Dissociation of the J/ψ by light mesons, and chiral symmetry

Alex Bourque and Charles Gale
Department of Physics, McGill University, 3600 University Street, Montreal, Quebec, Canada H3A 2T8
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The implication of chiral symmetry for the pion-induced dissociation of the J/ψ is examined in detail. It is shown how the low-energy dynamics of pions, constrained by chiral symmetry, affect the dissociation cross section. The derived soft-pion theorem is then integrated into a Lagrangian model that includes also abnormal parity content and chiral-symmetric form factors. Dissociation by the ρ meson is also considered.

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I. INTRODUCTION

It is predicted that at very high energy densities, confined hadronic matter melts into a novel form: the quark-gluon plasma (QGP). Several signatures to characterize its properties within the context of heavy-ion collisions have been proposed. One of these, initially championed by Matsui and Satz [1], is charmonium suppression. It rests on the observation that correlated $\bar{c}c$ pairs created in the earliest stage of the collisions through hard scatterings probe all subsequent stages of the system evolution. In particular, if a QGP is formed, they argued that the observed yield should be suppressed because of color screening [1]. The current view includes not only suppression but also the regeneration of charmonium [2–4]. Moreover, recent lattice data suggest that the J/ψ may survive in the plasma well above T_C [5,6], implying that that there could be no direct QGP suppression of this meson [7]; see, however, Ref. [8].

Before a claim of any definite QGP effects is made, it is essential to verify that the results cannot be reproduced by more mundane nuclear effects. Of all possible mechanisms, charmonium dissociation by nucleons is probably the most important one. Indeed, it is seen to be sufficient to explain the suppression patterns observed at the SPS not only for $p+A$ systems but also $O+U$ and $S+U$ collisions [9]. For heavier systems, nuclear suppression is not sufficient to account for experimental observations. For example, in Pb+Pb collisions at SPS, an abnormal suppression is observed. One possible cause of the charmonium suppression could of course be screening [10]. But dissociation by light-meson comovers can also go a long way in explaining the observed data [11–21].

In most phenomenological studies, the dissociation cross section by comovers is a model parameter, and little is said about the underlying microscopic mechanisms. Because experimental information about dissociation processes is scarce, one has to rely on theoretical studies. Several approaches are possible. One model calculates the dissociation cross sections by using constituent quarks and a nonrelativistic potential [21–24]. The dissociation processes then arise through the exchange of quarks. A fully relativistic constituent quark model can also be constructed based on an extension of the Nambu-Jona Lasino (NJL) model to the charm sector [25–30]. Dissociation then occurs through quark and meson exchanges. Being nonrenormalizable, the model requires the specification of an ultraviolet loop cutoff. One can circumvent the need of such a cutoff by introducing form factors at the quark level. This leads to the extended nonlocal NJL model of Ref. [31]. Another model relies rather on extrapolations of QCD sum rule (QCDSR) results to extract momentum-dependent vertices [32–39]. Finally, phenomenological Lagrangians [40–45] can be utilized. There, to account for the composite nature of the mesons, *ad hoc* form factors are often introduced.

These models then produce cross sections ranging from submillibarn to a few millibarns. Moreover, their energy behavior can be quite different [46]. This is compounded by the fact that in many models, chiral symmetry is not clearly implemented. As pointed out in Ref. [47] in the context of the Lagrangian-type models, chiral symmetry implies that for the normal parity content of the process $J/\psi+\pi\rightarrow(D^*+\bar{D})+(\bar{D}^*+D)$ the pion should decouple in the soft-momentum limit leading to a vanishing amplitude. Because this process is considered to be dominant, owing to the abundance of pions, quantifying this effect is therefore important.

In Ref. [48], the effect of implementing chiral symmetry in a simple Lagrangian model without form factors was considered in contrast with previous phenomenological Lagrangians [40–45]. It was shown there that for the $J/\psi+\pi\rightarrow(D^*+\bar{D})+(\bar{D}^*+D)$ process a reduction at threshold did occur. It is the purpose of this article to propose an improved Lagrangian model that incorporates not only chiral symmetry and form factors but also other dissociation channels, i.e., the so-called abnormal parity processes [45]. The ρ-induced dissociation cross sections will also be evaluated to assess the relative importance of dissociation by other light resonances.

This article is organized as follows: we first discuss the soft-pion theorem. The relevant degrees of freedom are then introduced, and these enable us to write down chiral Lagrangian densities. Inelastic cross sections are extracted and the soft-pion theorem is explicitly verified. Once parameters are fixed and symmetry preserving form factors are introduced, the relative strengths of chiral symmetry, abnormal parity content, and ρ-dissociation effects on the cross sections are presented and discussed.

II. DECOUPLING OF PIONS IN THE SOFT-MOMENTUM LIMIT

First, consider the case where the chiral symmetry is exact. The axial current for the Goldstone realization of the chiral symmetry [49] is

$$A_\mu(x) = f_\pi^0 \partial_\mu \pi + \cdots, \quad (1)$$

where f_π^0 is the pion decay constant in the chiral limit. Using the LSZ reduction formulas [50] and following Weinberg [51], the expectation value of the current between an arbitrary in and out state becomes

$$\int d^4x \langle \alpha | A_\mu(x) | \beta \rangle \, e^{-ip\cdot x} = \frac{p^\mu f_\pi^0}{p^2} i \mathcal{M}^{0\pi}_{\beta\to\alpha} + \mathcal{N}^\mu_{\beta\to\alpha}, \quad (2)$$

where $\mathcal{M}^{0\pi}_{\beta\to\alpha}$ is the transition amplitude in the chiral limit for the absorption of an incoming pion with momentum p, and $\mathcal{N}^\mu_{\beta\to\alpha}$ are the regular terms near the pion pole. Contracting the pion momentum on both side yields the current conservation condition:

$$\langle \alpha | p_\mu A^\mu(p) | \beta \rangle = f_\pi^0 i \mathcal{M}^{0\pi}_{\beta\to\alpha} + p_\mu \mathcal{N}^\mu_{\beta\to\alpha} = 0. \quad (3)$$

Under the assumption that $\mathcal{N}^\mu_{\beta\to\alpha}$ is regular near the pion pole, the pion then decouples in the soft-momentum limit giving

$$\mathcal{M}^{0\pi}_{\beta\to\alpha} \to 0. \quad (4)$$

This constraint, first studied by Adler [52], is an example of how the chiral symmetry manifests itself in the Goldstone mode for low-energy scattering. A general proof with many pions can be obtained [51].

Knowing that chiral symmetry is only partially conserved, let us now consider how the above theorem is modified. Under the PCAC hypothesis [49], the current matrix elements now become

$$\langle \alpha | A^\mu(p) | \beta \rangle = \frac{p^\mu f_\pi}{p^2 - m_\pi^2} i \mathcal{M}^{\pi}_{\beta\to\alpha} + \mathcal{N}^\mu_{\beta\to\alpha}, \quad (5)$$

where f_π is the decay constant for an explicitly broken chiral symmetry. Assuming that the explicit chiral breaking occurs only through a m_π^2 dependence the pion-absorption transition amplitude reduces to

$$\lim_{p\to 0} \mathcal{M}^{\pi}_{\beta\to\alpha} = \lim_{p\to 0} \mathcal{M}^{0\pi}_{\beta\to\alpha} \to 0. \quad (6)$$

This is a strong version of the smoothness assumption [53,54] that requires that the amplitude does not change significantly from $p^2 = m_\pi^2$ to $p^2 = 0$.

The above two derivations assume that no other singularities exist besides that provided by the pion pole, or in other words, that $\mathcal{N}^\mu_{\beta\to\alpha}$ is regular. This is in general not true [51,54]. Figure 1 shows the basic subdiagram where an initial off-shell particle absorbs a pion and then emits an on-shell particle. These two particles could for example be mesons. If the two particles have the same mass, then a kinematical singularity

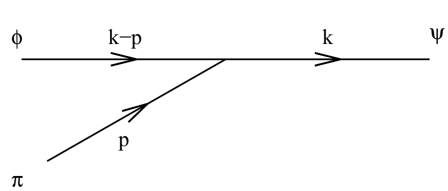

FIG. 1. Exception to the decoupling theorem due to a kinematic singularity.

will develop in the soft-pion limit, i.e., the denominator of the incoming particle propagator becomes

$$\lim_{p\to 0} (k-p)^2 - M^2 = \lim_{p\to 0} k^2 - 2k \times p + p^2 - M^2 \to k^2 - M^2 \to 0,$$

where M is the mass of the two particles and $k^2 = M^2$ because the outgoing particle is on shell.

This singularity can occur in two cases. First, when the incoming and outgoing particles have degenerate masses as it is sometimes realized for a chirally restored vacuum (e.g., σ and π mesons are degenerate). The other case manifests itself for abnormal parity interactions which permit the absorption/emission of a pion from a single particle. An example of such a process is between a pion and an isospin-doublet vector meson with negative parity, V, i.e.,

$$\mathcal{L}_{\pi VV} = g \epsilon^{\mu\nu\alpha\beta} \partial_\mu V_\nu \tau \cdot \pi \, \partial_\alpha V^\dagger_\beta, \quad (7)$$

where $\epsilon^{\mu\nu\alpha\beta}$ is the four-dimensional antisymmetric tensor. This interaction will then generate a singularity for soft-pion kinematics because the incoming and outgoing particles are identical.

III. DEGREES OF FREEDOM AND CHIRAL SYMMETRY

We now wish to build a chirally invariant Lagrangian. Doing so will ensure that the soft-pion limit is exhibited by the model. The principal difficulty is to identify the relevant degrees of freedom. In the final stage of a heavy-ion collision, the relative momentum of the J/ψ and a light meson is of the order of a few GeV: the charmonium dissociation is thus expected to be dominated by those processes with the smallest excitation threshold, i.e., cross sections with the lowest-mass final states. Therefore, it is sufficient to consider the dissociation processes into the lowest-mass open charmed mesons resulting from the interactions between the J/ψ, D, D^*, and the light mesons.

This point is incorrect if chiral symmetry is to be maintained. Indeed, as pointed out in Ref. [48] inclusion of the chiral partners is essential for the decoupling theorem to hold. It is thus expected that the chiral partners, even though they do not appear in the final states, still can play an important role through exchange diagrams. With this in mind, identifying the chiral partners of the D and D^* mesons is essential. Because they are pseudoscalar and vector mesons, respectively, their chiral partners are expected to be scalar and axial-vector particles. Moreover, under the heavy-quark spin symmetry, they should have similar masses. We see from Ref. [55] that the D_0^* and D_1 mesons are candidates for the scalar and axial partners, respectively.[1]

Introducing the chiral partners amounts to having a linear realization of chiral symmetry. One could also decide not to introduce these additional mesons, and consider a nonlinear realization of chiral symmetry by letting, for example, the chiral partner masses go to infinity. The D and D^* would have

[1]Ref. [55] states that the quantum numbers of the D_0^* and D_1 mesons have to be confirmed.

nonlinear transformation properties under the axial subgroup. This approach is the one used in building the Lagrangians incorporating heavy-quark spin-flavor symmetry [26]. For this study, the linear representation will be used. The open charmed mesons will be then the D, D^*, D_0^*, and D_1.

