New constraints
on NRQCD long-distance matrix elements
from $J/\psi$ plus $W/Z$ production at the CERN LHC

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Abstract

We study the associated production of prompt $J/\psi$ mesons and $W$ or $Z$ bosons within the factorization approach of nonrelativistic QCD (NRQCD) at next-to-leading order in $\alpha_s$, via intermediate color singlet $^3S_1^{[1]}$ and $^3P_J^{[1]}$ and color octet $^1S_0^{[8]}$, $^3S_1^{[8]}$ and $^3P_J^{[8]}$ states. Requiring for our predictions to be compatible with recent ATLAS measurements yields stringent new constraints on charmonium long-distance matrix elements (LDMEs) being nonperturbative, process-independent input parameters. Considering four popular LDME sets fitted to data of single $J/\psi$ inclusive production, we find that one is marginally compatible with the data, with central predictions typically falling short by a factor of three, one is unfavored, the factor of shortfall being about one order of magnitude, and two violate cross section positivity for direct $J/\psi + W/Z$ production. The large rate of prompt $J/\psi$ plus $W$ production observed by ATLAS provides strong evidence for the color octet mechanism inherent to NRQCD factorization, the leading color singlet contribution entering only at $\mathcal{O}(G_F\alpha_s^4)$, beyond the order considered here.
Although heavy quarkonia have been discovered already in 1974, the underlying mechanisms governing their production in high-energy collisions are still not fully understood. The most prominent approach is via the factorization theorem of nonrelativistic QCD (NRQCD) \cite{1,2}. According to it, the production cross section of quarkonium $H$ factorizes into perturbative short-distance cross sections of heavy quark-antiquark bound-state production and supposedly universal nonperturbative long-distance matrix elements (LDMEs) $\langle O^H(n) \rangle$, where $n = 2S+1L^J_s^{[1,8]}$ denotes the quarkonic Fock state, in color singlet “[1]” or octet “[8]” configuration. Velocity ($v$) scaling rules \cite{8} impose a strong hierarchy on the $\langle O^H(n) \rangle$ values, leading to a double expansion in the strong-coupling constant $\alpha_s$ and $v$.

For $H = J/\psi$ and $\psi(2S)$, the LDMEs of $n = 3S_1^{[1]}$ are leading in $v$ and those of $n = 1S_0^{[8]}$, $3S_1^{[8]}$ and $3P_2^{[8]}$ are subleading. For $H = \chi_{cJ}$, the LDMEs of $n = 3P_1^{[1]}$ and $3S_1^{[8]}$ are both leading in $v$.

The available charmonium LDME sets have all been extracted from data of inclusive single production. Thanks to the high luminosity meanwhile achieved by the LHC, also double production and associated production with bottomonia, $W$, and $Z$ bosons have been studied there, which can inject orthogonal information into LDME determinations. The goal of this letter is to provide the first complete analysis of prompt-$J/\psi$ plus $W$ or $Z$ hadroproduction at next-to-leading order (NLO) in $\alpha_s$, i.e. through $O(G_F\alpha_s^3)$. Invoking QCD and NRQCD factorization, we calculate the cross sections as

$$\sigma(pp \to J/\psi + W/Z + X) = \sum_H \text{Br}(H \to J/\psi) \sum_n \bar{\sigma}(pp \to c\bar{c}[n] + W/Z + X) \langle O^H(n) \rangle,$$  

(1)

$$\bar{\sigma}(pp \to c\bar{c}[n] + W/Z + X) = \sum_{a,b} \int dx_a dx_b f_{a/p}(x_a) f_{b/p}(x_b) \bar{\sigma}(ab \to c\bar{c}[n] + W/Z + X),$$  

(2)

where $H = J/\psi$, $\psi(2S)$ and $\chi_{cJ}$, with $J = 0, 1, 2$, and $n$ runs over all Fock states specified above. $\bar{\sigma}(ab \to c\bar{c}[n] + W/Z + X)$ are the partonic cross sections, evaluated as perturbative expansions in $\alpha_s$; $f_{a/p}(x)$ is the parton density function (PDF) of parton $a$ in the proton; $a, b$ include the up, down, strange (anti)quarks, and the gluon; $\text{Br}(H \to J/\psi)$ are the decay branching fractions, including $\text{Br}(J/\psi \to J/\psi) = 1$ for ease of notation. In the $W$ case, where $W^\pm$ is summed over, only $n = 3S_1^{[8]}$ contributes at leading order (LO) in $\alpha_s$.

