We present a new calculation of the $\pi$--$J/\psi$ dissociation cross sections within the Constituent Quark–Meson Model recently introduced. To discuss the absorption of $J/\psi$ in heavy-ion collisions, we assume the $J/\psi$ to be produced inside a thermalized pion gas, as discussed by Bjorken, and introduce the corrections due to absorption by nuclear matter as well. We fit the absorption length of the $J/\psi$ to the data obtained at the CERN SPS by the NA50 Collaboration for Pb-Pb collisions. Collisions of lower centrality allow us to determine the temperature and the energy density of the pion gas. For both these quantities we find values close to those indicated by lattice gauge calculations for the transition to a quark–gluon plasma. A simple extrapolation to more central collisions, which takes into account the increase of the energy deposited due to the increased nucleon flux, fails to reproduce the break in $J/\psi$ absorption indicated by NA50, thus lending support to the idea that an unconfined quark–gluon phase may have been produced. This conclusion could be sharpened by analysing in a similar way, as a function of centrality, other observables such as strange particle production.

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I. INTRODUCTION

In a recent paper the couplings of $J/\psi$ to $D^{(*)}$ mesons (where by $D^{(*)}$ we mean $D$ or $D^*$) and to $\pi$'s have been computed in the framework of a Constituent Quark–Meson Model introduced in \[\text{Ref. \cite{Maiani}}\]. The interest of this study stems from the possibility that $J/\psi$ absorption processes of the type:

$$\pi + J/\psi \rightarrow D^{(*)} + \bar{D}^{(*)},$$

play an important role in the relativistic heavy-ion scattering. Since a decrease in the $J/\psi$ production in these processes might signal the formation of Quark–Gluon Plasma (QGP), it is useful to have reliable estimates of the cross sections for the processes, which provide an alternative way to reduce the $J/\psi$ production rate. Previous studies of these effects can be found in \[\text{Ref. \cite{Stodolsky}}\]. In this paper we will address the problem of understanding how far we can go in explaining $J/\psi$ suppression data via an hadronic mechanism such as \[\text{Ref. \cite{Stodolsky}}\]. Does the hot gas of pions formed after the heavy-ion collision provide a source of attenuation of $J/\psi$ antagonist to the standard QGP suppression?

The relevant couplings needed to compute the cross section of process \[\text{Ref. \cite{Maiani}}\] are shown in Fig. 1. To compute the amplitudes \[\text{Ref. \cite{Maiani}}\], besides the $DD^*\pi$ coupling, Fig.\[\text{Ref. \cite{Maiani}}\], whose strength $g_{DD^*\pi}$ has been theoretically estimated and experimentally investigated \[\text{Ref. \cite{Stodolsky}}\], one would need also the $JD^{(*)}D^{(*)}$, Fig.\[\text{Ref. \cite{Maiani}}\], and the $JD^{(*)}D^{(*)}\pi$ couplings, Fig.\[\text{Ref. \cite{Maiani}}\]. In an effective Lagrangian approach the latter couplings provide direct four-body amplitudes, while the former enter the amplitude via tree diagrams with an exchange of a charmed particle $D^{(*)}$ in the $t$–channel.
These couplings have also been estimated by different methods, all presenting, in our opinion, different shortcomings. For example the use of the $SU(4)$ symmetry puts on the same footing the heavy quark $c$ and the light quarks, which is at odds with the results obtained within the Heavy Quark Effective Theory, where the opposite approximation $m_c \gg \Lambda_{QCD}$ is used. Similarly, the rather common approach based on Vector Meson Dominance should be considered critically, given the large extrapolation from $\rho^2 = 0$ to $\rho^2 = m_{J/\psi}^2$ that is involved. A different evaluation, based on QCD Sum Rules can be found in $[9]$ and presents the typical theoretical uncertainties of this method. It is useful to have diverse, although inevitably model-dependent, calculations to assess a reliable range of theoretical results. A comparison of our results with the cross sections obtained by other methods is given below, so to provide a further indication of the theoretical uncertainties (the theoretical uncertainties intrinsic to the model used in this paper have been discussed in $[9]$).

