-\mathrm{i}\pi d_U}$ in the new scenario. We note that M_{12}^U can have both positive and negative signs depending on the value of d_U due to the factor $\cot(\pi d_U)$, therefore if information about the sign can be obtained from other theoretical considerations or experimental data, the dimension d_U would be restricted. There may be a sizeable contribution to Γ_{12} at tree level, which is not possible for the usual mode where heavy particles are exchanged at tree level. For d_U equal to half integers, there is no contribution to M_{12}, but there is to Γ_{12}. Another salient feature is that the unparticle contribution to the ratio $M_{12}/(\Gamma_{12}/2)$ is related to the unparticle dimension parameter d_U by

$$\frac{M_{12}^U}{\Gamma_{12}^U/2} = \cot(\pi d_U) \tag{1}$$

If the unparticle contribution dominates meson and antimeson oscillation then the measurements of M_{12} and Γ_{12} would provide a possible way to determine the dimension parameter d_U. In this case, we may gain valuable information about the dimension of unparticle d_U which so far cannot be directly obtained when mapping the operator O_{BZ} onto $O_{\mathcal{U}}$ and remains as a free and adjustable parameter in most phenomenological research works.

Definitely, the contributions of new physics would be added to that of SM. The new physics effects also contribute to K^0-$\bar{K}^0$ and $B^0_{(s)}$-$\bar{K}^0_{(s)}$ mixing, however, in those cases the SM contribution obviously dominates that from new physics and the effects of new physics would be smeared out as the measurements are not very accurate. By contrast, the SM contribution to D^0-$\bar{D}^0$ is negligible and all contribution comes from the new physics, thus measurements on it may provide us with an ideal place for gaining valuable information about new physics which is indeed the goal of all high-energy physicists.

To make an accurate measurement, one needs a longer flight time before the D meson decays, thus the main labs for observing D^0-$\bar{D}^0$ would be the two B-factories, and the LHCb which will be running soon will offer us another ideal spot to carry out such measurements. Even though for this observation the BESIII does not have advantages in the kinematics because the linear momenta of the produced D mesons is small and the relativistic time dilation does not apply, as the builder promised, the database will be very large and may greatly enhance the statistics, so it may also be a possible lab for observing the mixing effects and explore new physics. Li and Yang [199, 200] studied how to properly extract information about the mixing from data and make an efficient analysis.

It is natural to ask if one can observe CP violation in the D-system. If D_H and D_L are not CP eigenstates, there could be an indirect CP violation, however, since D decays faster and there are many channels available, one cannot determine the indirect CP violation as easily as in the K system by measuring η_{+-} and η_{00} at different distances.

The direct CP violation is defined as

$$C_f(t) = \frac{\Gamma(D_p^0(t) \to f) - \Gamma(\bar{D}_p^0(t) \to \bar{f})}{\Gamma(D_p^0(t) \to f) + \Gamma(\bar{D}_p^0(t) \to \bar{f})}$$

which is a time-dependent measurable quantity. Some details are given in Ref. [197]. So far, there is still no report on the observation of CP violation at the D-system, even though the direct CP violation has been measured to be non-zero at the B-system. One may expect to make progress along the line in the future.

10 Charmed baryons

Let us present our notations for the excited charmed baryons. Inside a charmed baryon there are one charm quark and two light quarks (u, d or s). It belongs to either the symmetric 6_F or antisymmetric $\bar{3}_F$ flavor representation (see Fig. 1). For the S-wave charmed baryons, the total color-flavor-spin wave function and color wave function must be symmetric and antisymmetric, respectively. Hence, the spin of the two light quarks is $S = 1$ for 6_F or $S = 0$ for $\bar{3}_F$. The angular momentum and parity of the S-wave charmed baryons are $J^P = \frac{1}{2}^+$ or $\frac{3}{2}^+$ for 6_F and $J^P = \frac{1}{2}^+$ for $\bar{3}_F$. The S -wave charmed baryons are listed in Fig. 1, where we use the star to denote $\frac{3}{2}^+$ baryons and the prime to denote the $J^P = \frac{1}{2}^+$ baryons in the 6_F representation.