Partial results may be found in the literature. The LO results have already been obtained two decades ago \cite{4}. At NLO, the $1S_0^{[8]}$, $3S_1^{[1,8]}$, and $3P_2^{[8]}$ channels have been considered in the $W$ case \cite{5}, and the $3S_1^{[1]}$, $3S_1^{[8]}$, and $3P_2^{[8]}$ channels in the $Z$ case. We can reproduce these results, with noticeable differences only in Fig. 4 of Ref. \cite{5}, and fill all gaps by providing the NLO results for the $3P_1^{[1]}$ channels in the $W$ case and the $1S_0^{[8]}$, $3P_2^{[8]}$, and $3P_2^{[1]}$ channels in the $Z$ case. Unlike the preceding works, we have to consider virtual corrections to $P$-wave state production. Notice that the NLO $c\bar{c}[3P_2^{[8]}] + W$ contribution is just tree level. Albeit $P$-wave state virtual corrections have been tackled for single inclusive production, the additional $W/Z$ mass scale elevates the complexity of this NLO NRQCD
calculation to an unprecedented level.

Let us now review the main technical aspects of our calculation, starting with the treatment of $\gamma_5$, which appears in the $W$ and $Z$ axial-vector couplings and in the spin projection onto the $^1S_0^0$ state. Adopting the standard scheme \cite{8,9,10,11}, we use the axial-vector coupling $\gamma_\mu \gamma_5$ in its antisymmetric form 

$$\frac{i}{2} (\gamma_\mu \gamma_5 - \gamma_5 \gamma_\mu),$$

and employ the relation

$$\epsilon_{\mu_1 \mu_2 \mu_3 \mu_4} \epsilon_{\nu_1 \nu_2 \nu_3 \nu_4} = - \det (g_{\mu_\nu}),$$  \hspace{1cm} (3)

and apply the finite axial-vector coupling renormalization (see, e.g., Ref. \cite{11}). We explicitly verified that the final results are then independent on whether we choose a $D$- or four-dimensional metric $g$ in Eq. (3).

We generate, treat and square the amplitudes using FeynArts \cite{12} and custom FORM \cite{13} and Mathematica codes. We reduce the virtual loop integrals to a common set of master integrals using two methods. In the first one, we directly apply integration-by-parts relations generated with AIR \cite{14}, while in the second one, we first invoke a custom Passarino-Veltman-type tensor reduction, generalized for the case of arbitrary propagator powers and linearly dependent propagator momenta. We analytically check the agreement of both methods. As for the master integrals, we implement our own analytic expressions in combination with QCDLoop \cite{15}, checking everything against OneLoop \cite{16}. We analytically check the ultraviolet and infrared finiteness of our results and numerically compare our real corrections, after imposing infrared cutoffs, against HELACOnia \cite{18} output.

We organize the phase space integrations using the dipole subtraction procedure outlined in Ref. \cite{19}, changing only the momentum mapping of dipole term $V_{3,j}$ (into MapPW6($p_j,p_2$)) to cope with the presence of the massive non-QCD particle in the final state. We numerically check that all dipoles reproduce the real corrections in their respective limits and perform the check on the integrated dipoles outlined in section 4.3 of Ref. \cite{20}. As a further check, we also implement the phase space slicing procedure along section 3 of Ref. \cite{20} to find numerical agreement. We recover the notion \cite{20} that dipole subtraction significantly outperforms phase space slicing as for precision and speed.

We renormalize the charm quark mass in the on-shell scheme to be $m_c = 1.5$ GeV and take the charmonia to have mass $2m_c$ for definiteness. We express all electroweak couplings in terms of Fermi’s constant $G_F$ and the on-shell $W$ and $Z$ boson masses $M_W$ and $M_Z$.

At LO (NLO), we use the CTEQ6L1 (CTEQ6M) proton PDFs \cite{22} with asymptotic scale parameter $\Lambda_{QCD}^{(4)} = 215$ MeV (326 MeV) for $n_f = 4$ quark flavors, to be used in the one-loop (two-loop) formula for $\alpha_s^{(n_f)}(\mu_r)$, with renormalization scale $\mu_r$.