To apply our results to $J/\psi$ production in heavy-ion collisions, we follow the picture introduced by Bjorken some time ago $[7]$. Seen in the centre-of-mass frame, the colliding nuclei appear as Lorentz-contracted pancakes which traverse one another in a short time $t$ (say, 2 fm/c). After crossing, the nuclei leave behind them a fireball with essentially vanishing baryon number, where the produced $J/\psi$ is sitting (particles in the fireball are often referred to as comovers).

For low-centrality collisions, the fireball should be, to a good approximation, a pion gas that quickly goes to thermal equilibrium with a temperature corresponding to the energy deposited by the colliding nucleons in the central region. The energy density estimated by Bjorken is in the range of 0.3–3 GeV/fm$^3$ and is proportional to the number of nucleons per unit area that participate in the collision.

Our cross section calculation allows us to compute the absorption length of the $J/\psi$ in the pion gas, which we tentatively approximate with a perfect gas with vanishing chemical potential (at high temperature, say $T \gg 150$ MeV, there are enough inelastic collisions in the pion gas to provide chemical equilibrium as well).

To obtain a more accurate value of $T$, we introduce also the corrections due to $J/\psi$ absorption by nuclear matter, using the absorption cross sections obtained recently from $p-A$ collisions $[8]$.

We have fitted the absorption length to the data obtained by the NA50 collaboration at the SPS, which show a clear exponential behaviour as a function of the linear dimension of the collision region, $l$, up to $l \approx 5$ fm. We find not unreasonable values for the temperature ($T = 225$ MeV) and for the energy density ($\epsilon \approx 0.32$ GeV/fm$^3$). The absorption length decreases very steeply with increasing temperature.

The energy density is perhaps on the low side with respect to $a$ priori expectations based on the Bjorken formula. On the other hand, the temperature is somewhat higher than the critical temperature computed in lattice calculations (see $[8]$ for a recent review). The difference, if any, could be due to either an underestimate of the cross section, or to the perfect gas approximation being too crude, or to both. Large departures from perfect gas behaviour are found in numerical simulations below the critical temperature $[8]$. Further work to improve on the perfect Bose gas approximation and obtain a better calibration of the crucial parameters of the fireball, temperature and energy density, is needed.

To extrapolate to higher centrality, we keep into account the increase of the average nucleon number per unit area for decreasing impact parameter. This, in turn, makes the temperature and the deposited energy increase while the absorption length decreases, leading to a downward bending of the $J/\psi$ production as a function of $l$. The behaviour we find, however, is too smooth to reproduce the rather sharp break shown by the NA50 data for $l$ larger than 5 fm. This suggests that the simple pion gas description ceases to be valid at these values of $l$, thus lending support to the formation of a new phase. Our analysis implies that it would be useful to analyse in the same way, i.e. as a function of $l$, the data obtained for other observables such as the production of strange particles. Given the rather low temperatures that we find in the low-centrality region, we would expect very few strange particles to be produced.

![FIG. 1: Couplings involved in the tree–level calculations of processes $\pi$.](image-url)
in the central region as well. This would not apply to the unconfined phase, where the Boltzmann suppression of strange quarks with \( m_s = 150 \text{ MeV} \) would not be operative for temperatures around 225 MeV.

II. CROSS SECTIONS

The calculation of the total cross sections \( \sigma(\pi J/\psi \to D^{(*)}D^{(*)}) \) proceeds through the evaluation of the amplitudes listed in the Appendix and the two-body phase-space integration:

\[
\sigma_{\pi J/\psi \to D\bar{D}}(s) = \frac{M_{J/\psi}}{12\pi(s - M_{J/\psi}^2)} \int_{R} dp_D \frac{p_D}{E_D} \sum_{\text{pol}} |A(s, p_D)|^2 ,
\]

where \( s \) is the Mandelstam variable \( s = (p_\pi + p_{J/\psi})^2 \), and \( R \) is the phase-space interval \( R = (p_{D}^{\text{min}}, p_{D}^{\text{max}}) \). An overall factor of 2 is included to count all isospin quantum numbers in the final state. The laboratory frame is considered with high-momentum massless pions colliding on \( J/\psi \) at rest. En energy threshold of about \( E_\pi = 800 \text{ MeV} \) is required to open the reaction channel.

