Abstract

Values for the cross sections $\sigma_{\Psi N}^{tot}$ and $\sigma_{\Psi' N}^{tot}$ are of crucial importance, once one wants to understand nuclear suppression of $J/\Psi(\Psi')N$ in a search for the quark gluon plasma. Quoted values are usually extracted from $\gamma N \rightarrow J/\Psi(\Psi')N$ data via the vector dominance model (VDM) and are very small compared with expectations from QCD. The validity of QCD is questioned. When the VDM is extended to a multi-channel case, the $\gamma N$ data lead to a value $2.8 \pm 0.12 \text{ mb} < \sigma_{\Psi N}^{tot}(\sqrt{s} = 10 \text{ GeV}) < 4.1 \pm 0.15 \text{ mb}$, which is a factor 3 larger than the VDM prediction. Errors are due to experimental input, and the corridor shows the theoretical uncertainty. We also derive from the data $\sigma_{\Psi N}^{tot} / \sigma_{J/\Psi N}^{tot} \approx 4$, where we could not estimate the theoretical uncertainty.
1. Introduction

The success of the vector dominance model (VDM) for the description of high-energy
$\gamma p$ interactions, particularly for the production of light vector mesons \[1\] has led to a wide
spread confidence that the VDM also accurately describes the production of heavy flavoured
vector mesons. Using data on elastic photoproduction of charmonia, $J/\Psi$ and $\Psi'$, the use
of the VDM leads to a $J/\Psi$-nucleon total cross section $\sigma_{\text{tot}}^{\Psi N} \approx 1.3$ mb for $\sqrt{s} = 10$ GeV and
$\sigma_{\text{tot}}^{\Psi' N}/\sigma_{\text{tot}}^{\Psi N} \approx 0.8$ (see below), which are factors 2 to 4 smaller than predicted by QCD \[2, 3\].

Reliable values for the charmonium interaction cross sections are of crucial importance,
since $J/\Psi$ and $\Psi'$ suppressions in heavy ion collisions are possible signatures for the quark-
gluon plasma (QGP) formation. One needs to know the cross section $\sigma_{\text{tot}}^{\Psi N}$ in the absence
of a QGP-environment in order to predict a nuclear suppression of $J/\Psi$, which can be
used as a base line in search for new physics. An analysis of available data for proton-
nucleus collisions \[4, 5\] in a simple model assuming instantaneous production of the $J/\Psi$ \[4\]
(which has a reasonable accuracy at medium energies $\sqrt{s} \sim 10$ GeV \[6\]), leads to a value of
$\sigma_{\text{tot}}^{\Psi N} \approx 6$ mb \[7\], which is substantially higher than the VDM prediction (see also \[8\]).

In this paper we argue that the discrepancy between the values for $\sigma_{\text{tot}}^{\Psi N}$ from photopro-
duction and hadroproduction data is to be sought in a too naive analysis of photoproduction
data.

Also a disagreement of the VDM prediction for the $\phi N$ cross section with the prediction
of the additive quark model is claimed in \[9, 10\]. In this case more information about $\sigma_{\text{tot}}^{\phi N}$
is available than for charmonium. Facing such a problem with VDM in the strange sector,
one reasonably concludes \[10\] that predictions of the VDM for heavier flavours are even less
trustable.

The purpose of this paper is to clarify why the VDM fails for heavy flavours, and how
it can be modified so that one can extract reliable values for heavy quarkonium interaction
cross sections from photoproduction data.

2. Predictions of the VDM

The VDM establishes a relation between the elastic photoproduction $\gamma N \rightarrow VN$ cross
section of a vector meson $V$ and the total cross section $\sigma^{VN}_{tot}$ of this meson on a nucleon $||$, 

$$
(\sigma^{VN}_{tot})_{VDM} = \left[ \frac{16\pi\alpha_{em}M_V}{3} \frac{d\sigma(\gamma N \rightarrow VN)}{dt} \bigg|_{t=0} \right]^{1/2}.
$$

(1)

Here $\rho_V$ is the ratio of real to imaginary parts of the forward $VN$ elastic scattering amplitude. In most cases, data are available only for the angle integrated cross section, which is related to the forward elastic photoproduction cross section by $d\sigma(\gamma N \rightarrow VN)/dt|_{t=0} = B_{el}^{VN} \sigma(\gamma N \rightarrow VN)$. The slope $B_{el}^{VN}$ of the elastic $J/\Psi N$ differential cross section is determined mainly by the nucleon formfactor, since the radius of heavy quarkonium is small, $\langle r^2 \rangle_\Psi \ll \langle r^2 \rangle_N$. Therefore, $B_{el}^{VN}$ is only slightly larger than the half of the slope of elastic $pp$ scattering. Our estimate $B_{el}^{\Psi N} \approx 6 \text{ GeV}^{-2}$ [11] agrees well with experimental data [12].