Besides $\mu_r$, two more unphysical scales appear, namely, the factorization scales of QCD and NRQCD, $\mu_f$ and $\mu_A$. For definiteness, we put $\mu_A = m_c$ as default value and unify $\mu = \mu_r = \mu_f$, for which plausible default choices include $\mu_0 = m_{T,J/\psi}$ \cite{5,6}, $\mu_0 = \sqrt{m_{T,J/\psi} m_{T,W/Z}}$
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Table 1: LDME sets used in our analysis. Due to heavy-quark spin symmetry, \( \langle \mathcal{O}_{J/\psi}^{H}(3P_{J}^{[8]}) \rangle = (2J + 1)\langle \mathcal{O}_{J/\psi}^{H}(3P_{0}^{[8]}) \rangle \), \( \langle \mathcal{O}_{\chi}^{(^{3}S_{1}^{[8]})} \rangle = (2J + 1)\langle \mathcal{O}_{\chi}^{(^{3}S_{1}^{[8]})} \rangle \), and \( \langle \mathcal{O}_{\chi}^{J}(^{3}S_{1}^{[8]}) \rangle = (2J + 1)\langle \mathcal{O}_{\chi}^{(^{3}S_{1}^{[8]})} \rangle \).

| LDME Set | \( \langle \mathcal{O}_{J/\psi}^{H}(3P_{J}^{[8]}) \rangle / \text{GeV}^{3} \) | \( \langle \mathcal{O}_{\chi}^{(^{3}S_{1}^{[8]})} \rangle / \text{GeV}^{3} \) | \( \langle \mathcal{O}_{\chi}^{J}(^{3}S_{1}^{[8]}) \rangle / \text{GeV}^{3} \) | \( \langle \mathcal{O}_{\chi}^{J}(^{3}S_{1}^{[8]}) \rangle / \text{GeV}^{3} \) |
|-----------|----------------|----------------|----------------|----------------|
| \( \langle \mathcal{O}_{J/\psi}^{H}(3P_{J}^{[8]}) \rangle / \text{GeV}^{3} \) | 0.0304 ± 0.0035 | 0.017 ± 0.0005 | -0.0091 ± 0.0016 | -0.0091 ± 0.0016 |
| \( \langle \mathcal{O}_{\chi}^{(^{3}S_{1}^{[8]})} \rangle / \text{GeV}^{3} \) | 0.0100 ± 0.0003 | 0.00054 ± 0.00003 | 0.0049 ± 0.0001 | 0.0022 ± 0.0005 |
| \( \langle \mathcal{O}_{\chi}^{J}(^{3}S_{1}^{[8]}) \rangle / \text{GeV}^{3} \) | 0.017 | 0.017 | 0.017 | 0.017 |
| \( \langle \mathcal{O}_{\chi}^{J}(^{3}S_{1}^{[8]}) \rangle / \text{GeV}^{3} \) | 0.017 | 0.017 | 0.017 | 0.017 |

has been estimated to be 3% using Monte Carlo simulations in Ref. [28]. On the other hand, thanks to the much tighter \( m_{t^{+}t^{-}} \) cut, of just \( \pm 10 \) GeV around the \( Z \) peak, the measurement of Ref. [23] should hardly be contaminated by \( \gamma^{*} \) events. To correct for this mismatch, we subtract 3% from the result for \( \sigma_{Z} \times \text{Br}(Z \rightarrow l^{+}l^{-}) \) in Ref. [28]. To summarize, we have \( \sigma_{W} = (98.71 \pm 2.34) \) nb at \( \sqrt{s} = 7 \) TeV, and \( \sigma_{W} = (112.43 \pm 3.80) \) nb and \( \sigma_{Z} = (33.28 \pm 1.19) \) nb at \( \sqrt{s} = 8 \) TeV.

We employ four popular LDME sets, which we list in Table 1. The first one is a combination of (i) the \( J/\psi \) LDMEs obtained by a global fit to prompt production data after subtracting the estimated feed-down contributions [24]; (ii) the \( \psi' \) LDMEs recently determined from a global fit to data of unpolarized hadroproduction [30]; and (iii) the \( \chi_{cJ} \) LDMEs determined in Ref. [31]. The other three LDME sets are taken from Refs. [32][33][34], where the cut \( p_{T} > 9 \) GeV was imposed in the last case.