For example, in the case of the \( D\bar{D} \) final state, using the notation introduced in [1], we have \( A = A_{1a} + A_{2a} + A_{3a} \), where:

\[
\begin{align*}
A_{1a} &= i \frac{g_{\pi D D^*} g_{J/\psi D D^*}}{t - M_{D^*}^2} \epsilon(\eta, p_4, p_3 - p_2, \alpha) \Pi^{\mu \nu}(p_3 - p_2)(p_2)_{\nu} \\
A_{1b} &= i \frac{g_{\pi DD^*} g_{J/\psi DD^*}}{u - M_{D^*}^2} \epsilon(\eta, p_3, p_1 - p_3, \alpha) \Pi^{\mu \nu}(p_3 - p_1)(p_2)_{\nu} \\
A_{1c} &= i \frac{g_{J D D^*}}{M_D} \epsilon(\eta, p_2, p_4, p_3),
\end{align*}
\]

where \( \eta \) is the polarization vector of \( J/\psi \), \( \Pi^{\mu \nu}(q) \) is the sum-over-polarization-tensor

\[
\Pi^{\mu \nu}(q) = \frac{q^{\mu} q^{\nu}}{M^2} - g^{\mu \nu},
\]

\( M \) being the mass of the vector particle. The notation \( \epsilon(\mu, \nu, p, q) = \epsilon_{\mu \nu \sigma \rho} p^\sigma q^\rho \) is used. The coupling constants \( g_{\pi D D^*}, g_{J/\psi D D^*}, g_{J D D^*} \) have been discussed in [1] (the notation \( g_{J D D^*} = g_0 \) is used), and are:

\[
\begin{align*}
g_{\pi D D^*} &= \frac{13}{1 + \frac{q_{\pi} p_{D^*}}{A_1^2}} \\
g_{J/\psi D D^*} &= 4.05 \text{ (GeV)}^{-1} \\
g_{J D D^*} &= \frac{234}{1 + \frac{q_{\pi} p_{D^*}}{A_2^2}} \text{ (GeV)}^{-2},
\end{align*}
\]

where \( A_1 = 1.1 \text{ GeV} \) and \( A_2 = 0.8 \text{ GeV} \). Attempts to compute quadrilinear couplings using \( SU(4) \) symmetry can be found in [2]. The dependence of the cross sections on \( \sqrt{s} \) for the \( DD \) and \( DD^* \) final states are shown in Fig. 2.

The amplitudes for \( DD^*, D^* D^* \) final states are given in the Appendix. The results we obtain can be compared with those shown in Fig. 2 extracted by the authors of [1]. Note the sharp rise of \( DD^* \) and \( D^* D^* \), which is due to \( S \)-wave production.

III. THERMAL AVERAGE

In a first instance we neglect \( J/\psi \) absorption by the nuclear matter, and focus on the effect of the fireball left over after the nuclei have passed each other (i.e. the \textit{comoving particles} in the frame of the target nucleus). The quantity we are interested in is the \( J/\psi \) suppression \( \mathcal{A} \) as a function of the linear size \( l \) of the fireball. A \( J/\psi \) produced in any internal point of a fireball of diameter \( l \) has to travel on average a distance \( \approx 6/10l \) to reach the boundaries of the ball. On the other hand the mean free path of the \( J/\psi \) is \( \lambda \approx 1/\rho \sigma \), so the attenuation can be defined as:

\[
\mathcal{A}(x) = \exp \left[-x \langle \rho \cdot \sigma_{\pi J/\psi \rightarrow D^{(*)}D^{(*)}}T \rangle \right],
\]
FIG. 2: Cross sections $\sigma_{\pi J/\psi \rightarrow (DD, DD^*, D^*D^*)}$ as functions of $\sqrt{s}$. We show the error band for the $DD$ final state resulting from the theoretical uncertainty in the determination of the couplings [1]. The error bands on $DD^*$ and $D^*D^*$ are similar.