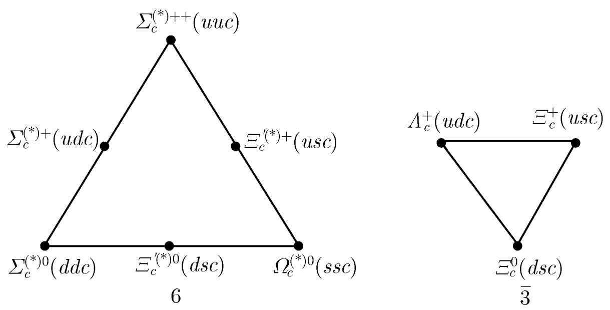

Fig. 1 The $SU(3)$ flavor multiplets of charmed baryons.

In Fig. 2 we present our notations and conventions for the P-wave charmed baryons. l_ρ is the orbital angular momentum between the two light quarks while l_λ denotes the orbital angular momentum between the charm quark and the two light quark system. We use the prime to label the Ξ_{cJ_l} baryons in the 6_F representation and the tilde to discriminate the baryons with $l_\rho = 1$ from that with $l_\lambda = 1$.

$$f_S(6):\ L=1\otimes S_{q_1q_2}=1 \longrightarrow \begin{cases} J_l=0:\ \Sigma_{c0}(\frac{1}{2}^-) & \Xi'_{c0}(\frac{1}{2}^-) \\ J_l=1:\ \Sigma_{c1}(\frac{1}{2}^-,\frac{3}{2}^-) & \Xi'_{c1}(\frac{1}{2}^-,\frac{3}{2}^-) \\ J_l=2:\ \Sigma_{c2}(\frac{3}{2}^-,\frac{5}{2}^-) & \Xi'_{c2}(\frac{3}{2}^-,\frac{5}{2}^-) \end{cases}$$

$$f_A(\bar{3}):\ L=1\otimes S_{q_1q_2}=0 \Longrightarrow J_l=1:\ \Lambda_{c1}(\tfrac{1}{2}^-,\tfrac{3}{2}^-)\quad \Xi_{c1}(\tfrac{1}{2}^-,\tfrac{3}{2}^-)$$

(a) $l_\rho=0,\ l_\lambda=1$

$$f_A(\bar{3}):\ L=1\otimes S_{q_1q_2}=1 \longrightarrow \begin{cases} J_l=0:\ \tilde{\Lambda}_{c0}(\frac{1}{2}^-) & \tilde{\Xi}_{c0}(\frac{1}{2}^-) \\ J_l=1:\ \tilde{\Lambda}_{c1}(\frac{1}{2}^-,\frac{3}{2}^-) & \tilde{\Xi}_{c1}(\frac{1}{2}^-,\frac{3}{2}^-) \\ J_l=2:\ \tilde{\Lambda}_{c2}(\frac{3}{2}^-,\frac{5}{2}^-) & \tilde{\Xi}_{c2}(\frac{3}{2}^-,\frac{5}{2}^-) \end{cases}$$

$$f_S(6):\ L=1\otimes S_{q_1q_2}=0 \Longrightarrow J_l=1:\ \tilde{\Sigma}_{c1}(\tfrac{1}{2}^-,\tfrac{3}{2}^-)\quad \tilde{\Xi}'_{c1}(\tfrac{1}{2}^-,\tfrac{3}{2}^-)$$

(b) $l_\rho=1,\ l_\lambda=0$

Fig. 2 The notations for P-wave charmed baryons. $f_S(6_F)$ and $f_A(\bar{3}_F)$ denote the $SU(3)$ flavor representation. $S_{q_1q_2}$ is the total spin of the two light quarks. L denotes the total orbital angular momentum of charmed baryon system.

The notation for D-wave charmed baryons is more complicated (see Fig. 3). Besides the prime, l_ρ and l_λ defined above, we use the hat and check to denote the charmed baryons with $l_\rho = 2$ and $l_\rho = 1$ respectively. For the baryons with $l_\rho = 1$ and $l_\lambda = 1$, we use the superscript L to denote the different total angular momentum in $\check{\Lambda}^L_{cJ_l}$, $\check{\Sigma}^L_{cJ_l}$ and $\check{\Xi}^L_{cJ_l}$.

The Babar and Belle collaborations observed several excited charmed baryons: $\Lambda_c(2880, 2940)^+$, $\Xi_c(2980, 3077)^{+,0}$ and $\Omega_c(2768)^0$ [203−208]. We collect the experimental information from these recently observed hadrons in Table 3. Their quantum numbers have not been determined, except for $\Lambda_c(2880)^+$ [209].