Refs. [13, 14, 15] give a collection of experimental data, which we fit with the expression,

$$
\sigma(\gamma N \rightarrow \Psi N) = \sigma_0 \left( \frac{\sqrt{s}}{10 \text{ GeV}} \right)^{2\lambda},
$$

(2)

with parameters $\sigma_0 = (10.3\pm0.7) \text{ nb}$, $\lambda = 0.40\pm0.02$. We can estimate $\rho_\Psi = (\pi/2)d[\ln(\sigma^{\Psi N}_{tot})]/d[\ln(s)] = \pi\lambda/4 \approx 0.3$.

Applying the VDM relation [1] to the parameterization [2] we extract a value

$$
(\sigma^{\Psi N}_{tot})_{VDM} = (1.24 \pm 0.13) \text{ mb} \times (\sqrt{s}/10 \text{ GeV})^\lambda.
$$

(3)

Using the VDM relation [1] one can also extract the total cross section for the radial excitation $\Psi'(2S)$. One finds for the ratio of the $\Psi'$ to $\Psi$ cross sections

$$
\frac{\sigma^{\Psi'N}_{tot}}{\sigma^{\Psi N}_{tot}} = \left( \frac{M_{\Psi'} \Gamma_{\Psi'}}{M_{\Psi'} \Gamma_{\Psi'}} \right)^{1/2}
$$

(4)

with the ratio of the photoproduction cross sections $R = \sigma(\gamma N \rightarrow \Psi'N)/\sigma(\gamma N \rightarrow \Psi N)$. For the experimental value $R = 0.195 \pm 0.03$ [16-18], eq. (4) gives

$$
\left( \frac{\sigma^{\Psi'N}_{tot}}{\sigma^{\Psi N}_{tot}} \right)_{VDM} = 0.8 \pm 0.2.
$$

(5)

This result contradicts dramatically the expectation based on QCD, *Actually, the exponent can be about 20% larger than $\lambda$ due to energy dependence of the slope $B_{el}^{\Psi N}$
\[
\frac{\sigma_{\Psi'N}}{\sigma_{\Psi N}} = \frac{< r^2 >_{\Psi'}}{< r^2 >_{\Psi}}.
\]

This ratio is 7/3 in the harmonic oscillator model for \( J/\Psi \) and \( \Psi' \), and equals 4 in a more realistic potential model [19].

3. Why does VDM fail for heavy flavours?

The basic assumption of the VDM is that via quantum mechanical fluctuations the photon converts into a vector meson \( V \), which interacts with the proton elastically and thereby comes on the energy shell. Is the hadronic fluctuation of the \( \gamma \) always a vector meson? In ref. [20, 21] it is proposed that the photoproduction of \( J/\Psi \) should be considered in perturbative QCD as an interaction of a \( c\bar{c} \) fluctuation (and not a \( J/\Psi \) fluctuation) of the photon with the target followed by projection of the produced \( c\bar{c} \) wave packet on the \( J/\Psi \) wave function. Only if the wave function of the \( c\bar{c} \) component of the photon is similar to that for the \( J/\Psi \), this model is close to the conventional VDM.

In order to see when and to what degree the wave functions of \( Q\bar{Q} \) fluctuations coincide with those of the vector mesons, considered as \( Q\bar{Q} \) bound states, we study the mean squared transverse radii. The wave function of the \( Q\bar{Q} \) component of the photon has the form [2],

\[
\Psi_{Q\bar{Q}}(r_T) \propto K_0(m_Qr_T)
\]

where \( K_0(x) \) is the modified Bessel function and \( r_T \) is the transverse \( Q\bar{Q} \) separation. Its mean transverse size squared is calculated from (7)

\[
\langle r_{T}^2 \rangle_\gamma = \frac{2}{3m_Q^2}.
\]

This value has to be compared with the mean transverse distance between the \( Q \) and \( \bar{Q} \) in the vector meson, considered as the lowest state in an harmonic oscillator potential with frequency \( \omega \):

\[
\langle r_{T}^2 \rangle_V = \frac{2}{m_Q \omega}.
\]
Empirically one finds that \( \omega = (M_{V(2S)} - M_{V(1S)})/2 \) has a value of roughly 300 MeV independent of the flavour.