FIG. 3: Cross sections computed with various approaches: QCD sum rules (band), short-distance QCD (dotted line), meson-exchange models (dot-dashed line), non-relativistic constituent quark model (dashed line) [10].

with $x \approx 6/10l$.

As a simple approximation, the thermal average $\langle \ldots \rangle_T$ is taken in a perfect Bose gas of charged and neutral pions with vanishing chemical potential:

$$\langle \rho \cdot \sigma_{\pi J/\psi \rightarrow D(\ast)D(\ast)} \rangle_T = \frac{3}{2\pi^2} \int_{E_{\pi}^\ast}^{\infty} dE_\pi \frac{E_\pi^2 \sigma(E_\pi)}{e^{E_\pi/kT} - 1}.$$  \hspace{1cm} (4)

For reference, we give also the number density of the relevant pions at temperature $T$

$$\rho(T) = \frac{3}{2\pi^2} \int_{E_{\pi}^\ast}^{\infty} dE_\pi \frac{E_\pi^2}{e^{E_\pi/kT} - 1}.$$  \hspace{1cm} (5)
where $E_{th}^\pi$ is the threshold energy required to open the reaction channel, and the total energy density:

$$\epsilon(T) = \frac{3}{2\pi^2} \int_{m_\pi}^\infty dE_\pi \frac{E_\pi^3}{e^{E_\pi/kT} - 1}. \quad (6)$$

We report in Table I the values of $\rho(T)$, $\epsilon(T)$ and $\langle \rho \cdot \sigma \rangle_T$ for different values of the temperature.

| $T$ (MeV) | $\rho$ (fm$^{-3}$) | $\epsilon$ (T) (MeV/fm$^3$) | $\langle \rho \cdot \sigma \rangle$ (fm$^{-2}$) |
|----------|--------------------|-----------------------------|-----------------------------|
| 205      | 0.0467             | 215                         | 0.037                       |
| 215      | 0.0629             | 261                         | 0.051                       |
| 225      | 0.0828             | 313                         | 0.070                       |
| 235      | 0.11               | 373                         | 0.093                       |
| 245      | 0.13               | 441                         | 0.12                        |
| 255      | 0.17               | 518                         | 0.16                        |
| 265      | 0.20               | 605                         | 0.20                        |
| 275      | 0.25               | 702                         | 0.25                        |

TABLE I: Values of $\rho(T)$, $\epsilon(T)$ and $\langle \rho \cdot \sigma \rangle_T$ for different values of the temperature.

FIG. 4: The overlap region of two colliding equal nuclei. Dimensions $l$ and $L$ determine the transverse size of the pion fireball and the longitudinal size of the nucleon column to be traversed by the $J/\psi$, respectively.

To connect to experimental data, we observe that, in the approximation of sharp-edged nuclei, the linear size of the fireball, $l$, is given by the transverse dimension of the region where nuclei overlap (see Fig. 4), which in turn is related to the impact parameter $b$ according to:

$$l = 2R - b. \quad (7)$$

The NA50 Collaboration has studied the attenuation of $J/\psi$ by studying the ratio:

$$A_{\text{expt}} = \frac{B_{\mu \mu} \sigma(J/\psi)}{\sigma(Drell-Yan)}. \quad (8)$$

as a function of the impact parameter $b$. Using Eq. (7), we can fit our Eq. (8) to the NA50 data, plotted as a function of $l$. The result is shown in Fig. 5 for different values of the fireball temperature. A good fit can be achieved in the region of lower centrality, up to $l \approx 5$ fm, thereby determining the fireball’s temperature. Observe that peripheral interactions agree with extrapolation from p- or d-induced reactions (the points at $l \approx 0$). From Fig. 5 we find $T \approx 255$ MeV.