To enable interested readers to perform comparisons with alternative LDME sets, we list in Table 1 the LO and NLO default cross sections \( d\sigma(pp \rightarrow c\bar{c}[n] + W/Z + X)/dp_{T,J/\psi} \times \text{Br}(J/\psi \rightarrow \mu^{+}\mu^{-}) \) of Eq. (2) assuming the ATLAS kinematic setup at \( \sqrt{s} = 8 \) TeV [23][24] including the binning in \( p_{T,J/\psi} \).

Figure 1 and Table 2 also usefully portray the anatomy of the NLO corrections in the various \( n \) channels as for sign and magnitude. For direct \( J/\psi \) production at NLO in the color singlet model (CSM), via \( c\bar{c}[^{3}S_{1}^{[1]}] + Z \), the perturbative expansion in \( \alpha_{s} \) appears to be well behaved, with a \( K \) factor of order unity. The \( K \) factors of \( \tilde{\sigma}(pp \rightarrow c\bar{c}[^{3}S_{0}^{[8]}] + Z + X) \) and \( \tilde{\sigma}(pp \rightarrow c\bar{c}[^{3}S_{1}^{[8]}] + W + X) \) range between 2 and 3. The \( c\bar{c}[^{3}S_{0}^{[8]}] \) contribution, which only appears at NLO, is positive and small against the LO \( c\bar{c}[^{3}S_{1}^{[8]}] \) contribution. The NLO \( P \)-wave contributions, both color singlet and octet, are throughout negative. The NLO NRQCD predictions are likely to be more reliable in the \( Z \) case than in the \( W \) case, where we expect large next-to-next-to-leading-order contributions due to the delayed...
\[ \sigma_{p p \rightarrow c[n] + Z + X} \text{ for all contributing Fock states } n \text{ in fb/GeV}^4 \text{ (fb/GeV}^6\text{) for } S \text{ (P) wave states, assuming the ATLAS kinematic conditions at } \sqrt{s} = 8 \text{ TeV} \]

| \( p_{T,J/\psi} \text{ [GeV]} \) | \( pp \rightarrow c[n] + Z + X \) | \( pp \rightarrow c[n] + W + X \) |
|-----------------|-----------------|-----------------|
| \( n = 3S_{0}^{[8]} \), LO | 0.0862 | 0.00472 |
| \( n = 3S_{0}^{[8]} \), NLO | 0.0806 | 0.00339 |
| \( n = 1S_{0}^{[8]} \), LO | 3.07 | 0.00220 |
| \( n = 1S_{0}^{[8]} \), NLO | 5.88 | 0.00107 |
| \( n = 3S_{1}^{[8]} \), LO | 127 | 365 |
| \( n = 3S_{1}^{[8]} \), NLO | 83.7 | 236 |
| \( n = 3P_{1}^{[8]} \), LO | 2.86 | 0.323 |
| \( n = 3P_{1}^{[8]} \), NLO | 2.86 | 0.206 |
| \( n = 3P_{1}^{[8]} \), NLO | 0.323 | 0.106 |
| \( n = 3P_{1}^{[8]} \), NLO | 0.0362 | 0.0948 |
| \( n = 3P_{1}^{[8]} \), NLO | 0.00102 | 0.432 |
| \( n = 3P_{2}^{[8]} \), LO | 0.751 | 0.0751 |
| \( n = 3P_{2}^{[8]} \), NLO | 0.699 | 0.317 |

Table 2: LO and NLO cross sections \( d\sigma(pp \rightarrow c[n] + W/Z + X)/dp_{T,J/\psi} \times Br(J/\psi \rightarrow \mu^+\mu^-) \) of Eq. (2) for all contributing Fock states \( n \) in fb/GeV\(^4\) (fb/GeV\(^6\)) for \( S \) (\( P \)) wave states, assuming the ATLAS kinematic conditions at \( \sqrt{s} = 8 \text{ TeV} \) including the binning in \( p_{T,J/\psi} \). The common shorthand notation \( d\sigma(pp \rightarrow c[n] + W/Z + X) \) implies \( \sum_{J=0}^{2}(2J+1)d\sigma(pp \rightarrow c[n] + W/Z + X) \). The integration accuracy is around 1%.
Figure 2: Comparison of the ATLAS data from Refs. [23,24,25] (rows), adjusted as described in the text, to our NLO predictions for $d\sigma(pp \to J/\psi + W/Z + X)/dp_{T,J/\psi} \times \text{Br}(J/\psi \to \mu^+\mu^-)$ in fb/GeV evaluated successively with the LDME sets of Table 1 (columns). The theoretical-uncertainty bands are evaluated as described in the text.

unfolding of the $n$ structure, with only $n = 3S_0^{[3]}$ being present at LO and $n = 3S_0^{[1]}$ not even at NLO. As for the prompt-$J/\psi$ plus $W/Z$ cross sections to be compared with ATLAS data in Fig. 2, we anticipate that their $K$ factors at the bin level range between 0.9 and 1.7 in the CSM and between 1.7 and 4.8 in full NRQCD.