To build a chiral invariant Lagrangian, it is convenient to define chiral fields. In Appendix A these are identified by considering the various possible quark bilinears. Knowing the transformation properties under chiral symmetry of the light and heavy quarks then permit to determine that of the mesons. The chiral fields are then found to be

$$W = \sigma + i\pi, \tag{8}$$

$$W^\dagger = \sigma - i\pi, \tag{9}$$

$$A_{R,L} = \rho \pm a_1, \tag{10}$$

$$D_{R,L} = D_0^* \pm iD, \tag{11}$$

$$D_{R,L}^* = D^* \pm D_1, \tag{12}$$

where W and $A_{R,L}$ are isospin triplets and $D_{R,L}$ and $D_{R,L}^*$ are isospin doublets.

IV. LAGRANGIAN DENSITIES

We first write down the free field Lagrangian by defining the following field strengths:

$$F_{R,L}^{\mu\nu} = \partial^\mu A_{R,L}^\nu - \partial^\nu A_{R,L}^\mu \tag{13}$$

for an arbitrary left- and right-handed vector field. Then, starting from the linear σ model, the free field Lagrangian reads

$$\begin{aligned}\mathcal{L}_0 = {} & \frac{1}{4}Tr[\partial_\mu W \partial^\mu W^\dagger] - \frac{1}{4}\mu^2 Tr[WW^\dagger] \\ & + \frac{f_\pi m_\pi^2}{4}Tr[W + W^\dagger] - \frac{1}{16}Tr\left[F_{\mu\nu}^L F_L^{\mu\nu} + F_{\mu\nu}^R F_R^{\mu\nu}\right] \\ & + \frac{m_0^2}{4}Tr\left[A_{L\mu}A_L^\mu + A_{R\mu}A_R^\mu\right] \\ & + \frac{1}{2}(\partial_\mu D_L \partial^\mu \bar{D}_L + \partial_\mu D_R \partial^\mu \bar{D}_R) \\ & - \frac{M^2}{2}(D_L\bar{D}_L + D_R\bar{D}_R) - \frac{1}{8}\left(F_{\mu\nu}^{D_L^*}F_{\bar{D}_L^*}^{\mu\nu} + F_{\mu\nu}^{D_R^*}F_{\bar{D}_R^*}^{\mu\nu}\right) \\ & + \frac{M^{*2}}{2}\left(D_L^{*\mu}\bar{D}_{L\mu}^* + D_R^{*\mu}\bar{D}_{R\mu}^*\right),\end{aligned} \tag{14}$$

where M and M^* are the degenerate masses of open charmed mesons and m_0 that of the ρ and a_1 mesons. Degeneracies will be lifted by spontaneous chiral symmetry breaking once the interactions are included, as in the linear σ model [49], which will result in mass splittings between D and D_0^* as well as between D_1. A pion mass has also been included with the third term, and thus chiral symmetry is explicitly broken.

For the interactions, the working assumption here will be that only the three- and four-point interactions with the lowest number of derivatives are to be considered.[2] Because the Lagrangian density is of dimension four and the mesonic fields are of dimension one, the three-point interactions will have couplings scaling as M^{1-n}, where M is an arbitrary mass scale and n is the number of derivatives, whereas the four-point interactions having one more field operator will scale as M^{-n}. Furthermore, only the minimal interactions with the chiral partners of the D and D^* mesons will be added to maintain chiral symmetry. Practically, this implies that all the interactions with D_0^* and D_1 fields will be generated by the spontaneous chiral symmetry breaking. Moreover, only the three- and four-point interaction terms necessary to contruct the amplitudes with the considered final states are explicitly written down. Finally, aside from the requirement that the Lagrangian density be real, the other tools used to construct the phenomenological Lagrangian are parity and charge conjugation invariances (which are valid symmetries of quantum chromodynamics). The effects of these discrete transformations on the field content, as well as the resulting interactions, are listed in Appendix B.

The next step is to make explicit the spontaneous chiral symmetry breaking by shifting the σ field in W by $\sigma \to \sigma + \sigma_0$ as in the linear sigma model. Doing so yields the new free field Lagrangian

$$\begin{aligned}\mathcal{L}_0 = {} & \tfrac{1}{2}\partial_\mu\pi\partial^\mu\pi - m_\pi^2\pi^2 + \tfrac{1}{2}\partial_\mu\sigma\partial^\mu\sigma - m_\sigma^2\sigma^2 \\ & - \tfrac{1}{8}Tr\left[F_\rho^{\mu\nu}F_{\mu\nu}^\rho\right] + \tfrac{1}{4}m_0^2 Tr\left[\rho_\mu^2\right] - \tfrac{1}{8}Tr\left[F_{a_1}^{\mu\nu}F_{\mu\nu}^{a_1}\right] \\ & + \tfrac{1}{4}m_0^2 Tr\left[a_{1\mu}^2\right] + \partial_\mu D\partial^\mu\bar{D} - (M^2 - 2\Delta\sigma_0)D\bar{D} \\ & + \partial_\mu D_0^*\partial^\mu\bar{D}_0^* - (M^2 + 2\Delta\sigma_0)D_0^*\bar{D}_0^* \\ & - \tfrac{1}{4}F_{\mu\nu}^{D^*}F_{\bar{D}^*}^{\mu\nu} + (M^{*2} - 2\Delta^*\sigma_0)D_\mu^*\bar{D}^{*\mu} \\ & - \tfrac{1}{4}F_{\mu\nu}^{D_1}F_{\bar{D}_1}^{\mu\nu} + (M^{*2} + 2\Delta^*\sigma_0)D_{1\mu}\bar{D}_1^\mu \\ & + ig_{WDD^*}^{(0)}\sigma_0(\partial_\mu D_0^*\bar{D}^{*\mu} - D^{*\mu}\partial_\mu\bar{D}_0^*) \\ & + g_{WDD^*}^{(0)}\sigma_0\left(\partial_\mu D\bar{D}_1^\mu - D_1^\mu\partial_\mu\bar{D}\right),\end{aligned} \tag{15}$$

where the expressions for m_π and m_σ are the same as for the linear σ model [49], and $g_{WDD^*}^{(0)}$, Δ, and Δ^* are coupling constants. We note that the introduction of interactions generate mass splittings between the D and D_0^* mesons and between the D^* and D_1 mesons respectively; thus lifting the mass degeneracies. The D meson masses then read

$$m_D^2 = M^2 - 2\Delta\sigma_0, \quad m_{D_0^*}^2 = M^2 + 2\Delta\sigma_0,$$
$$m_{D^*}^2 = M^{*2} - 2\Delta^*\sigma_0, \quad m_{D_1}^2 = M^{*2} + 2\Delta^*\sigma_0.$$

Moreover, the introduction of the interactions induces mixing between D_0^* and D^* fields and between D and D_1 fields. To cast the Lagrangian into a canonical form would thus require making field redefinitions. These are involved and would lead to additional interactions with higher powers of momentum, which is contrary to the original assumption of limiting possible interactions to those with the lowest powers of momentum. Moreover, the nonchiral invariant model of Ref. [45] with which we wish to make comparison has no such mixings. For this study, the coupling constant $g_{WDD^*}^{(0)}$ is thus set to zero removing the mixing. Finally, in this model, the ρ and a_1 mesons have degenerate masses. This is of no importance here because we wish to compute only the cross

[2]This assumption is strictly valid only if all terms with higher powers of derivative are suppressed for the considered kinematical regime [53].

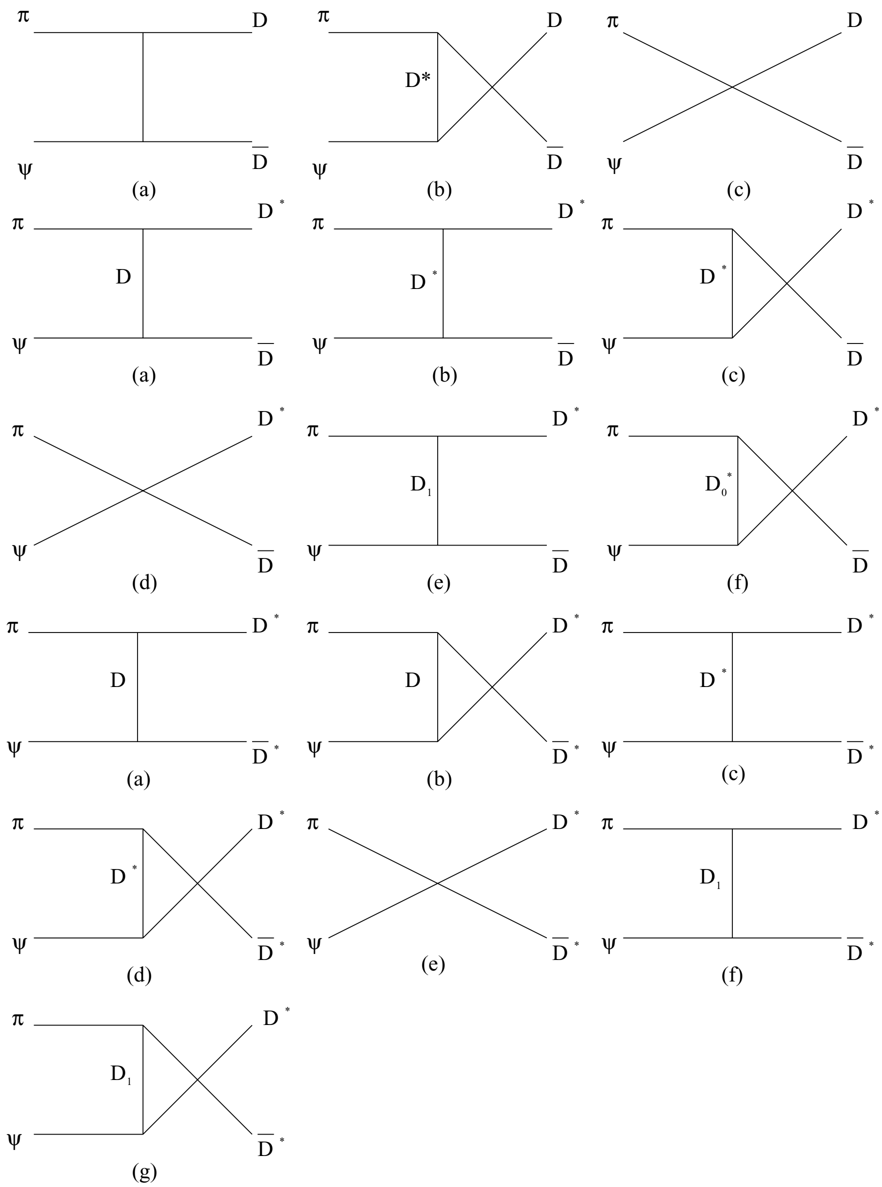

FIG. 2. Diagrams for $\pi + \psi$ dissociation.

sections with the two lightest mesons, namely the π and ρ. In Appendix B, the relevant interactions for the J/ψ meson by a pion or a ρ meson are listed. They include normal and abnormal interactions. As discussed in Sec. II, the latter are expected to circumvent the low-energy theorem.