To be more precise on the temperature, however, we must consider the effect of nuclear absorption, due to the nucleon column density, which the $J/\psi$ has to traverse during the initial phase of the collision. Neglecting the $J/\psi$ transverse momentum, the nuclear absorption factor is:

$$A_{\text{nuc.abs.}} = \exp[-L \cdot \rho_{\text{nuc.}} \cdot \sigma_{\text{nuc}}] \quad (9)$$
L is approximately equal to the longitudinal size of the overlap region of the colliding nuclei (Fig. 4). To avoid unwanted distortions, we have interpolated the values of $L$ given by NA50 as functions of the impact parameter $b$ to obtain $L$ as a function of our length $l$ via Eq. (7). A determination of the nuclear absorption cross section has been obtained by NA50 from the study of $J/\psi$ production in $p-A$ collisions. They find $\sigma_{\nucl} = 4.3 \pm 0.6$ mb

This cross section is smaller than what would explain the data in the low-centrality region with nuclear absorption only (this requires $\sigma_{\nucl} \approx 6$ mb [10]) and points to the importance of dissociation induced by the colliding particles, which in fact turns out to dominate. In conclusion, we compare the experimental quantity Eq. (8) to the expression:

$$A_{\nucl}(l) = C \cdot \exp[-L(l) \cdot \rho_{\nucl} \cdot \sigma_{\nucl}] \cdot \exp[-6/10 \cdot l \cdot (\rho \cdot \sigma)_{T}],$$

where $C$ is a normalization factor, $L$ and $l$ are obtained from the impact parameter $b$ following the NA50 prescription and Eq. (7), respectively. We use $\rho_{\nucl} = 0.17$ fm$^{-1}$, $\sigma_{\nucl} = 4.3$ mb.

The results are shown in Fig. 6. The fireball temperature that gives the best fit to the low-centrality data, up to $l \approx 5$ fm, is now $T = 225$ MeV, and the corresponding energy density is $\epsilon \approx 0.32$ GeV-fm$^{-3}$. We have studied how the theoretical errors in the coupling constants determined in [11] affect the determination of $T$. We find $T = 225 \pm 15$ MeV.

Going towards higher centrality, we expect an increase in the energy density of the fireball. One effect that is easy to take into account is the increase in the surface density of the nucleons participating in the collision. The energy density of the fireball is, in fact, proportional to the factor $\tilde{\sigma}$:

$$\tilde{\sigma} = \frac{\rho_{\nucl} \cdot V(b)}{S(b)} = A \cdot g(b/R),$$

$$g(b/R) = \frac{2}{3} \pi \left( \frac{1 - b/2R}{\arccos(b/2R) - (b/2R)\sqrt{1 - b^2/4R^2}} \right),$$

with $R = A^{1/3} r_0$, the nucleon radius $r_0 = 1.1$ fm [11].

Starting from $l_0 \approx 5$ fm and $T_0 = 225$ MeV, the geometrical factor $g(2 - l/R)$ is used to extrapolate the variation of temperature with $l$ according to:

$$T(l) = T_0 \left( \frac{g(2 - l/R)}{g(2 - l_0/R)} \right)^{1/4}. \quad (14)$$

Once we know $T(l)$ we compute the corresponding attenuation function defined as:

$$\tilde{A}(l) = C \cdot \exp[-L(l)\rho_{\nucl} \cdot \sigma_{\nucl}] \exp[-6/10 \cdot l \cdot (\rho \cdot \sigma)_{T(l)}]. \quad (15)$$

Fig. 4 shows the final best prediction of the conventional absorption effects due to the pion gas fireball and to the nuclear matter (solid curve).