In the past decades, there have been some research work on heavy baryons [210−216]. However, these new observation inspired new investigations of these states [217−222]. In Ref. [220], the authors studied the $\Lambda_c(2940)^+$ and its possible decay modes, assuming $\Lambda_c(2940)^+$ to be a $D^{*0}p$ molecular state [220]. Cheng *et al.* calculated the strong decays of newly observed charmed mesons in the framework of the heavy hadron chiral perturbation theory (HHChPT) [221]. In order to understand their structures using the present experimental information, in Refs. [223, 224], the strong decay pattern of the excited charmed baryons are studied systematically in the framework of the $^3\mathrm{P}_0$ strong decay model. After comparing the theoretical results with the available experimental data, their favorable quantum numbers and assignments are obtained in the quark model.

It is an interesting field. There was a puzzle in that the lifetime of B^0 is close to that of $B^\pm$ ($\tau(B^0) = (1.530 \pm 0.009) \times 10^{-12}$ s, $\tau(B^\pm) = (1.638 \pm 0.011) \times 10^{-12}$s) whereas the lifetime of $D^\pm$ is almost double of that of D^0 ($\tau(D^0) = (410.1 \pm 1.5) \times 10^{-15}$ s, $\tau(D^\pm) = (1040 \pm 7) \times 10^{-15}$ s). By the heavy quark effective theory, the total width of a meson is mainly determined by the total width of the heavy quark constituent in the meson, which is proportional to $m_Q^5 \times V_{CKM}$ where V_{CKM} are the corresponding CKM matrix elements. Obviously, the weak decay of the b-quark is Cabibbo suppressed while for the charm quark is Cabibbo favored. The obvious difference between the B and D lifetimes is excellently explained by Bigi *et al.* [225−230]. Another puzzle is the lifetime difference between the Λ_b and B meson, which used to be as large as 0.29, however the recent measurement alleviates the discrepancy as $\tau(\Lambda_b)/\tau(B^0) = 1.041 \pm 0.057$ [231], and is close to the theoretical evaluation [232−234]. However, in the theoretical calculation, one needs to evaluate the hadronic matrix element in analog to the discussion given in previous sections about the meson case, and the estimate still has a large uncertainty, so we do not think that the problem is completely solved.

A more intriguing issue is that the lifetime of $\tau(\Lambda_c)(\approx (200 \pm 6) \times 10^{-15}$ s) is much smaller than that of D^0 and $D^\pm$, and it implies that the light quark cannot com-

$f_S(6)$: $L=2\otimes S_{q_1q_2}=1$
- $J_l=1$: $\Sigma_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ $\Xi'_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$
- $J_l=2$: $\Sigma_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ $\Xi'_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$
- $J_l=3$: $\Sigma_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$ $\Xi'_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$

$f_A(\bar{3})$: $L=2\otimes S_{q_1q_2}=0 \Longrightarrow J_l=2$: $\Lambda_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ $\Xi'_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$

(a) $l_\rho=0,\ l_\lambda=2$

$f_S(6)$: $L=2\otimes S_{q_1q_2}=1$
- $J_l=1$: $\hat{\Sigma}_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ $\hat{\Xi}'_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$
- $J_l=2$: $\hat{\Sigma}_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ $\hat{\Xi}'_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$
- $J_l=3$: $\hat{\Sigma}_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$ $\hat{\Xi}'_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$

$f_A(\bar{3})$: $L=2\otimes S_{q_1q_2}=0 \Longrightarrow J_l=2$: $\hat{\Lambda}_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ $\hat{\Xi}_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$

(b) $l_\rho=2,\ l_\lambda=0$

$f_A(\bar{3})$
- $L=1\otimes S_{q_1q_2}=1$
  - $J_l=0$: $\Lambda^1_{c0}(\frac{1}{2}^+)$ $\Xi^1_{c0}(\frac{1}{2}^+)$
  - $J_l=1$: $\Lambda^1_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ $\Xi^1_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$
  - $J_l=2$: $\Lambda^1_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ $\Xi^1_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$
- $L=0\otimes S_{q_1q_2}=1 \Longrightarrow J_l=1$: $\Lambda^0_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ $\Xi^0_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$
- $L=2\otimes S_{q_1q_2}=1$
  - $J_l=1$: $\Lambda^2_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ $\Xi^2_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$
  - $J_l=2$: $\Lambda^2_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ $\Xi^2_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$
  - $J_l=3$: $\Lambda^2_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$ $\Xi^2_{c3}(\frac{5}{2}^+,\frac{7}{2}^+)$