V. INELASTIC SCATTERING AMPLITUDES

All amplitudes discussed in this section are explicitly written down in Appendix C.

A. $\pi + J/\psi$

The pion dissociation of the J/ψ proceeds through three processes, namely:

$$\mathcal{M}_1 = \sum_i \mathcal{M}^{\rho}_{1i} \epsilon_\rho(p_\psi), \tag{16}$$

$$\mathcal{M}_2 = \epsilon^*_\mu(p_{D^*}) \sum_i \mathcal{M}^{\mu\rho}_{2i} \epsilon_\rho(p_\psi), \tag{17}$$

$$\mathcal{M}_3 = \epsilon^*_\mu(p_{D^*}) \epsilon^*_\nu(p_{\bar{D}^*}) \sum_i \mathcal{M}^{\mu\nu\rho}_{3i} \epsilon_\rho(p_\psi), \tag{18}$$

where $\epsilon_\rho(p_\psi)$, $\epsilon^*_\mu(p_{D^*})$, and $\epsilon^*_\nu(p_{\bar{D}^*})$ are the polarization vectors for the J/ψ, D^*, and $\bar{D}^*$ mesons, respectively. The first and last amplitudes arise only due to abnormal parity interactions, whereas $\mathcal{M}_2$ contains one abnormal parity exchange process [Fig. 2(b)]. Note also that the amplitude for the final state $\bar{D}^* D$ is obtained from the conjugate of amplitude $\mathcal{M}_2$.

B. $\rho + J/\psi$

For the ρ-meson-induced dissociation, three processes are examined:

$$\mathcal{M}_4 = \sum_i \mathcal{M}^{\delta\rho}_{4i} \epsilon_\rho(p_\psi) \epsilon_\delta(p_\rho) \tag{19}$$

$$\mathcal{M}_5 = \epsilon^*_\mu(p_{D^*}) \sum_i \mathcal{M}^{\mu\delta\rho}_{5i} \epsilon_\rho(p_\psi)\epsilon_\delta(p_\rho), \quad (20)$$

$$\mathcal{M}_6 = \epsilon^*_\mu(p_{D^*})\epsilon^*_\nu(p_{\bar{D}^*}) \sum_i \mathcal{M}^{\mu\nu\delta\rho}_{6i} \epsilon_\rho(p_\psi)\epsilon_\delta(p_\rho), \quad (21)$$

where $\epsilon_\delta(p_\rho)$ is the polarization vector of the ρ meson. Again, the conjugate of $\mathcal{M}_5$ gives the amplitude for the $\bar{D}^*D$ final state. It is important to note that the chiral symmetry constraint does not introduce additional amplitudes involving the exchange of the D_0^* and the D_1 as in the case of the dissociation with pions. Consequently, the diagrams are the same as in Ref. [45], and we expect the results to agree.

C. Soft-pion limit

We now wish to demonstrate the soft-pion theorem for the dissociation of J/ψ meson by a pion into a D^*-$\bar{D}$ final state. It is expected that this property of chiral symmetry will soften the threshold behavior. Explicitly, this will be due to a cancellation of the contact term for the normal-parity subprocesses. The caveat here is of course that abnormal-parity interactions circumvent the theorem and it will still be possible to have a contact behavior near the threshold due to these [Eq. (C5)].

With this in mind and in the chiral limit, i.e, for massless pions, we let the pion momentum go to zero. It is trivial to see that the first subamplitude (Fig. 2), due to the exchange of a D meson [Eq. (C4)] decouples when the vector mesons are on–shell since their polarization vector then satisfies the orthogonality condition, i.e., $\epsilon(p) \times p = 0$. Similarly, the u-channel D^* exchange amplitude goes to zero. We are thus left with three normal parity amplitudes, including a contact term. In the soft-pion limit we have

$$\mathcal{M}_{2e} \rightarrow \frac{(2\Delta^*)(2g_{W\psi DD^*}\sigma_0)}{m^2_{D^*} - m^2_{D_1}} g^{\mu\alpha}\left(g_{\alpha\beta} - \frac{p_{D^*\alpha}p_{D^*\beta}}{m^2_{D_1}}\right) g^{\beta\rho}, \quad (22)$$

and

$$\mathcal{M}_{2f} \rightarrow \frac{(2\Delta)(2g_{W\psi DD^*}\sigma_0)}{m^2_D - m^2_{D_0^*}} g^{\mu\rho}. \quad (23)$$

Remembering that the the mass splittings between the D mesons are due to spontaneous chiral symmetry breaking, we can further write

$$\mathcal{M}_{2e} \rightarrow -\frac{(2\Delta)(2g_{W\psi DD^*}\sigma_0)}{4\Delta^*\sigma_0} g^{\mu\rho} = -g_{W\psi DD^*} g^{\mu\rho} \quad (24)$$

and

$$\mathcal{M}_{2f} \rightarrow -\frac{(2\Delta)(2g_{W\psi DD^*}\sigma_0)}{4\Delta\sigma_0} g^{\mu\rho} = -g_{W\psi DD^*} g^{\mu\rho} \quad (25)$$

where for the amplitude $\mathcal{M}_{2e}$ the orthogonality condition has been used to remove the term proportional to the product of four-vectors. Adding these two contributions to the contact term of $\mathcal{M}_{2d}$ leads to the desired result for the normal parity content. Because the contraction of the two ε-tensors results into a sum of products of the metric tensors, the leading behavior near the threshold for the process $\pi + \psi \rightarrow \bar{D} + D^*$ will be given by the amplitude $\mathcal{M}_{2b}$.

A remark is in order regarding the chiral limit. Relaxing this assumption will make the amplitudes $\mathcal{M}_{2e}$ and $\mathcal{M}_{2f}$ depend on the pion mass. It is trivial to see that these can be mapped smoothly into the chiral amplitudes considered above by letting the pion mass go to zero, thus satisfying the smoothness assumption of the decoupling theorem.

VI. CROSS SECTIONS FOR DISSOCIATION PROCESSES

A. Introducing symmetry conserving form factors

To complete the description of the phenomenological model, form factors must be introduced to account for the substructure of mesons. A Lorentz-invariant three-point form factor is introduced, namely

$$\mathcal{F}^M_3(q^2) = \frac{\Lambda^2}{\Lambda^2 + |q^2 - m^2_M|}, \quad (26)$$

where q^2 is the virtuality, m_M is the meson mass, and Λ is the range parameter. The cutoff parameter will be set to two different values, namely 1 and 2 GeV, as in previous studies [41,43,45]. These can be justified by noting that the typical hadronic scale is about 1 GeV and the exchanged mesons, which are open charmed mesons here, have masses of about 2 GeV. One could relax the universality condition by introducing a different cutoff parameter for each interaction, but the assumption of a common Λ is a realistic first approximation because the exchanged mesons are all D mesons.

The astute reader will note that the coupling constants should strictly be defined at the point where the form factor is one, i.e., $q^2 = m^2_M$. This is not the case for all the coupling values extracted in Ref. [45] that are used here. Indeed, the three-point couplings involving a ρ or a J/ψ meson are evaluated with these particles at zero virtuality. Nevertheless, it will still be assumed that the couplings extracted with the ρ or the J/ψ meson off-shell are the same as those on-shell.

A form factor for the four-point interactions is also introduced. Here, a dipole form is chosen, namely

$$\mathcal{F}_4(s,t) = \frac{\Lambda^2}{\Lambda^2 + |t - M_0^2|}\frac{\Lambda^2}{\Lambda^2 + |u - M_0^2|}, \quad (27)$$

where $s + t + u = m_1^2 + m_2^2 + m_3^2 + m_4^2$ and M_0 is a mass scale. This latter parameter is given by the average of the D and D^* masses, i.e., $M_0 = 1.94$ GeV. The four-point form factor is then equal to 1 when $t = u = M_0^2$. Strictly speaking the normalization, i.e., the coupling constant, is defined at this point.

The above discussion omits the constraint due to chiral symmetry. Indeed, some of the three-point form factors are determined by four-point form factors. This is the case for all three-point interactions generated by underlying four-point interactions, i.e., which have a W-field factor (see Appendix B for details). Specifically, let us consider the $W\psi DD^*$ interaction from which the $\pi\psi DD^*$, ψDD_1, and $\psi D^* D_0^*$ interactions are generated after spontaneous chiral symmetry breaking. The three-point form factors can then be extracted

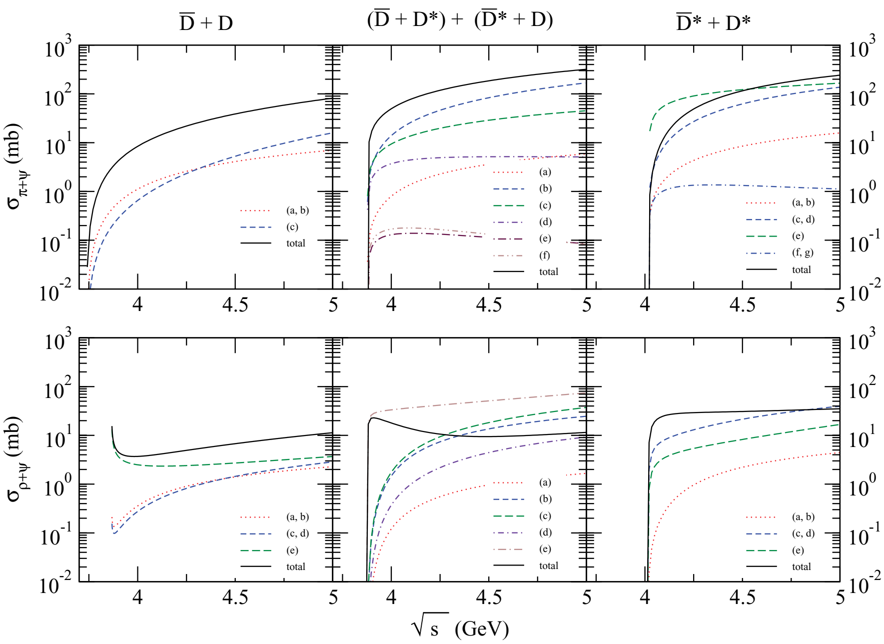

FIG. 3. (Color online) Dissociation cross sections without form factors.

from the four-point form factor by letting the pion momentum go to zero. Specifically, assume that the D^* and D mesons are off- and on-shell, respectively, then setting the pion momentum to zero yields the desired form factor for the $\psi D^* D_0^*$:

$$\lim_{p_\pi \to 0} \mathcal{F}_4(s,t) = \frac{\Lambda^2}{\Lambda^2 + |m_{D^*}^2 - M_0^2|} \frac{\Lambda^2}{\Lambda^2 + |u - M_0^2|} = \gamma_{D^*} \mathcal{F}_3^0(u) \tag{28}$$

where the index on the three-point form factor indicates that the parameter m_M is set to M_0, and $\gamma_M = \mathcal{F}_3^0(m_M^2)$. Taking the D meson off-shell and keeping the D^* on-shell gives the form factor for $\psi D D_1$ interaction. The same argument applies for the abnormal parity $\psi D^* D_1$ and $\psi D_0^* D$ interactions.