In Fig. 2, we compare the ATLAS data [23,24,25], modified as explained above, to our NLO predictions for $d\sigma(pp \to J/\psi + W/Z + X)/dp_{T,J/\psi} \times \text{Br}(J/\psi \to \mu^+\mu^-)$ with the same binning in $p_{T,J/\psi}$. The three rows in Fig. 2 correspond to $J/\psi + Z$ production at $\sqrt{s} = 8$ TeV [23] and $J/\psi + W$ production at 8 TeV [24] and 7 TeV [25], the four columns to the LDME sets in Table 1. In each frame, we break down the total result into the contributions from the individual channels $n$ of direct production and the combined feed-down contribution, and indicate theoretical uncertainties in the CSM and NRQCD results. The theoretical uncertainties are evaluated by adding in quadrature the errors from the following three sources: (i) variation of $\mu$ by a factor of 4 up and down relative to $\mu_0 = \sqrt{m_{T,J/\psi}m_{T,W/Z}}$; (ii) variation of $\mu_A$ by a factor of 2 up and down relative to $m_c$;
(iii) quadratic combination of the individual LDME errors quoted in Table 1. The large μ variation is to at least partially account for the fact that also μ0 = M_{W/Z} is a plausible reference scale. In want of specific information on the correlations between these errors, the theoretical uncertainties thus evaluated are likely to be conservative.

We are now in a position to assess the LDME sets in Table 1 with regard to their ability to usefully describe the ATLAS data [23,24,25] at NLO in NRQCD. We immediately observe that the second [32] and third [33] LDME sets lead to negative direct J/ψ + W/Z production cross sections, which is physically unacceptable, excluding them from the valid LDME sets. They are only rescued into the positive by the feed-down contributions. Next, we observe that the first three LDME sets considered (plus the ones of Refs. [35,36,37], for which we refrain from showing results for lack of space) lead to predictions that throughout undershoot the data by about one order of magnitude. To attribute such a sizable gap to underestimated DPS contributions would require the cross sections to be overwhelmingly dominated by DPS, in contrast to the J/ψ–W/Z azimuthal-angle analyses of Refs. [23,24,25], which all support SPS dominance. This renders the first LDME set unfavorable, albeit not invalid. On the other hand, the fourth LDME set [34] leads to an underestimation of the data by only a factor of about three, with experimental and theoretical uncertainties typically touching or overlapping. We note in passing that this and the other LDME set determined in Ref. [34] have, however, their own problems in applications beyond the scope of this paper, including negative NLO predictions for the LHCb measurement of prompt η_c production [38] and underestimation of HERA photoproduction data by one order of magnitude. Furthermore, they involve a delicate fine tuning of negative 1S_0[8] and positive 3S_1[8] J/ψ hadroproduction channels canceling to around 90%.

To summarize, we have presented the first complete NLO NRQCD predictions of prompt-J/ψ plus W/Z associated hadroproduction, tackling P-wave loop contributions with an additional large mass scale. Requiring consistency with ATLAS data [23,24,25] provides valuable new information on the interplay of the J/ψ, χ_{cJ}, and ψ′ LDMEs, orthogonal to the one encoded in one-particle-inclusive charmonium production data previously fitted to [29,30,31,32,33,34,35,36,37], which has allowed us to critically assess — and in some cases even to rule out — the resulting LDME sets. The decent description of the ATLAS measurement of prompt-J/ψ plus W production [24,25] by at least the LDME set [34] provides strong evidence for the color octet mechanism and, once again, exposes the deficiency of the CSM to describe charmonium production. Our analysis thus marks an important milestone on the path of scrutinizing NRQCD factorization.

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References

[1] W. E. Caswell and G. P. Lepage, Phys. Lett. B 167, 437 (1986).
[2] G. T. Bodwin, E. Braaten, and G. P. Lepage, Phys. Rev. D 51, 1125 (1995); [erratum: Phys. Rev. D 55, 5853 (1997)].