The two sizes eqs. (8) and (9) depend differently on \( m_Q \), they cannot coincide for all flavours, in particular, the discrepancy increases the heavier the quark mass is.

We introduce the factor \( N_\Psi \) and \( N_{\Psi'} \) for the transition from the VDM prediction to the real value

\[
\sigma_{tot}^{\Psi N} = \frac{1}{N_\Psi} \sigma_{tot}^{\Psi N}_{VDM} \tag{10}
\]

(and similarly for \( \Psi' \)) and attempt to calculate these factors by two different methods in the next sections.

4. Hadronic representation: a multi-channel approach

The above observation that the \( c\bar{c} \) fluctuation is quite different from the \( J/\Psi \) can be phrased formally in writing an expansion of the \( c\bar{c} \) fluctuation in terms of a complete set of charmonium states

\[
|\gamma\rangle_{c\bar{c}} = \alpha_1 |J/\Psi\rangle + \alpha_2 |\Psi'\rangle + \ldots . \tag{11}
\]

With this expansion of the \( c\bar{c} \) fluctuation, photoproduction of \( \Psi' \) and \( J/\Psi \) gains contributions from different mechanisms as is shown schematically in Fig. 2 for a two channel example: i) direct production as in the VDM \( \gamma \rightarrow \Psi, \Psi N \rightarrow \Psi N \) and ii) indirect production, \( \gamma \rightarrow \Psi', \Psi'N \rightarrow \Psi N \), which contains the off-diagonal diffractive interaction \( \Psi N \leftrightarrow \Psi'N \). The two amplitudes which lead to the final state must be added before squaring. The non-diagonal process \( \Psi \rightarrow \Psi' \) is the correction to the VDM expression which we propose.

The virtual photoproduction amplitude in the naive VDM reads,

\[
f_{VDM}(\gamma^*N \rightarrow \Psi N) = \frac{C_{\gamma\Psi}^2}{M_{\Psi}^2 + Q^2} f(\Psi N \rightarrow \Psi N) , \tag{12}
\]

where the photon of virtuality \( Q^2 \) converts into \( \Psi \) with a probability amplitude \( C_{\gamma\Psi}^2 \), the \( J/\Psi \) propagates with a usual propagator and \( f(\Psi N \rightarrow \Psi N) \) denotes the elastic scattering amplitude. The extension of eq. (12) to the multi-channel case is straightforward:
Figure 1: Diagrams for the elastic photoproduction of the vector mesons, an example of two channels. Photoproduction of the ground state $V$ (a), and of the radial excitation $V'$ (b).

\[ f(\gamma^*N \rightarrow \Psi N) = N_\Psi(Q^2) f_{VDM}(\gamma^*N \rightarrow \Psi N), \]  

(13)

where we have defined the correction factor $N_\Psi(Q^2)$

\[ N_\Psi(Q^2) = 1 + \sum_{i \geq 1} \frac{C_{\gamma\Psi_i} f(\Psi_i N \rightarrow \Psi N)}{C_{\gamma\Psi} f(\Psi N \rightarrow \Psi N)} \frac{M_\Psi^2 + Q^2}{M_{\Psi_i}^2 + Q^2}. \]  

(14)

It follows from perturbative QCD (and is shown in the next section) that $N_\Psi \propto 1/Q^2$ for $Q^2 \rightarrow \infty$. This property of color transparency implies for the expression (14) a sum rule,

\[ \sum_{i \geq 1} \frac{C_{\gamma\Psi_i} f(\Psi_i N \rightarrow \Psi N)}{C_{\gamma\Psi} f(\Psi N \rightarrow \Psi N)} = -1. \]  

(15)

For heavy quarkonia the mass splitting is much smaller than the mass, therefore, we expect a strong cancellation in (14) even at low values of $Q^2$. Indeed, the zero order term in an expansion of $N_\Psi(0)$ in the small parameter $\Delta M_\Psi/M_\Psi$ vanishes according to (15). The first nonzero term in the expansion is of the order of

\[ N_\Psi(0) \sim \frac{\Delta M_\Psi}{M_\Psi} \approx 0.2. \]  

(16)