There is also another subtlety when it comes to the interactions generating the mass splittings of the D mesons, i.e., those coming from $\mathcal{L}_{WDD}$ and $\mathcal{L}_{WD^*D^*}$. Indeed, the interaction form factors will now appear in the mass shifts leading to self-consistent equations. For example, for the D^*-D_1 mass splitting, we have

$$m_{D_1}^2 - m_{D^*}^2 = 2\Delta^* \sigma_0 \lim_{p_\pi \to 0} \left(\mathcal{F}_3^{D_1}(q^2) + \mathcal{F}_3^{D^*}(q^2)\right) = 4\Delta^* \sigma_0 \frac{\Lambda^2}{\Lambda^2 + |m_{D_1}^2 - m_{D^*}^2|}. \tag{29}$$

From these, we see that the values of the interaction strengths, Δ^* and Δ, are functions of both the cutoff parameter and the mass scale.

In light of these modifications, we re-examined the soft-pion limit for the $\mathcal{M}_2$ amplitude. $\mathcal{M}_{2d}$ is now given by

$$\lim_{p_\pi \to 0} \mathcal{M}_{2d} = 2 g_{W\psi DD^*} \gamma_{D^*} \gamma_D g^{\mu\rho}, \tag{30}$$

whereas $\mathcal{M}_{2e}$ and $\mathcal{M}_{2f}$ reduce to

$$\lim_{p_\pi \to 0} \mathcal{M}_{\{2e,2f\}} = \lim_{p_\pi \to 0} \left(\gamma_D \mathcal{F}_3^0(t) \mathcal{F}_3^{D_1}(t)\right) \times \frac{(2\Delta^*)(2 g_{W\psi DD^*} \sigma_0)}{m_{D^*}^2 - m_{D_1}^2} g^{\mu\rho} = -\gamma_D \gamma_{D^*} g_{W\psi DD^*} g^{\mu\rho}, \tag{31}$$

where Eq. (29) has been used to go from the first line to the second.

B. Results

The cross sections are first studied without form factors. The parameters used in the calculation can be found in Appendix D. The six cross sections are presented in Fig. 3 where the solid curves are the cross sections including all subamplitudes. Overall, near threshold both dissociation by a pion and by a ρ meson are of the same order of magnitude; the pion dissociation starts to dominate over the ρ dissociation beyond 4 GeV.

The effect of introducing chiral symmetry can be assessed by considering the pion absorption into the $(\bar{D} + D^*) + (\bar{D}^* + D)$ final state. The leading contribution to this process is due to the subamplitude $\mathcal{M}_{2b}$, which arises because of the abnormal parity content in the Lagrangian. This is made clear in both Figs. 3 and 4. Indeed, at a value of the center-of-mass energy of 3.9 GeV, excluding this subamplitude reduces the cross section by 65%. In contrast, removing $\mathcal{M}_{2e}$ and $\mathcal{M}_{2f}$, i.e., the subamplitudes necessary to maintain chiral symmetry, increases the cross section by only 27%. Although omitting the chiral constraint increases the cross section, as expected from the soft-pion theorem, the effect is subleading compared to the inclusion of abnormal parity interactions. Moreover, the presence of abnormal parity content makes the dissociation into $\bar{D}D$ and $\bar{D}^* D^*$ pairs possible, which further increases the total pion-absorption cross section.

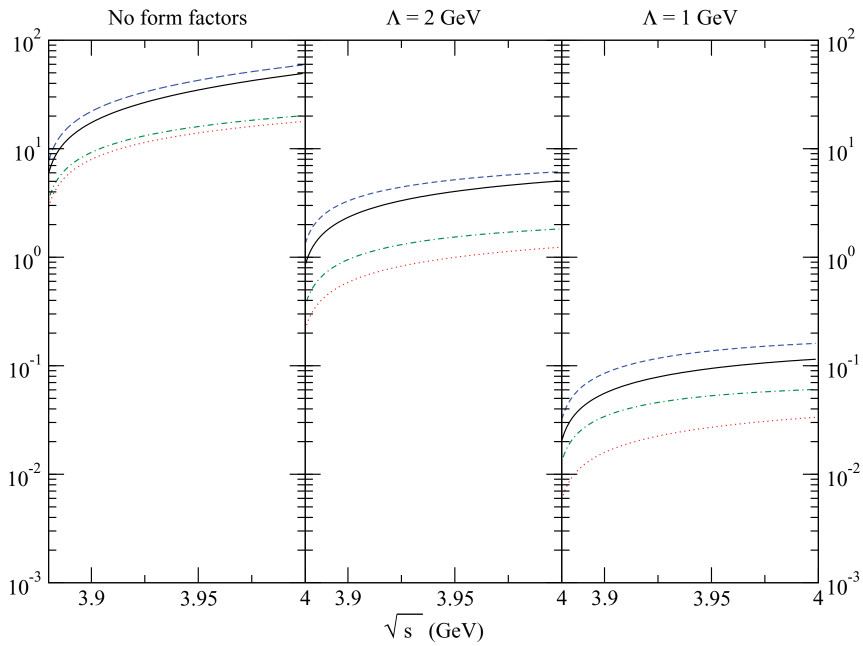

FIG. 4. (Color online) Effects of chiral symmetry and abnormal parity content on the $\pi + J/\psi \to (\bar{D} + D^*) + (\bar{D}^* + D)$ cross section. The dotted, dashed, and dot-dashed lines correspond to cross sections without the abnormal parity subamplitude, without the two subamplitudes due to chiral symmetry and without all three subamplitudes. The total inclusive cross section with all contributions is shown with the solid lines.

Turning to ρ dissociation, the initial expectation is that the results of Ref. [45] should be reproduced because no additional interactions are introduced in applying the chiral symmetry constraint. All three pion-absorption cross sections are monotonically increasing with $\sqrt{s}$ and featureless as in Ref. [45]. However, the three ρ-dissociation cross sections differ qualitatively in shape when compared to results from Ref. [45]. In spite of the fact that the interactions and the squared subamplitudes are the same, the relative phases and, consequently, the interference patterns are different, leading to the observed dissimilarities.

The above discussion is valid only when form factors are omitted. In Fig. 5, cross sections with and without form factors are compared. Two cases of the cutoff parameter are considered, namely 1 and 2 GeV. As expected, for decreasing Λ the suppression is increased. Overall, it is clear that the magnitudes of the two dissociation channels are set by the functional forms of the form factors and the values of the parameters. With this caveat, the inclusive pion-dissociation cross section is of the order of a few millibarns near threshold for a cutoff of 2 GeV, and a fraction of millibarn for $\Lambda = 1$ GeV.

Finally, the relative effect of chiral symmetry as the cutoff is lowered is shown in Fig. 4. At $\sqrt{s} = 3.9$ GeV, the cross section for $\Lambda = 1$ GeV increases by 51% when the subamplitudes due to chiral symmetry are neglected, whereas

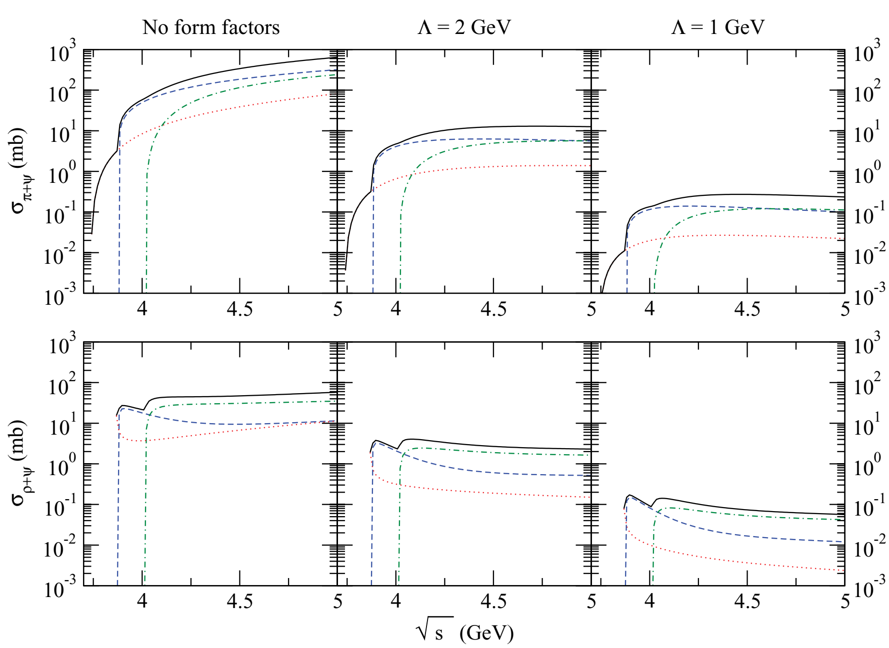

FIG. 5. (Color online) Comparison of the dissociation cross sections with and without form factors. The dotted, dashed, and dot-dashed lines correspond to dissociation into $\bar{D} + D$, $(\bar{D} + D^*) + (\bar{D}^* + D)$, and $\bar{D}^* + D^*$. The total inclusive cross sections are given by the solid lines; cusps are due to channels opening.

it decreases by 72% when the abnormal parity content is omitted.

VII. CONCLUSION AND OUTLOOK

We presented an extension of the work done in Ref. [48] where, in addition introducing chiral symmetric interactions, abnormal parity content and ρ mesons are also included. The former is important because the soft-pion theorem is circumvented in this case, whereas the latter is a first step toward assessing the relative importance of the J/ψ dissociation by other light resonances. To account for the quark substructure of mesons, *ad hoc* mesonic form factors were also added. Comparing the ρ-induced dissociation with the pion ones shed no more light than what was found in Ref. [45]. Any statements about the relative strength between π- and ρ-induced dissociation depend heavily on the choice of form factors and the techniques used to fix their absolute normalizations and are thus model dependent.

We also conclude that there are some indications that the introduction of chiral symmetry does reduce the cross section of $\pi + J/\psi \to (\bar{D} + D^*) + (D + \bar{D}^*)$ but also that the implementation of abnormal parity content is probably as or even more important because it increases not only the maximum reached by the $\pi + J/\psi \to (\bar{D} + D^*) + (D + \bar{D}^*)$ cross section, but also it allows new decay channels, such as $\pi + J/\psi \to \bar{D}^* + D^*$, to open.

In the future, the J/ψ-dissociation rates will be integrated in an evolving hot and dense medium. Introducing other light mesons, such as the ω, as well as higher charmonium resonances will also be considered to improve the phenomenological description. Moreover, adding final states incorporating D_0^* and D_1 mesons and evaluating the cross sections for the inverse reactions will also figure with the additional developments. Then, contact with the phenomenology measured at the CERN Super Proton Synchrotron and at the Relativistic Heavy Ion Collider will be made.