$f_S(6)$
- $L=0\otimes S_{q_1q_2}=0 \Longrightarrow J_l=0$: $\Sigma^0_{c0}(\frac{1}{2}^+)$ $\Xi'^0_{c0}(\frac{1}{2}^+)$
- $L=1\otimes S_{q_1q_2}=0 \Longrightarrow J_l=1$: $\Sigma^1_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$ $\Xi'^1_{c1}(\frac{1}{2}^+,\frac{3}{2}^+)$
- $L=2\otimes S_{q_1q_2}=0 \Longrightarrow J_l=2$: $\Sigma^2_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$ $\Xi'^2_{c2}(\frac{3}{2}^+,\frac{5}{2}^+)$

(c) $l_\rho=1,\ l_\lambda=1$

Fig. 3 The notations for the D-wave charmed baryons.

Table 3 A summary of recently observed charmed baryons by Babar and Belle collaborations.

State	Mass and Width (MeV)	Decay channels in experiments	Other information
$\Lambda_c(2880)^+$	2881.5 ± 0.3, <8 [209]	$\Lambda_c\pi^+\pi^-$	J^P favors $\frac{5}{2}^+$ [204],
	$2881.9\pm0.1\pm0.5$, $5.8\pm1.5\pm1.1$ [203]	D^0p	$\frac{\Gamma(\Sigma_c^\star(2520)\pi^\pm)}{\Gamma(\Sigma_c(2455)\pi^\pm)}=0.225\pm0.062\pm0.025$ [204]
	$2881.2\pm0.2^{+0.4}_{-0.3}$, $5.5^{+0.7}_{-0.5}\pm0.4$ [204]	$\Sigma_c^{\star0,++}(2520)\pi^{+,-}$	
$\Lambda_c(2940)^+$	$2939.\pm1.3\pm1.0$, $17.5\pm5.2\pm5.9$ [203]	D^0p	–
	$2937.9\pm1.0^{+1.8}_{-0.4}$, $10\pm4\pm5$ [204]	$\Sigma_c(2455)^{0,++}\pi^{+,-}$	
$\Xi_c(2980)^+$	$2967.1\pm1.9\pm1.0$, $23.6\pm2.8\pm1.3$ [205]	$\Lambda_c^+K^-\pi^+$	–
	$2978.5\pm2.1\pm2.0$, $43.5\pm7.5\pm7.0$ [206]	$\Lambda_c^+K^-\pi^+$	
$\Xi_c(2980)^0$	$2977.1\pm8.8\pm3.5$, 43.5 [206]	$\Lambda_c^+K_S^0\pi^-$	–
$\Xi_c(3077)^+$	$3076.4\pm0.7\pm0.3$, $6.2\pm1.6\pm0.5$ [205]	$\Lambda_c^+K^-\pi^+$	–
	$3076.7\pm0.9\pm0.5$, $6.2\pm1.2\pm0.8$ [206]	$\Lambda_c^+K^-\pi^+$	
$\Xi_c(3077)^0$	$3082.8\pm1.8\pm1.5$, $5.2\pm3.1\pm1.8$ [206]	$\Lambda_c^+K_S^0\pi^-$	–
$\Xi_c(3055)^+$ [207]	$3054.2\pm1.2\pm0.5$, $17\pm6\pm11$	$\Lambda_c^+K^-\pi^+$	–
$\Xi_c(3122)^+$ [207]	$3122.9\pm1.3\pm0.3$, $4.4\pm3.4\pm1.7$	$\Lambda_c^+K^-\pi^+$	–
$\Omega_c(2768)^0$	2768.3 ± 3.0 [208]	$\Omega_c^0\gamma$	$J^P=\frac{3}{2}^+$

pletely be treated as a spectator and the binding effects may also be important. In the heavy quark language, the $\frac{1}{m_c^n}$ $(n\geqslant1)$ corrections may play important roles when the total width is evaluated. The question is, for exclusive decays, can one use the pQCD or a similar method where the $1/m_Q$ expansion is adopted? It seems that one can apply the pQCD method to analyze the Λ_b decays [235−239], but so far, there is no work devoted to pQCD application in calculating Λ_c decays yet. The reason is that the energy scale of M_{Λ_c} is relatively low and the exchanged gluons in the processes are not hard enough to enable us to consider only the leading, at most the NLO, contributions. How to involve a great number of higher terms in the expansion may be a very challenging task for theorists.