The observed strong cancellation between different channels in (13) is a general property of heavy quarkonia, and we do not expect the factor $N_\Psi(Q^2)$ to depend much on the set of orthogonal states chosen for the hadronic basis. In one case the function $N_\Psi(Q^2)$ can
be calculated exactly. If the operator for the elastic scattering amplitude of a $c\bar{c}$ pair on a nucleon is proportional to the transverse quark separation squared $r^2_T$, and harmonic oscillator wave functions are chosen for the hadronic basis, one has $f(\Psi_i N \rightarrow \Psi N) \propto \langle \Psi_i | r^2_T | \Psi_0 \rangle = 0$ for $i \geq 2$. Then, according to the sum rule (13), one finds the surprisingly simple result,

$$N_{\Psi}(Q^2) = \frac{M^2_{\Psi} - M^2_{\Psi'}}{M^2_{\Psi'} + Q^2}, \quad (17)$$

which is independent of any matrix element. With $N_{\Psi}(0) = 0.30$, one arrives at a factor of 3.3 correction to the naive VDM,

$$\sigma_{tot}^{\Psi N} = (4.3 \pm 0.5) \, mb \times \left( \frac{\sqrt{s}}{10 \, GeV} \right)^{0.4}. \quad (18)$$

How good in the correction factor (17)? We will give an independent calculation of $N_{\Psi}(Q^2)$ in the next sections. In fig. 2 we compare our predictions with the factor $N_{\Psi}(Q^2)$ Eq. (17) and with $N_{\Psi}(Q^2) = 1$, which is the conventional VDM, and available data from the EMC experiment [17]. The comparison obviously supports the form (17), although the error bars are rather large.

For the $\Psi'$ production we have to introduce at least three intermediate states $\Psi$, $\Psi'$ and $\Psi''$, which exhaust the sum in the hadronic oscillator basis and have

$$f(\gamma^* p \rightarrow \Psi' p) = f_{VDM}(\gamma^* p \rightarrow \Psi' p) \times \left[ 1 + \sum_{i=0,2} C^\gamma \Psi_i \frac{f(\Psi_i N \rightarrow \Psi' N)}{f(\Psi' N \rightarrow \Psi' N)} \frac{M^2_{\Psi'} + Q^2}{M^2_{\Psi_i} + Q^2} \right], \quad (19)$$

where the squared bracket defines $N_{\Psi'}(Q^2)$. In an oscillatory basis $f(\Psi' N \rightarrow \Psi' N)/f(\Psi N \rightarrow \Psi N) = 7/3$. Together with the color transparency sum rule (13) all the parameters are determined

$$N_{\Psi'}(Q^2) = 1 - \frac{2}{7} \frac{M^2_{\Psi'} + Q^2}{M^2_{\Psi'} + Q^2} - \frac{5}{7} \frac{M^2_{\Psi''} + Q^2}{M^2_{\Psi''} + Q^2}, \quad (20)$$

which expression leads to $N_{\Psi'}(0) = 0.064$. Then the ratio of the corrected values for the
Figure 2: Data on the $Q^2$-dependence of the exclusive muoproduction of $J/\Psi$ \cite{17}. The dashed curve is the expectation of the VDM, the solid curve includes the correction factor \cite{17}.

The ratio of $\Psi'$ to $\Psi$ total cross sections is

$$\frac{\sigma_{\Psi'}_{\text{tot}}}{\sigma_{\Psi}_{\text{tot}}} = \frac{N_{\Psi'}(0)}{N_{\Psi}(0)} \left( \frac{\sigma_{\Psi'}_{\text{tot}}}{\sigma_{\Psi}_{\text{tot}}} \right)_{VDM} = 3.75.$$ \hspace{1cm} (21)

This ratio corresponds nearly exactly to the QCD prediction, namely, the ratio of radii squared, which calculated for a realistic potential \cite{19} gives a value 4.