Fitting the data of low-centrality, we are able to determine with some accuracy the temperature of the fireball, assumed to be an ideal pion gas. The temperature is remarkably close to, perhaps somewhat higher than, the critical temperature computed in lattice calculations [4], $T_c \approx 180$ MeV. We recall that in SU(3) gauge theory, $T_c = 260$ MeV while lattice results with $n_f = 2$ suggest $T_c = 170$ MeV. If the same high temperature limit is assumed for $n_f > 0$ and SU(3), then for $n_f = 2$ or 3, the $\psi'$ breaks at 190–200 MeV and the $J/\psi$ itself would not break up until $T > 2T_c$ [12]. Departures of the estimated temperature from the real one could be due either to an underestimation of the cross section, or to the inadequacy of the perfect gas approximation, or to both. Large deviations from the perfect gas are found in numerical simulations below the critical temperature [4]. Further work to improve on the perfect Bose gas approximation and to obtain a better calibration of the crucial parameters of the fireball, temperature and energy density, is needed.

The energy density can be compared with the formula given in ref. [4]:

$$\epsilon = \frac{A(b)}{S(b)} \frac{dE}{dy} \frac{1}{ct}, \quad (16)$$

$A(b)$ and $S(b)$ are the number of nucleons that participate in the collision and the overlap area, respectively, as functions of the impact parameter $b$; $dE/dy$ is the energy deposited per unit rapidity and $ct$, the longitudinal length.
of the fireball, is related to the time \( t \) it takes for the nuclei to separate completely. The energy deposited per unit rapidity is estimated as follows \( \text{(17)} \):

\[
\frac{dE}{dy} = \frac{dN_{ch}}{dy} \left( 1 + \frac{N_{\text{neutr}}}{N_{ch}} \right) \langle E \rangle \simeq 3 \times 1.5 \times 400 \text{ MeV} = 1.8 \text{ GeV} .
\]

For central Pb-Pb collisions, one has:

\[
\frac{A(b = 0)}{S(b = 0)} = \frac{A}{\pi R^2} = \frac{A^{1/3}}{\pi r_0^2} = 1.5 \text{ fm}^{-2} , \text{(18)}
\]

For Pb and \( L = 5 \text{ fm} \) we find \( g = 0.7 \). In conclusion, for Pb-Pb collisions we find:

\[
\epsilon = 0.95 \text{ GeV/fm}^3 \cdot \left( \frac{2 \text{ fm}}{ct} \right) \quad (L = 5 \text{ fm}) \text{(19)}
\]

\[
\epsilon = 1.35 \text{ GeV/fm}^3 \cdot \left( \frac{2 \text{ fm}}{ct} \right) \quad (L = 2R = 13 \text{ fm}) . \text{(20)}
\]

**FIG. 5:** The attenuation function \( A \) given in Eq. (3) as a function the linear size of the fireball. Lines correspond to a fixed temperature of the \( \pi \) gas expressed in MeV (isotherms). Experimental data: diamonds= \( p-(p,d) \) (NA51), boxes=Pb-Pb (NA50). The value of \( l \) for each experimental point is obtained from the measured value of the impact parameter \( b \) via Eq. (7). Lower-centrality interactions agree with extrapolation from \( p \)-induced reactions.

### IV. CONCLUSIONS AND OUTLOOK

Our final results are displayed in Figs. 6 and 7. In the lower-centrality region, where we expect to be still in the deconfined phase, we find \( T = 225 \pm 15 \text{ MeV} \) corresponding to the energy density \( \epsilon \simeq 0.32 \text{ GeV/fm}^3 \). The value of the energy density is not inconsistent with the Bjorken formula. The temperature, however, is somewhat higher than the critical temperature computed in lattice calculations.

The extrapolation to higher centrality keeps into account only the increase of the average nucleon number per unit area (an admittedly crude approximation). This leads to a downward bending of the \( J/\psi \) production as a function of \( l \). However, the behaviour is too smooth to reproduce the rather sharp break shown by the NA50 data.