A great sign of progress in the field is that the double-charmed baryon Ξ_{cc} was observed by SELEX [240−244] at the Fermi Lab. Since it contains two charm quarks which would decay via weak interaction independently, one may expect to gain some information about the binding effects. In analog with the method given by Bigi *et al.* [225−230], we evaluated [245] the lifetimes of Ξ_{cc}^+, Ξ_{cc}^{++} and Ω_{cc}^+ which contain different light flavors, to order of

$1/m_c^3$. The hadronic matrix elements were evaluated in terms of the simplest model, i.e. the harmonic oscillator model [61]. The theoretical predictions can be tested in the future experiments especially the LHCb. Indeed, we lay our hope on the LHCb data which may produce a large database on such doubly charmed baryons, and maybe also baryons containing two b-quarks Ξ_{bb}, or b and c quarks Ξ_{bc} etc. By studying them, we will gain valuable information on the non-perturbative QCD effects which bind the heavy quarks together with a light flavor. Moreover, one may expect to observe baryons which are composed of three heavy quarks (b and/or c). There has been some theoretical work in this field [246], but phenomenological studies on their production rates and decay modes require more serious works in order to investigate their structures and governing mechanisms in some details. Some details of baryons containing two heavy quarks have been reviewed by Kiselev and Likholded in an enlightening paper [247].

We have also proposed a special decay mode of Ξ_{cc} to investigate the new physics. Namely, we suggest to measure a direct decay of Ξ_{cc} into non-charm final states [248], which is definitely a signal for new physics beyond SM. Concretely, in the SM, the direct decay of Ξ_{cc} into non-charm final states are very restricted and the sequential decays, such as $\Xi_{cc} \to \Lambda_c + \bar{D} \to \Lambda + K + X$ would produce final states which involve at least more than two hadrons, so these can be easily distinguished from the expected decay modes with two non-charmed hadrons. We have tried to employ two models which have recently become more popular to the theorists, the Z' and unparticle models which can induce the flavor-changing neutral current at tree level, so we can realize the reaction $cc \to qq'$ where q and q' are light flavors. However, our numerical results show that the two models cannot cause sizable fractions for practical measurement, even though one can have a great luminosity at LHCb. Our conclusion is that if such decay modes are observed, one can claim that as a definite evidence of new physics, but it is neither the Z' model, nor the unparticle model. This would motivate us to explore other new physics models beyond the SM.

The charmed pentaquark is a hot topic and the QCD theory does not exclude the pentaquark, the question is do we need to take it more seriously and where should we search for them and how do we identify them from the regular baryons [249–252]? As we discussed in previous sections, many theorists favor tetraquarks as a plausible explanation for some newly observed X, Y and Z resonances, therefore why should we firmly reject the existence of pentaquarks? A full scan of the pentaquark is not intended in this review work, but we would like to analyze some aspects concerning charm physics. The first proposal on the existence of pentaquarks is based on the group theory analysis [253], where in the pentquark is composed of four u and d quarks and one anti-strange quark with the strangeness baryon number being +1. Then a claim of finding the charmed pentaquark was made where a $\bar{c}$ replaced $\bar{s}$. However, later, most major labs in the world reported negative results in the search for pentaquarks and the society of high energy physics tended to deny discovery of the pentaquark. We also made efforts to look for traces of the pentaquark via some processes such as photo-production and radiative decay of pentaquarks, as well as the hadronic decays involving pentaquarks [254, 255], and the conclusion is still that more accurate measurements are necessary. On an other aspect, we seriously consider that the pentaquark may mix with the other baryonic states [256].

It is interesting to investigate the quark-structure of pentaquarks which is favored by QCD. Besides assuming that all the four quarks and an anti-quark all mix together in a big hadronic bag, more favorable structures are proposed by several authors, for example, Jaffe and Wilczek [257, 258] suggested the diquark-diquark-antiquark structure whereas Karliner and Lipkin [259] favored the diqaurk-triquark configuration. The key point is why the pentaquark has not been observed in any colliders but was seen in fixed target experiments. Lipkin tried to explain the situation by considering the initial quark configuration in the beam particle or target. This should draw the serious attention of experimentalists and theorists and design new experiments to search for evidence of pentaquarks. Since in the $p\bar{p}$ collisions at LHC, the rich quark contents and high luminosity may be a good source to produce pentaquarks, and until then we can draw a definite conclusion if the pentaquark can exist as a real particle. In this line the charmed pentaquark may be more favorable because it contains a heavy flavor. The most interesting subject would be that if the pentaquark is strictly prohibited, there must be a symmetry to restrict it, it would require a new understanding on the mechanism beside the general theory of QCD.