ACKNOWLEDGMENTS

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APPENDIX A: FIELD REPRESENTATIONS AND CHIRAL SYMMETRY

To write down all the possible invariant interactions between the mesons, it is essential to know their chiral transformation properties. Obviously, for the J/ψ meson this is trivial as it is a singlet of the chiral group. For the π, ρ, D, D^*, and their chiral partners, it is convenient to define chiral fields.

The field representations of the π and σ mesons are given by

$$W = \sigma + i\pi \tag{A1}$$

$$W^\dagger = \sigma - i\pi. \tag{A2}$$

Their transformation properties under the $SU_L(2) \times SU_R(2)$ group can be assessed by coupling the chiral meson fields to quark bilinears of corresponding parity giving

$$\bar{q}\,(\sigma + i\gamma_5\pi)\,q, \tag{A3}$$

where $\pi = \tau^a\pi^a$. Projecting the quark fields into their left- and right-handed representations yields

$$\bar{q}(\sigma + i\gamma_5\pi)q = \bar{q}_L W q_R + \bar{q}_R W^\dagger q_L. \tag{A4}$$

Under a chiral transformation of the light quark fields as defined by

$$q_{R,L} \to U_{R,L} q_{R,L} = e^{-i\tau^i\epsilon^i_{R,L}} q_{R,L}, \tag{A5}$$

where τ^i are the SU(2) Pauli matrices satisfying the normalization condition $Tr\left(\tau^i\tau^j\right) = 2\delta^{ij}$, the chiral mesonic fields have to transform as

$$W \to U_L W U_R^\dagger \tag{A6}$$

$$W^\dagger \to U_R W U_L^\dagger \tag{A7}$$

for the interaction to be invariant.

The spin-1 light mesons will not be introduced as gauge bosons as in Ref. [48]. Applying the same technique as for the σ and π fields yields the interaction

$$\bar{q}\,(\not{\rho} + \not{a}_1\gamma_5)\,q = \bar{q}_L A_L q_L + \bar{q}_R A_R q_R, \tag{A8}$$

where now $\rho_\mu = \rho^a_\mu\tau^a = \frac{1}{2}(A_R + A_L)$ and $a_{1\mu} = a^a_{1\mu}\tau^a = \frac{1}{2}(A_R - A_L)$. From these we infer that

$$A^L_\mu \to U_L A^L_\mu U_L^\dagger \tag{A9}$$

$$A^R_\mu \to U_R A^R_\mu U_R^\dagger, \tag{A10}$$

which do not transform as gauge bosons.

Turning now to the open charmed mesons, we consider first the D and D_0^* isospin doublet fields and their conjugates that are written as

$$\begin{aligned} \bar{D}^T &= (\bar{D}^0, D^-), \quad D = (D^0, D^+) \\ \bar{D}_0^{*T} &= \left(\bar{D}_0^{*0}, D_0^{*-}\right), \quad D_0^* = \left(D_0^{*0}, D_0^{*+}\right), \end{aligned} \tag{A11}$$

where T is the transposition operator. These can be rearranged into

$$\bar{D}_{R,L} = \left(\bar{D}_0^* \mp i\bar{D}\right), \quad D_{R,L} = (D_0^* \pm iD). \tag{A12}$$

Because the open charmed mesons have only one light valence quark, they are expected to transform under chiral symmetry according to

$$\begin{aligned} \bar{D}_{R,L} &\to U_{R,L}\bar{D}_{R,L} \\ D_{R,L} &\to D_{R,L}U_{R,L}^\dagger, \end{aligned} \tag{A13}$$

which can be made explicit by considering the coupling to the quark bilinears:

$$\begin{aligned} \bar{Q}(D_0^* + iD\gamma_5)q &= \bar{Q}_L D_R q_R + \bar{Q}_R D_L q_L \\ \bar{q}(\bar{D}_0^* + i\bar{D}\gamma_5)Q &= \bar{q}_L \bar{D}_L Q_R + \bar{q}_R \bar{D}_R Q_L. \end{aligned} \tag{A14}$$

Similarly, the D^* and D_1 fields can be cast into chiral forms yielding

$$\bar{D}^*_{R,L} = (\bar{D}^* \pm \bar{D}_1), \quad D^*_{R,L} = (D^* \pm D_1) \tag{A15}$$

TABLE I. Field transformation properties under parity and charge conjugation.

	$\psi^\mu(J/\psi)$	$W(W^\dagger)$	$D_{R,L}(\bar{D}_{R,L})$	$D^{*\mu}_{R,L}(\bar{D}^{*\mu}_{R,L})$	$A^\mu_{R,L}$	∂_μ	$\epsilon^{\mu\nu\alpha\beta}$
P	$-\psi^\mu$	$W^\dagger(W)$	$D_{L,R}(\bar{D}_{L,R})$	$-D^{*\mu}_{L,R}(-\bar{D}^{*\mu}_{L,R})$	$-A^\mu_{L,R}$	$-\partial_\mu$	$-\epsilon^{\mu\nu\alpha\beta}$
C	$-\psi^\mu$	$W^*(W^T)$	$\bar{D}^T_{R,L}(D^T_{R,L})$	$-\bar{D}^{*\mu T}_{R,L}(-D^{*\mu T}_{R,L})$	$-A^{\mu T}_{R,L}$	$+\partial_\mu$	$-\epsilon^{\mu\nu\alpha\beta}$

and the quark-meson interactions then read

$$\begin{aligned}\bar{Q}(\not{D}^* + \not{D}_1\gamma_5)q &= \bar{Q}_L D^*_R q_R + \bar{Q}_R D^*_L q_L \\ \bar{q}(\bar{\not{D}}^* + \bar{\not{D}}_1\gamma_5)Q &= \bar{q}_L \bar{D}^*_L Q_L + \bar{q}_R \bar{D}^*_R q_R\end{aligned} \tag{A16}$$

from which transformation properties similar to Eq. (A13) are deduced.

APPENDIX B: CHIRAL INVARIANT INTERACTIONS

Table I lists the chiral field properties under discrete transformations. They are particularly useful to fix the relative signs of the interaction terms. Moreover, the concepts of normal and abnormal parity interactions are also introduced as a classification. Abnormal parity interactions have an ε-tensor factor.

The three-point normal-parity interactions are then:

$$\mathcal{L}_{WDD} = -\Delta(D_L W \bar{D}_R + D_R W^\dagger \bar{D}_L) \tag{B1}$$

$$\mathcal{L}_{WD^*D^*} = -\Delta^*\left(D^{*\mu}_L W \bar{D}^*_{R\mu} + D^{*\mu}_R W^\dagger \bar{D}^*_{L\mu}\right) \tag{B2}$$

$$\begin{aligned}\mathcal{L}_{WDD^*} = {}& i g^{(0)}_{WDD^*}\left(\partial_\mu D_L W \bar{D}^{*\mu}_R + \partial_\mu D_R W^\dagger \bar{D}^{*\mu}_L\right) \\ &+ i g^{(1)}_{WDD^*}\left(D_L \partial_\mu W \bar{D}^{*\mu}_R + D_R \partial_\mu W^\dagger \bar{D}^{*\mu}_L\right) \\ &+ \text{h.c}\end{aligned} \tag{B3}$$

$$\begin{aligned}\mathcal{L}_{\psi D^*D^*} = {}& i g^{(0)}_{\psi D^*D^*}\partial_\mu \psi_\nu\left(D^{*\mu}_R \bar{D}^{*\nu}_R + D^{*\mu}_L \bar{D}^{*\nu}_L\right) \\ &+ i g^{(1)}_{\psi D^*D^*}\psi_\mu\left(\partial^\mu D^*_{R\nu}\bar{D}^{*\nu}_R + \partial^\mu D^*_{L\nu}\bar{D}^{*\nu}_L\right) \\ &+ i g^{(2)}_{\psi D^*D^*}\psi_\mu\left(\partial_\nu D^{*\mu}_R \bar{D}^{*\nu}_R + \partial_\nu D^{*\mu}_L \bar{D}^{*\nu}_L\right) \\ &+ \text{h.c.},\end{aligned} \tag{B4}$$

$$\begin{aligned}\mathcal{L}_{AD^*D^*} = {}& i g^{(0)}_{AD^*D^*}\left(D^{*\nu}_R \partial_\mu A_{R\nu}\bar{D}^{*\mu}_R + D^{*\nu}_L \partial_\mu A_{L\nu}\bar{D}^{*\mu}_L\right) \\ &+ i g^{(1)}_{AD^*D^*}\left(D^*_{R\nu} A_{R\mu}\partial^\mu \bar{D}^{*\nu}_R + D^*_{L\nu} A_{L\mu}\partial^\mu \bar{D}^{*\nu}_L\right) \\ &+ i g^{(2)}_{AD^*D^*}\left(D^{*\mu}_R A_{R\nu}\partial_\mu \bar{D}^{*\nu}_R + D^{*\mu}_L A_{L\nu}\partial_\mu \bar{D}^{*\nu}_L\right) \\ &+ \text{h.c.},\end{aligned} \tag{B5}$$

$$\begin{aligned}\mathcal{L}_{\psi DD} &= i g_{\psi DD}\psi_\mu(\partial^\mu D_R \bar{D}_R + \partial^\mu D_L \bar{D}_L) + \text{h.c.}, \\ \mathcal{L}_{ADD} &= i g_{ADD}\left(D_R A^\mu_R \partial_\mu \bar{D}_R + D_L A^\mu_L \partial_\mu \bar{D}_L\right) \\ &\quad + \text{h.c.}\end{aligned} \tag{B6}$$

whereas the four-point interactions read

$$\mathcal{L}_{WWWW} = -\tfrac{1}{16}\lambda^2(Tr[WW^\dagger])^2, \tag{B7}$$

$$\begin{aligned}\mathcal{L}_{W\psi DD^*} = {}& g_{W\psi DD^*}\psi^\mu(D_L W \bar{D}^*_{R\mu} + D_R W^\dagger \bar{D}^*_{L\mu}) \\ &+ \text{h.c.},\end{aligned} \tag{B8}$$

$$\mathcal{L}_{A\psi DD} = g_{A\psi DD}\psi_\mu\left(D_R A^\mu_R \bar{D}_R + D_L A^\mu_L \bar{D}_L\right) \tag{B9}$$

$$\begin{aligned}\mathcal{L}_{A\psi D^*D^*} = {}& g^{(0)}_{A\psi D^*D^*}\psi_\mu\left(D^{*\nu}_R A^\mu_R \bar{D}^*_{R\nu} + D^{*\nu}_L A^\mu_L \bar{D}^*_{L\nu}\right) \\ &+ g^{(1)}_{A\psi D^*D^*}\psi_\mu\left(D^{*\nu}_R A^\nu_R \bar{D}^{*R\mu} + D^{*\nu}_L A^\nu_L \bar{D}^{*L\mu} \right. \\ &\left. + \text{h.c.}\right)\end{aligned} \tag{B10}$$

where h.c. refers to the Hermitian conjugate. All the coupling constants are dimensionless with the exception of Δ and Δ^*, which have dimension of mass.