Our analysis suggests that the simple pion gas description ceases to be valid at higher values of \( l \), thus lending some support to the formation of a new phase. To confirm or disprove this, it would be crucial to analyse in the same way, i.e. as a function of \( l \), the data obtained at the SPS and RHIC for other observables, such as strange particles production. In a pion gas, given the relatively low temperatures that we find in the low-centrality region, we would expect very few strange particles to be produced in the central region. This would not apply to the deconfined phase, where the Boltzmann suppression of strange quarks with \( m_s = 150 \text{ MeV} \) would not be operative for temperatures around 225 MeV.
FIG. 6: The attenuation function $A$ given in Eq. (11) as a function of the linear size of the fireball, computed with $\rho_{\text{nucl}} = 0.17 \text{ fm}^{-3}$ and $\sigma_{\text{nucl}} = 4.3 \text{ mb}$. Lines correspond to a fixed temperature of the $\pi$ gas expressed in MeV. Experimental data: Pb-Pb collisions from NA50 [8]. Diamonds represent $p-(p,d)$ (NA51) data points [13].

FIG. 7: Complete calculation of the attenuation function. $T = 225 \text{ MeV}$, $\rho_{\text{nucl}} = 0.17 \text{ fm}^{-3}$ and $\sigma_{\text{nucl}} = 4.3 \text{ mb}$. The solid line includes the effect of the increase of nucleon density per unit surface with decreasing impact parameter. The dashed line, given for reference, is the same as in Fig. 6. Data from NA50 Pb-Pb collisions [8].

Appendix

We report the amplitudes corresponding to the tree–level diagrams for the $DD^*$ and $D^*D^*$ final states. The diagrams and the corresponding amplitudes are as follows (dashed lines for spinless particles, continuous lines for spin-1 particles); we use the same notation as in [1].
In the case in which the final state is $DD^*$:

\[
A_{2a} = \frac{2g_\pi DD \cdot g_\Gamma \phi DD}{t - M_B^2} (p_2 \cdot (\epsilon(p_3))(p_4 \cdot (\epsilon(p_1)))
\]

\[
A_{2b} = \frac{g_\pi D^* D^* \cdot g_\Gamma \phi DD}{M_D(t - M_D^2)} (\epsilon(p_1), p_4, p_1, p_2, \rho) \times (\epsilon(p_3), p_2, p_3, \nu) \Pi^\nu (p_3 - p_2)
\]

\[
A_{2c} = \frac{g_\pi DD \cdot g_\Gamma \phi D^* D^*}{u - M_B^2} p_{2\rho} ((p_3 - (p_1 \cdot (\epsilon(p_1)))\epsilon_\mu (p_3) + (\epsilon(p_1) \cdot (\epsilon(p_3)))\epsilon_\mu (p_3) + (\epsilon(p_1) \cdot (p_3))\epsilon_\mu (p_3)) \times \Pi^\nu (p_3 - p_1, M_B^2) + ((p_3 - (p_1) \cdot (\epsilon(p_3)))\epsilon_\mu (p_3) + (\epsilon(p_1) \cdot (p_3))\epsilon_\mu (p_3)) \times \Pi^\nu (p_3 - p_1, M_B^2)
\]

\[
A_{2d} = -p_2^\nu \epsilon_\mu (p_1) (g_3 M_D^{-2}(p_{3\nu} p_{4\mu} (p_4 \cdot (\epsilon(p_3))) - p_{3\nu} p_{4\mu} (p_4 \cdot (\epsilon(p_3)))) - g_1 (g_{\mu\nu} (p_4 \cdot (\epsilon(p_3))) + 2p_{3\mu} \epsilon_\nu (p_3)))
\]