11 Diquark

It is an extremely interesting topic, not only for charmed baryons, but since it is widely applied in studying physical processes where baryons are involved, the diquark is worth careful investigation. By the $SU(3)$ theory, two quarks can reside in a color-anti-triplet $\bar{3}$ to constitute a loose bound state and then combine with the rest quark to make the baryon a color singlet. The QCD-induced potential is proportional to the Casimir factor $\langle\psi|\lambda^a\lambda^a|\psi\rangle$ where λ^a is the $SU(3)$ generator and $|\psi\rangle$ is the state, thus in a color-anti-triplet (or a triplet for two anti-quarks) the two quarks attract each other because the Casimir factor is negative. It seems that the bound state is plausible, but for light quarks, their linear momenta in the diquark are rather large and may spread out in space, thus we cannot make a real physical subject in common sense. Therefore whether the diquark structure is only a mathematical description for a baryon or it can be a physical subject is still in dispute. That is how to understand that the subject "diquark" is an intriguing topic in theoretical physics.

However, if the baryon contains a heavy quark, it becomes simpler than is usually assumed; the heavy quark may sit near the geometric center of the hadron (almost at the center) and one can further postulate that the linear momentum of the heavy quark is zero. Then the two light quarks would compose a diquark even though its spatial spreading is still large. Therefore, when we apply the concept of the diquark in this case, we indeed do not treat it as a point-like particle, but a loosely bound state with a common momentum, and interact with the heavy quark as a whole object. Moreover, when some reaction takes place, it may also perform as a whole object, but to compensate for the effects of its inner motion, one needs to introduce a form factor(s) whose phenomenological form was given in Ref.anselmino1 as $F(Q^2)=\dfrac{Q_0^2}{Q_0^2+Q^2}$, where Q_0 is a parameter and is determined as $Q_0^2\sim 3.2$ GeV^2 by fitting data [260−263]. The form is obviously understandable. The form factors should be normalized to unity as $Q^2\to 0$, i.e. as one looks at the diquark from a far distance, the form factor becomes a unity, whereas as $Q^2\to\infty$, the inspector then penetrates into the diquark, so that he would see the individual quarks instead of the whole and the diquark picture no longer holds and mathematically, it is required to approach zero as $Q^2\to\infty$. The form factors should be calculated in terms of a more fundamental way, i.e. based on quantum field theory. However, the non-perturbative QCD effects can by no means be treated completely based on an underlying theory so far, instead we need to invoke phenomenological models. The Bethe-Salpeter equation obviously is a reasonable approach where the kernel depends on concrete models. The form factor obtained in terms of the BSE generally coincides with the picture described above. Guo, Thomas and Williams [264, 265] studied the $1/m_Q$ corrections for the B-S equations for Λ_Q, Ω_Q in the diquark picture. Then in the framework of the B-S equation, we calculate the form factors for various types of diquarks [266], and the form factors obtained based on the B-S theory are qualitatively consistent with that given in Refs. [260−263]. Applying the diquark picture, we calculate several processes where heavy baryons participate [267−269]. The results are somehow competent to be compared with data, and they indicate that the diquark picture is vigorous and robust, even though it still needs further study.

It may be worth mentioning a special situation, namely that the baryons containing two heavy quarks should fit in the diquark picture well. Falk *et al.* [270] considered that the two charm quarks in a doubly heavy baryon constitute a "perturbatively bound diquark" whose wavefunction at origin follows:

$$|R_{(cc)}(0)|^2\approx\frac{1}{8}|R_\psi(0)|^2\approx(0.41\ \mathrm{GeV})^3$$

and the factor 1/8 is due to the difference of the color structures of color anti-triplet and singlet. Georgi and Wise proposed a very special symmetry called the superflavor symmetry [39, 40]. Simulating the supersymmetry where the SM fermions have their supersymmetric scalar partners and the SM bosons also have their SUSY fermionic partners, in the superflavor symmetry, the heavy quark is written in a doublet with its superflavor partner (scalar or vector) of color triplet as