5. Quark representation

In order to estimate the theoretical uncertainty of the correction factor $N_{\Psi}(0)$ and $N_{\Psi'}(0)$ evaluated using the harmonic oscillator basis, we calculate these factors by yet another approach.
The multi-channel equations (13) - (14) can be represented in the quark basis as,

\[ f(\gamma^* N \rightarrow \Psi N) = \langle \Psi | \sigma(r_T) | \gamma^{*}_{c\bar{c}} \rangle \]  

(22)

with the wave function for the photon \( |\gamma^{*}_{c\bar{c}}\rangle \) in the quark basis. In the same representation, the expression for the VDM reads

\[ f(\gamma^* N \rightarrow \Psi N) = \langle \Psi | \sigma(r_T) | \Psi \rangle \langle \Psi | \gamma^{*}_{c\bar{c}} \rangle . \]  

(23)

Then the correction factor takes the form

\[ N_\Psi(Q^2) = \frac{\langle \Psi | \sigma(r_T) | \gamma^{*}_{c\bar{c}} \rangle}{\langle \Psi | \sigma(r_T) | \Psi \rangle \langle \Psi | \gamma^{*}_{c\bar{c}} \rangle} , \]  

(24)

where \( \Psi \) may stand for the wave functions of the \( J/\Psi \) or the \( \Psi' \). The perturbative wave function of the \( c\bar{c} \) fluctuation of the virtual photon reads in non-relativistic approximation,

\[ |\gamma^{*}_{c\bar{c}}\rangle \propto K_0(\epsilon r_T) \]  

\[ \epsilon^2 = m_c^2 + Q^2/4. \]  

This perturbative QCD expression for \( |\gamma^{*}_{c\bar{c}}\rangle \) neglects the interaction between \( c \) and \( \bar{c} \) and therefore is questionable for low \( Q^2 \). In this respect the hadronic representation is preferable, since effectively it takes into account all the non-perturbative effects.

The mean squared \( c\bar{c} \) separation in the photon is calculated to

\[ \langle r^2_T \rangle_{\gamma^*} = 2/3(m_c^2 + Q^2/4) \]  

and goes to zero for \( Q^2 \rightarrow \infty \). In this limit and using \( \sigma(r_T) = ar_T^2 \), the expression for \( N_\Psi(Q^2) \) becomes very simple, because \( \Psi(r_T) \) can be replaced by \( \Psi(0) \) in the integrals involving \( K_0(r_T) \). Then,

\[ N_\Psi \rightarrow \frac{\langle r^2_T \rangle_{\gamma^*}}{\langle r^2_T \rangle_{\Psi}} \quad (Q^2 \rightarrow \infty) . \]  

(25)

This result is exact. We note, that \( \langle r^2_T \rangle_{\gamma^*} \) goes to zero with the \( Q^2 \rightarrow \infty \) and this leads to the sum rule (15). If one uses this expression also at \( Q^2 = 0 \) together with \( \langle r^2_T \rangle_\Psi = 2/\omega m_c \) eq. (9) one finds \( N_\Psi(0) = 2\omega/m_c \), which agrees with (17) to lowest order in \( \omega/2m_c \). In the same limit

\[ \frac{N_{\Psi'}(Q^2)}{N_{\Psi}(Q^2)} = \frac{\langle r^2_T \rangle_{\Psi'}}{\langle r^2_T \rangle_{\Psi}} . \]  

(26)

The expression of \( N_\Psi(Q^2) \) in (25) is based on the approximation \( \langle r^2_T \rangle_{\gamma^*} \ll \langle r^2 \rangle_{\Psi} \), which is not well satisfied for charmonium. In order to test the accuracy, we give up this approximation, and use another trial wave functions \( |\Psi\rangle \), among them the harmonic oscillator ones.
Be specific, we use a mixture of the harmonic oscillator (HO) and Coulomb (C) wave functions,

$$\langle \vec{r} | \Psi \rangle = \alpha \langle \vec{r} | \Psi_{HO} \rangle + \beta \langle \vec{r} | \Psi_C \rangle$$ (27)

for 1S and 2S states. The Coulomb wave function is calculated with a running QCD coupling as is explained in [23]. We vary $\alpha$ ($\beta = \sqrt{1 - \alpha^2}$) and adjust the parameters of $|\Psi_{HO}\rangle$ and $|\Psi_C\rangle$ so that one has always the same $\langle r^2_T \rangle_\Psi$ as taken from [13].

Then we obtain the following results:

i) for the $J/\Psi$

$$2.3 \leq N_{\Psi}^{-1} \leq 3.3,$$ (28)

where the upper limit corresponds to (14) and the lower one is calculated from eq. (19bb) with $\alpha = 0$ in (27). The case of pure Coulomb wave function ($\alpha = 0$) lies in the interval. Thus,

$$2.8 \pm 0.12 \leq \sigma_{\Psi N}^{\Psi N}(\sqrt{s} = 10 \, GeV) \leq (4.1 \pm 0.15) \, mb$$ (29)

with the energy dependence

$$\sigma_{\Psi N}^{\Psi N}(\sqrt{s}) = \sigma_{\Psi N}^{\Psi N}(10 \, GeV) \left( \frac{\sqrt{s}}{10 \, GeV} \right)^{0.4}$$ (30)

ii) The results for $N_{\Psi'}(0)$ are very sensitive to the wave function (27) (because of the node), so that we prefer to give the result in the form (21) without a corridor of theoretical uncertainty.