The amplitudes for the $D^* D^*$ final state are:

\[
A_{3a} = i \frac{g_\pi DD \cdot g_\Gamma \phi DD^*}{t - M_B^2} (\epsilon(p_1), (p_4), p_3 - p_2, p_4) (p_2 \cdot (\epsilon(p_3)))
\]

\[
A_{3b} = i \frac{g_\pi D^* D^* \cdot g_\Gamma \phi DD^*}{u - M_B^2} (\epsilon(p_1), p_4 - p_2, (p_3) (p_2 \cdot (\epsilon(p_3)))
\]

\[
A_{3c} = \frac{g_\pi D^* D^* \cdot g_\Gamma \phi DD^*}{t - M_B^2} (\epsilon(p_3 - p_2, \tau, p_3, \epsilon(p_3)) \times \Pi^{\lambda \tau} (p_3 - p_2, M_B^2) (p_{3\lambda} (\epsilon(p_1) \cdot (\epsilon(p_4))) - \epsilon_\lambda (p_1) (p_3 - p_2) \cdot (\epsilon(p_4))) - \epsilon_\lambda (p_4) (p_3 - p_2) \cdot (\epsilon(p_1)))
\]

\[
A_{3d} = \frac{g_\pi D^* D^* \cdot g_\Gamma \phi DD^*}{u - M_B^2} (\epsilon(p_4, (p_4), p_3 - p_1, \beta) \times \Pi^{\beta \lambda} (p_3 - p_1, M_B^2) (-\epsilon(p_3) \cdot (\epsilon(p_1)) + p_{3\lambda} (\epsilon(p_1))) \epsilon_\lambda (p_3) - (p_3 - p_1) \cdot (\epsilon(p_3)) \epsilon_\lambda (p_1))
\]

\[
A_{3e} = i M_D p_{2\nu} \epsilon_\mu (p_1) \left( \frac{g_8 M_D}{g_6} \epsilon(\mu, (p_3), (p_4), p_3) p_5^\nu - \frac{g_6 M_D}{g_8} \epsilon(\mu, (p_3), (p_4), p_3) p_5^\nu + \frac{g_1 + 2g_7}{M_D^2} \epsilon(\mu, (p_4), (p_4), p_3) p_5^\nu + \frac{g_9}{M_D^2} \epsilon(\mu, (p_3), (p_4), p_3) p_5^\nu + \frac{g_1}{M_D^2} \epsilon(\nu, (p_4), (p_3), (p_4)) p_4^\nu + \frac{g_5}{M_D^2} \epsilon(\mu, (p_3), (p_4), (p_4)) p_4^\nu + \frac{g_1 + g_7}{M_D^2} \epsilon(\nu, (p_4), (p_3), (p_4)) p_4^\nu - g_4 \epsilon(\mu, (\nu, (p_3), (p_4)) - \frac{g_5 + g_7}{M_D^2} \epsilon(\mu, (p_4), (p_3), (p_4)) p_4^\nu - \frac{g_7}{M_D^2} (2\epsilon(\nu, (p_3), (p_4), (p_4)) - \epsilon(\nu, (p_4), (p_3), (p_4) p_4^\nu + \frac{g_9}{M_D^2} \epsilon(\mu, (p_3), (p_4), (p_4)) p_4^\nu + \frac{g_7}{M_D^2} (2\epsilon(\nu, (p_3), (p_4), (p_4)) - \epsilon(\nu, (p_4), (p_3), (p_4)) p_4^\nu)
\]

\[
2\epsilon(\mu, (p_4), (p_3), (p_3)) - 3g_7 \epsilon(\mu, (p_3), (p_4), (p_4)) + \epsilon(\mu, (p_4), (p_3), (p_4)) p_4^\nu)
\]

\[
2\epsilon(\mu, (p_4), (p_3), (p_3)) - 3g_7 \epsilon(\mu, (p_3), (p_4), (p_4)) + \epsilon(\mu, (p_4), (p_3), (p_4)) p_4^\nu)
\]

\[
2\epsilon(\mu, (p_4), (p_3), (p_3)) - 3g_7 \epsilon(\mu, (p_3), (p_4), (p_4)) + \epsilon(\mu, (p_4), (p_3), (p_4)) p_4^\nu)
\]
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