$$\varPsi_s=(h_v^+,\chi_v)^T$$

and

$$\varPhi_v=(h_v^+,A_v^{mu})^T$$

where h_v^+ is a heavy quark and obeys $\psi h_v^+=h_v^+$, χ_v and A_v^μ are the heavy color triplet scalar and vector, respectively, with constraint $v_\mu A_v^\mu=0$. With this symmetry, we can estimate the production rates of heavy baryons which include two heavy quarks, i.e. a heavy diquark. In this scenario, the heavy scalar or vector diquarks can be treated as the superflavor symmetric partners of the heavy quark and then one can compare the baryon processes with the corresponding meson cases where two heavy mesons are produced. For both cases of the baryon and meson, the light quark would play the same role. In the meson cases, at the leading order, there is only one Isgur-Wise function to manifest the non-perturbative QCD effects, thus by the superflavor symmetry, we would also apply the same function [41] to calculate the production rates of X_{cc}, X_{bc}, X_{bb}, etc. in e^+e^- collisions. Definitely the same procedure can be applied to calculate the heavy baryon production rates at LHC, even though the situation is a bit more complicated.

12 Production of charm

The production of charm, i.e. the production of charmed mesons, baryons and charmonia, is a large subject, and we are not going to cover this field in this short review, but indicate some tricky and challenging issues.

The subject of production of J/ψ, ψ' and Υ was reviewed by Lansberg [272] in some detail. In this incomplete review, we do not intend to cover the whole subject, but indicate that this field is still tricky and intriguing by two examples which are widely discussed in the physics community.

As a long known question, the inclusive J/ψ production in $e^+e^- \to J/\psi c\bar{c}$ raises a challenge to our theory. At LO, the QCD prediction on the cross section is smaller than the observed value by a factor of 5. This discrepancy may come from the wavefunctions of the produced J/ψ [271] or the NLO corrections. In an enlightening work [276], the authors suggested that the NLO may enhance the cross section by a factor of about 1.8. This may greatly change the situation.

However, another serious question is raised in that if the NLO can change this cross section so much, what mechanism or symmetry determines this unusual enhancement? Generally, the loop suppression is proportional to α_s/π, and at this energy scale α_s is smaller than 0.2. If the number of diagrams (at NLO, the number of diagrams is rather large) is the reason, as long as the interference among the diagrams are mostly constructive, one could ask, how about the contribution of NNLO? Definitely, even though the diagrams at NLO are constructive, it does not suggest that they are constructive at NNLO at all. If the NNLO contribution decreases, we can happily declare that the theory works well, however if it continues to increase or the correction at NNLO is even larger than that of NLO, the unitarity would confront a serious challenge and one needs to take the hint more seriously. Indeed, to give a firm answer to the question, one must calculate the NNLO corrections. However, there are over 1000 two-loop diagrams and a complete calculation is an extremely difficult project. Since it is necessary, some serious work must be done and it is indeed worthwhile.

The problem about the production of J/ψ was investigated by several groups [273−275]. The authors studied the hadronprocution of J/ψ, ψ', Υ in association with the heavy quark pair which was the subject we listed above and investigated further, with a possible solution to the J/ψ production puzzle. Their result was confirmed by Gong and Wang [277, 278]. In the papers [277, 278], the authors further studied the p_t distribution of the J/ψ polarization in the $pp \to J/\psi + X$ process. There is also serious discrepancy between the theoretical prediction and data. They calculated the NLO p_t distribution of the J/ψ polarization and found an amazing change. Namely, they found that the polarization status of J/ψ at large p_t changes from the transverse-polarization dominance at LO into the longitudinal polarization dominance at NLO. This correction indeed partly makes up for the gap between the theoretical result and the data. In their first work, they only considered the color-singlet contribution and got improvement, so there was a hope that when the color-octet contribution was involved, the situation would be improved further and the data could be explained eventually. But unfortunately, when the color-octet contribution is taken into account, the new correction is not so drastic and does not change the whole picture much, therefore the large gap between theoretical prediction and data still exists. The authors [277, 278] claimed that to explain the data, new mechanisms or new physics or new understanding of the J/ψ structure etc. are needed.

The two examples do not intend to cover the large-scale field of production of charm, but indicate that there are still many problems which are not fully understood yet, and a lot of theoretical efforts are required. This is also associated with the puzzle we discussed in this review as to why the pentaquark has not been observed, and so on.