Furthermore, we test the sensitivity of the results to the form of the transition operator $\sigma(r_T)$. We choose the form $\sigma(r_T) = a[1 - \exp(-br_T^2)]/b$ [23], which gives the linear $r_T^2$ dependence at small $r_T$. We found a very small deviation of the order of few percent.

6. Conclusion and discussions

The conventional expressions of the VDM applied to the photoproduction data of charmonium production lead to inconsistencies both with experimental data ($\sigma_{\Psi N}^{\Psi N}$ values extracted from $hA$ data) as well as predictions from QCD ($\sigma_{\Psi N}^{hN} \propto \langle r^2 \rangle_h$). A generalization of
the VDM to a multi-channel problem leads to a value $\sigma_{tot}^\Psi N$ about three times as large as the one predicted by VDM, and the discrepancies relating to $\sigma_{tot}^\Psi N$ are largely removed.

Few more comments are in order.

- Theoretical accuracy of the correction factor $N_\Psi(Q^2)$ is estimated in the quark representation by varying the shape of the $J/\Psi$ wave function and the form of the dipole cross section. We conclude that the uncertainty does not exceed 20%.

- Another source of information about $\sigma_{tot}^\Psi N$ is production of $J/\Psi$ off nuclei, which give a slightly higher value than we predict. This is because the simplified analysis does not take into account extension of the charmonium length of path in nuclear matter, due to the coherence length, which is about 2 fm for the kinematics of $J/\Psi$. Therefore a reduction of about 20% should be applied to $\sigma_{tot}^\Psi N$. Besides, dynamics of charmonium hadroproduction is far more complicate than for photoproduction.

In order to be free of the effects of the coherence and formation lengths one should produce a low energy charmonium off nuclei. A unique opportunity is annihilation $\bar{p}p \rightarrow \Psi$ on a bound protons.

- Using the same formulas for photoproduction of bottomia we arrive at the result, $\sigma_{tot}^\Upsilon N = 8 (\sigma_{tot}^\Upsilon N)_VDM$. The cross section $(\sigma_{tot}^\Upsilon N)_VDM$ as well as the ratio $r_\Upsilon$ cannot be predicted, since no data on bottomium photoproduction is available yet.

- We expect a weak rising energy-dependence of the correction factor $N_\Psi(Q^2)$. Indeed, it is known from HERA data on the proton structure function $F_2(x, Q^2)$ that the smaller is the size of a photon fluctuation, the steeper is the energy dependence. This is why the $J/\Psi$ photoproduction cross section growth so steeply compared to what is known from soft hadronic interactions. According to the QCD evolution equations this effect is due to fast grows of the gluon cloud of the $c\bar{c}$ fluctuation of the photon, which steeper than that of $J/\Psi$.

- One may wonder why VDM works so well for $\rho$ mesons. Indeed, the same expression
(17) leads to deviation of 40% from the VDM prediction. This estimate contains an additional uncertainty since the dipole approximation $\sigma(r_T) \propto r_T^2$ is justified only at small $r_T$, while large distances are important for $\rho$ mesons. A harmonic oscillator is also too rough an approximation in this case.

- One should be cautious applying VDM to the photoproduction of $J/\Psi$ on nucleons at low energy. If the lifetime of the $J/\Psi$ fluctuation in the photon is much shorter than the nucleon radius, $t_c \approx M_{\Psi}^2/2E_\gamma \ll R_N$, the photoproduction cross section should be much smaller than VDM predicts, since the fluctuation can interact only during its short lifetime. This effect is taken into account by the nucleon form factor (compare with [25]).

Acknowledgement: We are thankful to M.G. Ryskin and B.G. Zakharov for useful discussions and suggestions. This work has been supported by a grant from the Gesellschaft für Schwerionenforschung Darmstadt (GSI), grant no. HD HÜ T, and the Bundesministerium für Bildung, Wissenschaft, Forschung und Technologie (BMBF), grant no. 06 HD 856.

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