Abnormal-parity interactions cannot be written down directly at this point as there remains an ambiguity in their definitions. Indeed, the interaction forms are not unique as there is a nontrivial relation called the Schouten's identity, relating different matrix elements [31,56]. To build the interactions, the gauged Wess-Zumino Lagrangian is used as a guide as in Ref. [45]. The three-point interactions are then:

$$\begin{aligned}&\mathcal{L}_{WD^*D^*} \\ &\quad = i g_{WD^*D^*}\epsilon^{\mu\nu\alpha\beta}(\partial_\mu D^*_{L\nu} W \partial_\alpha \bar{D}^*_{R\beta} - \partial_\mu D^*_{R\nu} W^\dagger \partial_\alpha \bar{D}^*_{L\beta}),\end{aligned} \tag{B11}$$

$$\begin{aligned}&\mathcal{L}_{\psi DD^*} \\ &\quad = i g_{\psi DD^*}\epsilon^{\mu\nu\alpha\beta}\partial_\mu \psi_\nu(\partial_\alpha D^*_{L\beta}\bar{D}_L - \partial_\alpha D^*_{R\beta}\bar{D}_R) + \text{h.c.},\end{aligned} \tag{B12}$$

$$\begin{aligned}&\mathcal{L}_{ADD^*} \\ &\quad = i g_{ADD^*}\epsilon^{\mu\nu\alpha\beta}(\partial_\mu D^*_{L\nu}\partial_\alpha A_{L\beta}\bar{D}_L - \partial_\mu D^*_{R\nu}\partial_\alpha A_{R\beta}\bar{D}_R) \\ &\qquad + \text{h.c.}\end{aligned} \tag{B13}$$

and the four-point interactions are given by

$$\begin{aligned}&\mathcal{L}_{W\psi DD} \\ &\quad = g_{W\psi DD}\epsilon^{\mu\nu\alpha\beta}\psi_\mu(\partial_\nu D_L \partial_\alpha W \partial_\beta \bar{D}_R - \partial_\nu D_R \partial_\alpha W^\dagger \partial_\beta \bar{D}_L),\end{aligned} \tag{B14}$$

$$\begin{aligned}&\mathcal{L}_{W\psi D^*D^*} \\ &\quad = -g^{(0)}_{W\psi D^*D^*}\epsilon^{\mu\nu\alpha\beta}\psi_\mu(D^*_{L\nu}\partial_\alpha W \bar{D}^*_{R\beta} - D^*_{R\nu}\partial_\alpha W^\dagger \bar{D}^*_{L\beta}) \\ &\qquad - g^{(1)}_{W\psi D^*D^*}\epsilon^{\mu\nu\alpha\beta}\partial_\mu \psi_\nu(D^*_{L\alpha} W \bar{D}^*_{R\beta} - D^*_{R\alpha} W^\dagger \bar{D}^*_{L\beta}),\end{aligned} \tag{B15}$$

$$\begin{aligned}&\mathcal{L}_{A\psi DD^*} \\ &\quad = g^{(0)}_{A\psi DD^*}\epsilon^{\mu\nu\alpha\beta}\psi_\mu(\partial_\nu D_R A_{R\alpha}\bar{D}^*_{R\beta} - \partial_\nu D_L A_{L\alpha}\bar{D}^*_{L\beta}) \\ &\qquad - g^{(1)}_{A\psi DD^*}\epsilon^{\mu\nu\alpha\beta}\psi_\mu(D_R A_{R\nu}\partial_\alpha \bar{D}^*_{R\beta} - D_L A_{L\nu}\partial_\alpha \bar{D}^*_{L\beta}) \\ &\qquad + \text{h.c.}\end{aligned} \tag{B16}$$

where all couplings scale as M^{-1} with the exception of $g_{W\psi DD}$, which behaves as M^{-3}.

Once chiral symmetry is spontaneously broken the relevant normal-parity interactions become

$$\mathcal{L}_{\pi DD^*_0} = -2\Delta(D^*_0\pi\bar{D} + D\pi\bar{D}^*_0), \tag{B17}$$

$$\mathcal{L}_{\pi D^*D_1} = -2\Delta^* i\left(D^*_\mu \pi \bar{D}^\mu_1 - D^\mu_1 \pi \bar{D}^*_\mu\right), \tag{B18}$$

$$\begin{aligned}\mathcal{L}_{\pi DD^*} = {}& 2i g^{(0)}_{WDD^*}(\partial_\mu D\pi\bar{D}^{*\mu} - D^{*\mu}\pi\partial_\mu \bar{D}) \\ &+ 2i g^{(1)}_{WDD^*}(D\partial_\mu \pi \bar{D}^{*\mu} - D^{*\mu}\partial_\mu \pi \bar{D}),\end{aligned} \tag{B19}$$

$$\mathcal{L}_{\psi D^*D^*} = 2i g^{(0)}_{\psi D^*D^*}\partial_\mu \psi_\nu(D^{*\mu}\bar{D}^{*\nu} - D^{*\nu}\bar{D}^{*\mu})$$

$$+2ig^{(1)}_{\psi D^*D^*}\psi_\mu(\partial^\mu D^*_\nu \bar{D}^{*\nu} - D^*_\nu\partial^\mu \bar{D}^{*\nu})$$
$$+2ig^{(2)}_{\psi D^*D^*}\psi_\mu(\partial_\nu D^{*\mu}\bar{D}^{*\nu} - D^{*\nu}\partial_\nu \bar{D}^{*\mu}), \quad \text{(B20)}$$

$$\mathcal{L}_{\rho D^*D^*} = 2ig^{(0)}_{AD^*D^*}(D^{*\nu}\partial_\mu\rho_\nu\bar{D}^{*\mu} - D^{*\mu}\partial_\mu\rho_\nu\bar{D}^{*\nu}) + 2ig^{(1)}_{AD^*D^*}(D^*_\nu\rho_\mu\partial^\mu\bar{D}^{*\nu} - \partial^\mu D^*_\nu\rho_\mu\bar{D}^{*\nu}) + 2ig^{(2)}_{AD^*D^*}(D^{*\mu}\rho_\nu\partial_\mu\bar{D}^{*\nu} - \partial_\mu D^{*\nu}\rho_\nu\bar{D}^{*\mu}), \quad \text{(B21)}$$

$$\mathcal{L}_{\psi DD} = 2ig_{\psi DD}\psi_\mu(\partial^\mu D\bar{D} - \partial^\mu D\bar{D}), \quad \text{(B22)}$$

$$\mathcal{L}_{\rho DD} = 2ig_{ADD}(D\rho^\mu\partial_\mu\bar{D} - \partial_\mu D\rho^\mu\bar{D}), \quad \text{(B23)}$$

$$\mathcal{L}_{\psi D_0^* D^*} = 2g_{W\psi DD^*}\sigma_0\psi^\mu(D_0^*\bar{D}^*_\mu + D^*_\mu\bar{D}^*_0), \quad \text{(B24)}$$

$$\mathcal{L}_{\psi DD_1} = 2ig_{W\psi DD^*}\sigma_0\psi_\mu\left(D^\mu_1\bar{D} - D\bar{D}^\mu_1\right), \quad \text{(B25)}$$

for three-point normal-parity interactions and

$$\mathcal{L}_{\pi\psi DD^*} = 2g_{W\psi DD^*}\psi^\mu(D\pi\bar{D}^*_\mu + D^*_\mu\pi\bar{D}), \quad \text{(B26)}$$

$$\mathcal{L}_{\rho\psi DD} = 2g_{A\psi DD}\psi^\mu D\rho_\mu\bar{D}, \quad \text{(B27)}$$

$$\mathcal{L}_{\rho\psi D^*D^*} = 2g^{(0)}_{A\psi D^*D^*}\psi^\mu D^{*\nu}\rho_\mu\bar{D}^*_\nu + 2g^{(1)}_{A\psi D^*D^*}\psi^\mu(D^*_\mu\rho^\nu\bar{D}^*_\nu + D^{*\nu}\rho_\nu\bar{D}^{*\nu}). \quad \text{(B28)}$$

for the four-point normal-parity interactions. The last two three-point interactions are induced from $\mathcal{L}_{W\psi DD^*}$. These play an essential role in showing the decoupling of the pion from the dissociation amplitude in the soft-momentum limit. As mentioned in Sec. IV, the coupling constant $g^{(0)}_{WDD^*}$ is set to zero to remove the mixing between the various D mesons. Furthermore, we drop the index on the remaining coupling constant $g^{(1)}_{WDD^*} \to g_{WDD^*}$. For the sake of making more transparent the correspondence with Ref. [45], we further set $g^{(0,2)}_{\{\psi,A\}D^*D^*} = -g^{(1)}_{\{\psi,A\}D^*D^*} \to g_{\{\psi,A\}D^*D^*}$ and $g^{(i)}_{A\psi D^*D^*} \to g_{A\psi D^*D^*}$.

Similarly, the abnormal parity content is

$$\mathcal{L}_{\pi D^*D^*} = -2g_{WD^*D^*}\epsilon^{\mu\nu\alpha\beta}\partial_\mu D^*_\nu\pi\partial_\alpha\bar{D}^*_\beta, \quad \text{(B29)}$$

$$\mathcal{L}_{\psi DD^*} = -2g_{\psi DD^*}\epsilon^{\mu\nu\alpha\beta}\partial_\mu\psi_\nu(\partial_\alpha D^*_\beta\bar{D} + D\partial_\alpha\bar{D}^*_\beta), \quad \text{(B30)}$$

$$\mathcal{L}_{\rho DD^*} = -2g_{ADD^*}\epsilon^{\mu\nu\alpha\beta}(\partial_\alpha D^*_\beta\partial_\mu\rho_\nu\bar{D} + D\partial_\mu\rho_\nu\partial_\alpha\bar{D}^*_\beta), \quad \text{(B31)}$$

$$\mathcal{L}_{\psi D^*D_1} = 2g_{W\psi D^*D_1}\sigma_0\epsilon^{\mu\nu\alpha\beta}\partial_\mu\psi_\nu(D_{1\alpha}\bar{D}^*_\beta - D^*_\alpha\bar{D}_{1\beta}), \quad \text{(B32)}$$

where the last interaction is generated by $\mathcal{L}_{WAD^*D^*}$ and

$$\mathcal{L}_{\pi\psi DD} = -2ig_{W\psi DD}\epsilon^{\mu\nu\alpha\beta}\psi_\mu\partial_\nu D\partial_\alpha\pi\partial_\beta\bar{D}, \quad \text{(B33)}$$

$$\mathcal{L}_{\pi\psi D^*D^*} = -2ig^{(0)}_{W\psi D^*D^*}\epsilon^{\mu\nu\alpha\beta}\psi_\mu D^*_\nu\partial_\alpha\pi\bar{D}^*_\beta - 2ig^{(1)}_{W\psi D^*D^*}\epsilon^{\mu\nu\alpha\beta}\partial_\mu\psi_\nu D^*_\alpha\pi\bar{D}^*_\beta, \quad \text{(B34)}$$

$$\mathcal{L}_{\rho\psi D^*D^*} = 2ig^{(0)}_{A\psi DD^*}\epsilon^{\mu\nu\alpha\beta}\psi_\mu(\partial_\nu D\rho_\alpha D^*_\beta + D^*_\nu\rho_\alpha\partial_\beta\bar{D}) - 2ig^{(1)}_{A\psi DD^*}\epsilon^{\mu\nu\alpha\beta}\psi_\mu(D\rho_\nu\partial_\alpha D^*_\beta - \partial_\nu D^*_\alpha\rho_\beta\bar{D}), \quad \text{(B35)}$$

for the three- and four-point interactions. This completes the list of all the required interactions.