13 Discussion

This are by no means a definite conclusion because in this review, we raise many unsolved problems and list some references where many authors devote great efforts in this fascinating field. The theoretical difficulty is obvious in that charm the quark is heavy, but not sufficiently heavy, so non-perturbative QCD effects still play important roles, and moreover, the relativistic effects of quark-quark in hadrons are also not negligible. Therefore, the heavy quark effective theory (HQET) may apply to this field, but needs to consider higher orders in the $1/m_c$ expansions, which it brings unexpected uncertainties. The good example is the lifetime difference between Λ_c and D mesons, compared to the B-cases. On an other aspect, careful studies in the charm field may greatly enrich our understanding of the basic principles; concretely, how to properly deal with the non-perturbative QCD effects is one of the most important issues. Moreover, the work by the authors [276−278] indicates that higher order QCD corrections may be very important, and even play a crucial role. This indeed further intrigues the field of charm physics.

The main task in this field seems not to be closely related to searching for new physics beyond SM, as is generally considered because of its energy scale, but there is indeed the possibility to investigate new physics. For example, the D^0-$\bar{D}^0$ mixing, as aforementioned, and the lepton flavor violation processes besides the neutrino oscillation experiments. Namely, one can examine the reactions like [279]

$$e^+e^- \to \psi \to e^{-(+)}\mu^{+(-)}$$

for which the charm-tau energy region may be the ideal observation place. On the theoretical aspect, in the SM, this FCNC can occur at one-loop order where neutrino masses need to be large enough to produce observational effects. However, it contradicts all the data of the solar neutrino and atmospheric neutrino experiments, and one can claim that observation of such a reaction is a clear signal of new physics beyond SM. There are so many new physics models which can produce this reaction, can the reaction, if it indeed occurs, distinguish among them? Generally it cannot, but one can at least tell which model is a possible one and maybe, such a reaction would help to get rid of a few other proposed models.

There are general discussions on the possibilities of measuring the production rates from the theoretical angle [280] and set bounds on the model parameters. Very recently, we suggested [281] that the unparticle model may induce such FCNC. By adopting the reasonable region of the concerned parameters, our calculation shows that one can expect a large production rate near the resonance peak of J/ψ. Carefully analyzing possible backgrounds, one has confidence to extract a clear signal for the luminosity of PEPC II and detection accuracy of BESIII. If it is observed, the unparticle model at least is one of the favorable candidates of new physics. On the other hand, if it is not observed, the unparticle model would be somehow disfavored.

There are too many questions to be answered in this energy region, and as we discussed above, there are several very enlightening review papers which shed light on this exciting field [12, 282−286].

Indeed, besides the major line, there are some other subjects which should be researched in this field, for example, it is proposed to test the Bell inequality in high energy processes, $\eta_c \to \Lambda\bar{\Lambda} \to p\pi^- + \bar{p}\pi^+$ [287] and decays of charmonia into kaon [288−291]. As a matter of fact, the scheme proposed in Ref. [287] was carried out in terms of the DM2 data, however, because the database was too small, poor statistics made the work meaningless. Today, the situation will be different, the BESIII will provide a tremendously large database of J/ψ, so that one may expect to make accurate measurements to validate the Bell inequality.

We believe that we are standing at a prosperous epoch for high energy physics. The two B-factories continue to produce more data where one may also have a chance to study charm-related physics, such as CP violation and higher excited states of charmonia and charmed mesons and baryons. Moreover, the LHC will be running soon and not only a lot of information towards underlying physics such as the Higgs mechanism and new physics beyond SM will be collected, but also some details about bottom and charm physics can be achieved, especially, we may lay great hope on the LHCb. Even the ALICE can tell us some interesting subjects about charm physics, for example, if the J/ψ suppression in quark-gluon-plasma (QGP) really exists, can the higher temperature and pressure change the potential between quarks, can we expect an observable phase transition, like deconfinement and chiral symmetry restoration, etc. More than anything else which can enrich our knowledge on charm physics, the BEPC II and BESIII will be operating very soon, and a great database on J/ψ and other members of the ψ family, as well as on the D-mesons, baryons, and even on tensors will be available, so that we are sure, many new challenges will be waiting for us. So we feel much encouraged than at any other time, because so many questions need to be solved and more data will help to do the job.

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