[3] G. P. Lepage, L. Magnea, C. Nakhleh, U. Magnea and K. Hornbostel, Phys. Rev. D 46, 4052-4067 (1992).

[4] B. A. Kniehl, C. P. Palisoc and L. Zwirner, Phys. Rev. D 66, 114002 (2002).

[5] G. Li, M. Song, R. Y. Zhang and W. G. Ma, Phys. Rev. D 83, 014001 (2011).

[6] M. Song, W. G. Ma, G. Li, R. Y. Zhang and L. Guo, JHEP 02, 071 (2011) [erratum: JHEP 12, 010 (2012)].

[7] B. Gong, J. P. Lansberg, C. Lorce and J. Wang, JHEP 03, 115 (2013).

[8] G. 't Hooft and M. J. G. Veltman, Nucl. Phys. B 44, 189-213 (1972).

[9] P. Breitenlohner and D. Maison, Commun. Math. Phys. 52, 11-38 (1977).

[10] S. A. Larin, Phys. Lett. B 303, 113-118 (1993).

[11] J. C. Collins, “Renormalization: An Introduction to Renormalization, The Renormalization Group, and the Operator Product Expansion,”, Cambridge University Press (1984).

[12] T. Hahn, Comput. Phys. Commun. 140, 418-431 (2001).

[13] J. A. M. Vermaseren, arXiv:math-ph/0010025.

[14] C. Anastasiou and A. Lazopoulos, JHEP 07, 046 (2004).

[15] R. K. Ellis and G. Zanderighi, JHEP 02, 002 (2008).

[16] A. van Hameren, Comput. Phys. Commun. 182, 2427-2438 (2011).

[17] G. P. Lepage, J. Comput. Phys. 27, 192 (1978).

[18] H. S. Shao, Comput. Phys. Commun. 198, 238-259 (2016).

[19] M. Butenschoen and B. A. Kniehl, Nucl. Phys. B 950, 114843 (2020).

[20] M. Butenschoen and B. A. Kniehl, Nucl. Phys. B 957, 115056 (2020).

[21] P. A. Zyla et al. [Particle Data Group], Prog. Theor. Exp. Phys. 2020, 083C01 (2020).

[22] J. Pumplin et al. (CTEQ Collaboration), JHEP 0207, 012 (2002).

[23] G. Aad et al. [ATLAS], Eur. Phys. J. C 75, 229 (2015).

[24] M. Aaboud et al. [ATLAS], JHEP 01, 095 (2020).
[25] G. Aad et al. [ATLAS], JHEP 04, 172 (2014).
[26] G. Aad et al. [ATLAS], New J. Phys. 15, 033038 (2013).
[27] M. Aaboud et al. [ATLAS], Eur. Phys. J. C 77, 367 (2017).
[28] S. Chatrchyan et al. [CMS], Phys. Rev. Lett. 112, 191802 (2014).
[29] M. Butenschoen and B. A. Kniehl, Phys. Rev. D 84, R051501 (2011).
[30] M. Butenschoen and B. A. Kniehl, DESY-22-119
[31] Y. Q. Ma, K. Wang and K. T. Chao, Phys. Rev. D 83, 111503 (2011).
[32] B. Gong, L. -P. Wan, J. -X. Wang and H. -F. Zhang, Phys. Rev. Lett. 110, 042002 (2013).
[33] G. T. Bodwin, K. T. Chao, H. S. Chung, U. R. Kim, J. Lee and Y. Q. Ma, Phys. Rev. D 93, no.3, 034041 (2016).
[34] N. Brambilla, H. S. Chung, A. Vairo and X. P. Wang, Phys. Rev. D 105, no.11, L111503 (2022).
[35] K. -T. Chao, Y. -Q. Ma, H. -S. Shao, K. Wang and Y. -J. Zhang, Phys. Rev. Lett. 108, 242004 (2012).
[36] H. Han, Y. Q. Ma, C. Meng, H. S. Shao and K. T. Chao, Phys. Rev. Lett. 114, no.9, 092005 (2015).
[37] H. F. Zhang, Z. Sun, W. L. Sang and R. Li, Phys. Rev. Lett. 114, no.9, 092006 (2015).
[38] R. Aaij et al. [LHCb], Eur. Phys. J. C 75, no.7, 311 (2015).