APPENDIX C: DISSOCIATION AMPLITUDES

A. $\pi + J/\psi$

We first investigate the dissociation process into two D mesons illustrated in the first set of diagrams in Fig. 2. The subamplitudes are explicitly:

$$\mathcal{M}^\rho_{1a} = \frac{4g_{WDD^*}g_{\psi DD^*}}{t - m^{*2}_D}p^\alpha_\pi\epsilon^{p_\psi p_{\bar{D}}\beta\rho} \times \left(g_{\alpha\beta} - \frac{(p_\pi - p_D)_\alpha(p_\pi - p_D)_\beta}{m^{*2}_D}\right), \quad \text{(C1)}$$

$$\mathcal{M}^\rho_{1b} = -\frac{4g_{WDD^*}g_{\psi DD^*}}{u - m^{*2}_D}p^\alpha_\pi\epsilon^{p_\psi p_D\beta\rho} \times \left(g_{\alpha\beta} - \frac{(p_\pi - p_{\bar{D}})_\alpha(p_\pi - p_{\bar{D}})_\beta}{m^{*2}_D}\right), \quad \text{(C2)}$$

$$\mathcal{M}^\rho_{1c} = g_{W\psi DD}\epsilon^{p_\pi p_\psi p_{\bar{D}}\rho} \quad \text{(C3)}$$

where $t = (p_\pi - p_D)^2$ and $u = (p_\pi - p_{\bar{D}})^2$. Note that there are no additional diagrams compared to Ref. [45].

Next we consider the absorption process that has been considered dominant in the literature, namely $\pi + \psi \to \bar{D} + D^*$. As seen in Fig. 2, because of chiral symmetry, the number of subprocesses is higher than in a theory where the chiral partners are disregarded. Specifically, the list of subamplitudes for this process is

$$\mathcal{M}^{\mu\rho}_{2a} = -\frac{4g_{WDD^*}g_{\psi DD}}{t - m^2_D}p^\mu_\pi\left(2p^\rho_{\bar{D}} - p^\rho_\psi\right), \quad \text{(C4)}$$

$$\mathcal{M}^{\mu\rho}_{2b} = -\frac{4g_{WD^*D^*}g_{\psi DD^*}}{t - m^{*2}_D}\epsilon^{p_\psi p_{D^*}\mu\alpha}\epsilon^{p_\psi p_{\bar{D}}\beta\rho} \times \left(g_{\alpha\beta} - \frac{(p_\pi - p_{D^*})_\alpha(p_\pi - p_{D^*})_\beta}{m^{*2}_D}\right), \quad \text{(C5)}$$

$$\mathcal{M}^{\mu\rho}_{2c} = -\frac{4g_{WDD^*}g_{\psi D^*D^*}}{u - m^2_{D^*}}p^\alpha_\pi\left(2g^{\beta\rho}p^\mu_\psi - g^{\mu\rho}\left(p^\beta_\psi + p^\beta_{D^*}\right) + 2g^{\mu\beta}p^\rho_{D^*}\right)\left(g_{\alpha\beta} - \frac{(p_\pi - p_{\bar{D}})_\alpha(p_\pi - p_{\bar{D}})_\beta}{m^{*2}_D}\right), \quad \text{(C6)}$$

$$\mathcal{M}^{\mu\rho}_{2d} = g_{W\psi DD^*}g^{\mu\rho}, \quad \text{(C7)}$$

$$\mathcal{M}^{\mu\rho}_{2e} = -\frac{4\Delta^* g_{W\psi D^*D^*}\sigma_0}{t - m^2_{D_1}}g^{\mu\alpha}g^{\beta\rho} \times \left(g_{\alpha\beta} - \frac{(p_\pi - p_{D^*})_\alpha(p_\pi - p_{D^*})_\beta}{m^2_{D_1}}\right), \quad \text{(C8)}$$

$$\mathcal{M}^{\mu\rho}_{2f} = \frac{4\Delta g_{W\psi DD^*}\sigma_0}{u - m^2_{D^*_0}}g^{\mu\rho}, \quad \text{(C9)}$$

where $t = (p_\pi - p^*_D)^2$ and $u = (p_\pi - p_{\bar{D}})^2$. We note that the the last two amplitudes arise because of the exchange of the D_1 and D^*_0 mesons.

Finally, the last pion-absorption process is that which leads to the heaviest final state considered in this study, i.e., D^*-$\bar{D}^*$. The subamplitudes related to the diagrams in Fig. 2 are:

$$\mathcal{M}^{\mu\nu\rho}_{3a} = \frac{4g_{WDD^*}g_{\psi D^*D}}{t - m^2_D}p^\nu_\pi\epsilon^{p_\psi p_{\bar{D}^*}\mu\rho}, \quad \text{(C10)}$$

$$\mathcal{M}_{3b}^{\mu\rho} = -\frac{4g_{WDD^*}g_{\psi D^*D}}{u-m_D^2}p_\pi^\mu\epsilon^{p_\psi p_{D^*}\nu\rho}, \tag{C11}$$

$$\mathcal{M}_{3c}^{\mu\nu\rho} = \frac{4g_{WD^*D^*}g_{\psi D^*D^*}}{t-m_{D^*}^2}\epsilon^{p_\psi p_{D^*}\alpha\nu} \times\left(2g^{\beta\rho}p_\psi^\mu - g^{\mu\rho}\left(p_\psi^\beta + p_{\bar{D}^*}^\beta\right) + 2g^{\mu\beta}p_{\bar{D}^*}^\rho\right) \times\left(g_{\alpha\beta} - \frac{(p_\pi - p_{D^*})_\alpha(p_\pi - p_{D^*})_\beta}{m_D^{*2}}\right), \tag{C12}$$

$$\mathcal{M}_{3d}^{\mu\nu\rho} = \frac{4g_{WD^*D^*}g_{\psi D^*D^*}}{u-m_{D^*}^2}\epsilon^{p_\psi p_{\bar{D}^*}\alpha\mu} \times\left(2g^{\beta\rho}p_\psi^\nu - g^{\nu\rho}\left(p_\psi^\beta + p_{D^*}^\beta\right) + 2g^{\nu\beta}p_{D^*}^\rho\right) \times\left(g_{\alpha\beta} - \frac{(p_\pi - p_{\bar{D}^*})_\alpha(p_\pi - p_{\bar{D}^*})_\beta}{m_D^{*2}}\right), \tag{C13}$$

$$\mathcal{M}_{3e}^{\mu\nu\rho} = 2g_{W\psi D^*D^*}^{(0)}\epsilon^{p_\pi\mu\nu\rho} + 2g_{W\psi D^*D^*}^{(1)}\epsilon^{p_\psi\mu\nu\rho}, \tag{C14}$$

$$\mathcal{M}_{3f}^{\mu\rho} = -\frac{4\Delta^* g_{W\psi D^*D^*}\sigma_0}{t-m_{D_1}^2}g^{\alpha\nu}\epsilon^{p_\psi\mu\beta\rho} \times\left(g_{\alpha\beta} - \frac{(p_\pi - p_{D^*})_\alpha(p_\pi - p_{D^*})_\beta}{m_{D_1}^2}\right), \tag{C15}$$

$$\mathcal{M}_{3g}^{\mu\rho} = -\frac{4\Delta^* g_{W\psi D^*D^*}\sigma_0}{u-m_{D_1}^2}g^{\alpha\mu}\epsilon^{p_\psi\beta\nu\rho} \times\left(g_{\alpha\beta} - \frac{(p_\pi - p_{\bar{D}^*})_\alpha(p_\pi - p_{\bar{D}^*})_\beta}{m_{D_1}^2}\right) \tag{C16}$$

where $t=(p_\pi - p_{D^*})^2$ and $u=(p_\pi - p_{\bar{D}^*})^2$.

B. $\rho + J/\psi$

The amplitudes for the dissociation into the lowest mass state given in Fig. 6 are

$$\mathcal{M}_{4a}^{\delta\rho} = -\frac{4g_{ADD}g_{\psi DD}}{t-m_D^2}\left(2p_D^\delta - p_\rho^\delta\right)\left(2p_{\bar{D}}^\rho - p_\psi^\rho\right) \tag{C17}$$

$$\mathcal{M}_{4b}^{\delta\rho} = -\frac{4g_{ADD}g_{\psi DD}}{u-m_D^2}\left(2p_{\bar{D}}^\delta - p_\rho^\delta\right)\left(2p_D^\rho - p_\psi^\rho\right), \tag{C18}$$

$$\mathcal{M}_{4c}^{\delta\rho} = -\frac{4g_{ADD^*}g_{\psi DD^*}}{t-m_D^{*2}}\epsilon^{p_\rho p_D\alpha\delta}\epsilon^{p_\psi p_{\bar{D}}\beta\rho} \times\left(g_{\alpha\beta} - \frac{(p_\rho - p_{D^*})_\alpha(p_\rho - p_{D^*})_\beta}{m_D^{*2}}\right), \tag{C19}$$

$$\mathcal{M}_{4d}^{\delta\rho} = -\frac{4g_{ADD^*}g_{\psi DD^*}}{u-m_D^{*2}}\epsilon^{p_\rho p_{\bar{D}}\alpha\delta}\epsilon^{p_\psi p_D\beta\rho} \times\left(g_{\alpha\beta} - \frac{(p_\rho - p_{\bar{D}})_\alpha(p_\rho - p_{\bar{D}})_\beta}{m_D^{*2}}\right), \tag{C20}$$

$$\mathcal{M}_{4e}^{\delta\rho} = -2g_{A\psi DD}g^{\delta\rho}. \tag{C21}$$

where $t=(p_\rho - p_D)^2$ and $u=(p_\rho - p_{\bar{D}})^2$. Of all the six processes studied, this is the only one that is exothermic, i.e., the initial state is more massive than the final one. This kinematical constraint will give rise to a divergent cross section behavior at low $\sqrt{s}$.

The amplitudes of the second process (depicted in Fig. 6) are

$$\mathcal{M}_{5a}^{\mu\delta\rho} = \frac{4g_{ADD^*}g_{\psi DD}}{t-m_D^2}\epsilon^{p_\rho p_{D^*}\mu\delta}\left(2p_{\bar{D}}^\rho - p_\psi^\rho\right), \tag{C22}$$

$$\mathcal{M}_{5b}^{\mu\delta\rho} = \frac{4g_{ADD^*}g_{\psi DD}}{u-m_D^2}\left(2p_{\bar{D}}^\delta - p_\rho^\delta\right)\epsilon^{p_\psi p_{D^*}\mu\rho}, \tag{C23}$$

$$\mathcal{M}_{5c}^{\mu\delta\rho} = \frac{4g_{AD^*D^*}g_{\psi DD^*}}{t-m_{D^*}^2} \times\left(2g^{\alpha\delta}p_\rho^\mu - g^{\mu\delta}\left(p_\rho^\alpha + p_{D^*}^\alpha\right) + 2g^{\alpha\mu}p_{D^*}^\delta\right) \times\left(g_{\alpha\beta} - \frac{(p_\rho - p_{D^*})_\alpha(p_\rho - p_{D^*})_\beta}{m_D^{*2}}\right)\epsilon^{p_\psi p_{\bar{D}}\beta\rho}, \tag{C24}$$

$$\mathcal{M}_{5d}^{\mu\delta\rho} = \frac{4g_{ADD^*}g_{\psi D^*D^*}}{u-m_{D^*}^2}\epsilon^{p_\rho p_{\bar{D}}\alpha\delta} \times\left(g_{\alpha\beta} - \frac{(p_\rho - p_{\bar{D}})_\alpha(p_\rho - p_{\bar{D}})_\beta}{m_D^{*2}}\right) \times\left(2g^{\beta\rho}p_\psi^\mu - g^{\mu\rho}\left(p_\psi^\beta + p_{D^*}^\beta\right) + 2g^{\mu\beta}p_{D^*}^\rho\right), \tag{C25}$$

$$\mathcal{M}_{5e}^{\mu\delta\rho} = 2g_{A\psi DD^*}^{(0)}\epsilon^{p_{\bar{D}}\mu\delta\rho} + 2g_{A\psi DD^*}^{(1)}\epsilon^{p_{D^*}\mu\delta\rho}, \tag{C26}$$

where $t=(p_\rho - p_{D^*})^2$ and $u=(p_\rho - p_{\bar{D}})^2$.

And, finally, the set of amplitudes for the final dissociation processes, given in Fig. 6, have the corresponding expressions:

$$\mathcal{M}_{6a}^{\mu\nu\delta\rho} = -\frac{4g_{ADD^*}g_{\psi DD^*}}{t-m_D^2}\epsilon^{p_\rho p_{D^*}\mu\delta}\epsilon^{p_\psi p_{\bar{D}^*}\nu\rho}, \tag{C27}$$

$$\mathcal{M}_{6b}^{\mu\nu\delta\rho} = -\frac{4g_{ADD^*}g_{\psi DD^*}}{u-m_D^2}\epsilon^{p_\rho p_{\bar{D}^*}\nu\delta}\epsilon^{p_\psi p_{D^*}\mu\rho}, \tag{C28}$$

$$\mathcal{M}_{6c}^{\mu\nu\delta\rho} = -\frac{4g_{AD^*D^*}g_{\psi D^*D^*}}{t-m_{D^*}^2} \times\left(g_{\alpha\beta} - \frac{(p_\rho - p_{D^*})_\alpha(p_\rho - p_{D^*})_\beta}{m_D^{*2}}\right) \times\left(2g^{\alpha\delta}p_\rho^\mu - g^{\mu\delta}\left(p_\rho^\alpha + p_{D^*}^\alpha\right) + 2g^{\alpha\mu}p_{D^*}^\rho\right) \times\left(2g^{\beta\rho}p_\rho^\nu - g^{\nu\rho}\left(p_\rho^\beta + p_{\bar{D}^*}^\beta\right) + 2g^{\beta\nu}p_{\bar{D}^*}^\rho\right), \tag{C29}$$

$$\mathcal{M}_{6d}^{\mu\nu\delta\rho} = -\frac{4g_{AD^*D^*}g_{\psi D^*D^*}}{t-m_{D^*}^2} \times\left(g_{\alpha\beta} - \frac{(p_\rho - p_{D^*})_\alpha(p_\rho - p_{D^*})_\beta}{m_D^{*2}}\right) \times\left(2g^{\alpha\delta}p_\rho^\nu - g^{\nu\delta}\left(p_\rho^\alpha + p_{\bar{D}^*}^\alpha\right) + 2g^{\alpha\nu}p_{\bar{D}^*}^\rho\right) \times\left(2g^{\beta\rho}p_\rho^\nu - g^{\nu\rho}\left(p_\rho^\beta + p_{D^*}^\beta\right) + 2g^{\beta\nu}p_{D^*}^\rho\right), \tag{C30}$$

$$\mathcal{M}_{6e}^{\mu\nu\delta\rho} = g_{A\psi D^*D^*}^{(0)}\left(2g^{\mu\nu}g^{\delta\rho} - g^{\mu\delta}g^{\nu\rho} - g^{\mu\rho}g^{\nu\delta}\right) \tag{C31}$$

where $t=(p_\rho - p_{D^*})^2$ and $u=(p_\rho - p_{\bar{D}^*})^2$.

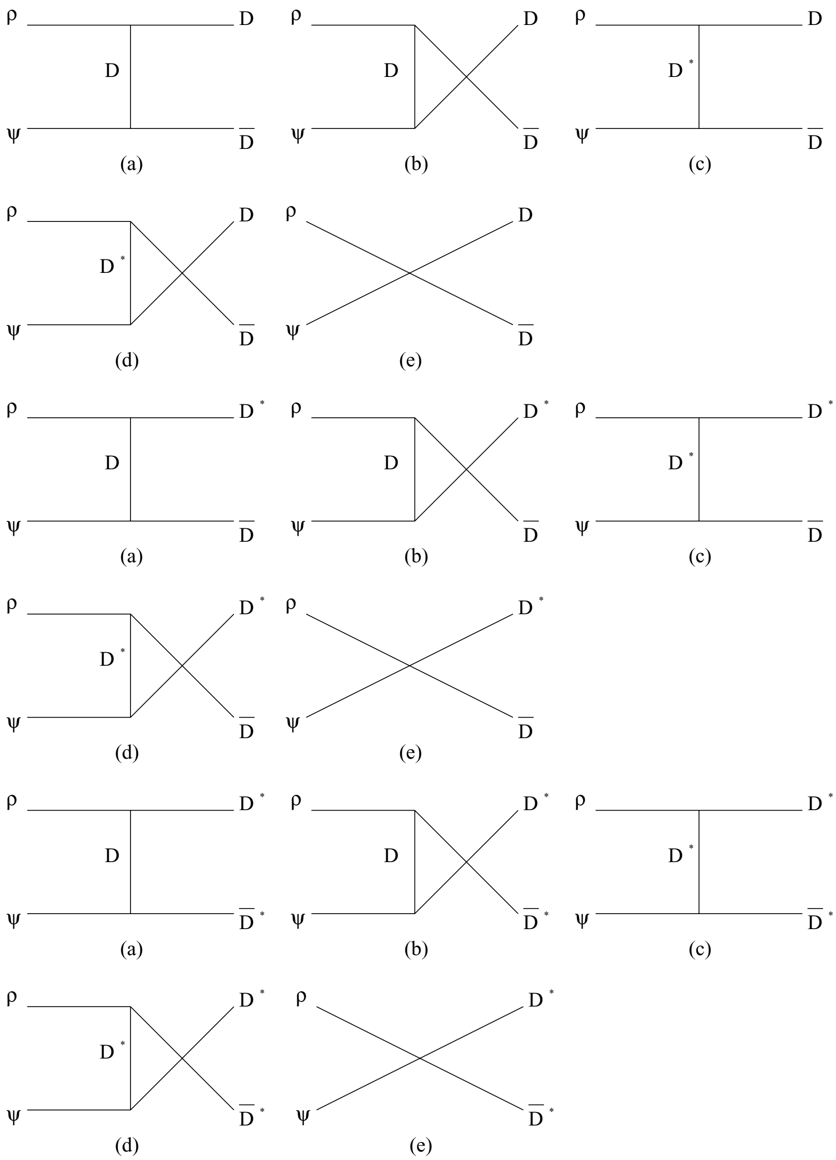

FIG. 6. Diagrams for $\rho + \psi$ dissociation.

APPENDIX D: PARAMETER FIXING

The coupling constants used here are fixed to those of Ref. [45]. There, in addition to fitting the available experimental data, they invoked the vector meson dominance hypothesis, the heavy quark spin-flavor symmetry, and the underlying SU(4) symmetry on which the Lagrangian is built. Each of these assumptions is problematic. Unfortunately, because experimental data are lacking to fix, for example, the four-point couplings, the only other way would be to use other model calculations with varying degrees of sophistication. Table II lists the coupling constant values used.

TABLE II. Coupling constants of the phenomenological Lagrangian.

Three-point couplings		Four-point couplings	
g_{WDD^*}	4.40	$g_{\psi DD^*}$	16.96
$g_{\psi D^*D^*}$	3.86	$g_{A\psi DD}$	19.43
$g_{\psi DD}$	3.86	$g_{A\psi D^*D^*}$	9.72
$g_{AD^*D^*}$	1.26	$g_{W\psi DD}$	8.00 GeV^{-3}
g_{ADD}	1.26	$g^{(i)}_{W\psi D^*D^*}$	19.10 GeV^{-1}
$g_{WD^*D^*}$	4.54 GeV^{-1}	$g^{(i)}_{A\psi DD^*}$	10.89 GeV^{-1}
$g_{\psi DD^*}$	4.32 GeV^{-1}		
g_{ADD^*}	1.41 GeV^{-1}		

Setting the coupling constants to those of Ref. [45] is not sufficient to determine all the parameters. Five parameters: M, M^*, Δ, Δ^*, and σ_0 have to be determined. The last one is the decay constant, $f_\pi = 93$ MeV. The four remaining

TABLE III. Cutoff-dependent coupling constants.

	Δ (GeV)	Δ^* (GeV)	γ_D	γ_{D^*}
No form factors	6.10	5.01	1	1
$\Lambda = 1$ GeV	19.85	14.36	0.79	0.78
$\Lambda = 2$ GeV	9.53	7.35	0.94	0.94

parameters have to reproduce the masses [3] of the D, D^*, D_0^*, and D_1 mesons, namely $m_D = 1.87$ GeV, $m_{D^*} = 2.01$ GeV, $m_{D_0^*} = 2.40$ GeV, and $m_{D_1} = 2.43$ GeV, respectively [55]. This leads to values of $M = 2.15$ GeV and $M^* = 2.23$ GeV. Table III lists the values of Δ and Δ^* and γ_D and γ_{D^*} used. Finally, the pion, ρ, and J/ψ masses are taken to be 0.138 GeV, 0.770 GeV, and 3.10 GeV, respectively.

[3]For the D and D^* the isopin averaged masses are used.

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