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HEAVY-QUARK PRODUCTION IN PHOTON–NUCLEON AND PHOTON–PHOTON COLLISIONS*

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I discuss mechanisms of heavy quark production in (real) photon–nucleon and (real) photon–(real) photon collisions. In particular, I focus on application of the Saturation Model. In addition to the main dipole–nucleon or dipole–dipole contribution included in recent analyses, I propose how to calculate within the same formalism the hadronic single-resolved contribution to heavy quark production. At high photon–photon energies this yields a sizeable correction of about 30–40 % for inclusive charm production and 15–20 % for bottom production.

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1. Introduction

The total cross section for virtual photon–proton scattering in the region of small x and intermediate Q^2 can be well described by the Saturation Model (SM) [1]. The very good agreement with experimental data can be extended even to the region of rather small Q^2 by adjusting an effective quark mass. At present there is no deep understanding of the fit value of the parameter as we do not understand in detail the confinement and the underlying nonperturbative effects related to large size QCD contributions.

In this presentation I shall limit to the production of heavy quarks which is simpler and more transparent for real photons. Here one can partially avoid the problem of the poor understanding of the effective light quark mass, i.e. the domain of the large (transverse) size of the hadronic system emerging from the photon.

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It was shown recently that the simple SM description can be successfully extended also to the photon–photon scattering [2]. The heavy quark production in photon–photon collisions is interesting in the context of a deficit of standard QCD predictions relative to the experimental data as observed recently for b quark production.

2. Heavy quark production in photon–nucleon scattering

In the picture of dipole scattering the cross section for heavy quark–antiquark ($Q\bar{Q}$) photoproduction on the nucleon can be written as

$$\sigma_{\gamma N \to Q\bar{Q}}(W) = \int d^2\rho dz \, |\Phi_{\mathrm{T}}^{Q\bar{Q}}(\vec{\rho}, z)|^2 \sigma_{dN}(\rho, z, W) \,, \tag{1}$$

where Φ_{T} is (transverse) quark–antiquark photon wave function (see for instance [4]) and σ_{dN} is the dipole–nucleon total cross section. Inspired by its phenomenological success [1] we shall use the SM parametrisation for σ_{dN}. Because for real photoproduction the Bjorken-x is not defined we are forced to replace x by x_g [3].

In Fig. 1(a) we show predictions of SM for charm photoproduction. The dotted line represents calculations based on Eq.(1). The result of this calculation exceeds considerably the fixed target experimental data. One should remember, however, that the simple formula (1) applies at high energies only. At lower energies one should include effects due to kinematical threshold. In the momentum representation this can be done by requiring: $M_{Q\bar{Q}} < W$, where $M_{Q\bar{Q}}$ is the invariant mass of the final $Q\bar{Q}$ system. This upper limit still exceeds the low energy experimental data. There are phase space limitations in the region $x_g \to 1$ which have been neglected so far. Those can be estimated using naive counting rules. Such a procedure leads to a reasonable agreement with the fixed target experimental data.

The deviation of the solid line from the dotted line gives an idea of the range of the safe applicability of SM for the production of the charm quarks/antiquarks. The cross section for $W > 20$ GeV is practically independent of the approximate treatment of the threshold effects. SM seems to slightly underestimate the H1 collaboration data [5]. For comparison in Fig. 1 we show the result of similar calculations in the collinear approach (thick dash-dotted line) with details described in [3].

The calculation above is not complete. For real photons a vector dominance contribution due to photon fluctuation into vector mesons should be included on the top of the dipole contribution. In the present calculation we include only the dominant gluon–gluon fusion component. Then

$$\sigma_{\gamma N \to Q\bar{Q}}^{\mathrm{VDM}}(W) = \sum_V \frac{4\pi}{f_V^2} \int dx_V dx_N g_V(x_V, \mu_F^2) g_N(x_N, \mu_F^2) \sigma_{gg \to Q\bar{Q}}(\hat{W}) \,. \tag{2}$$

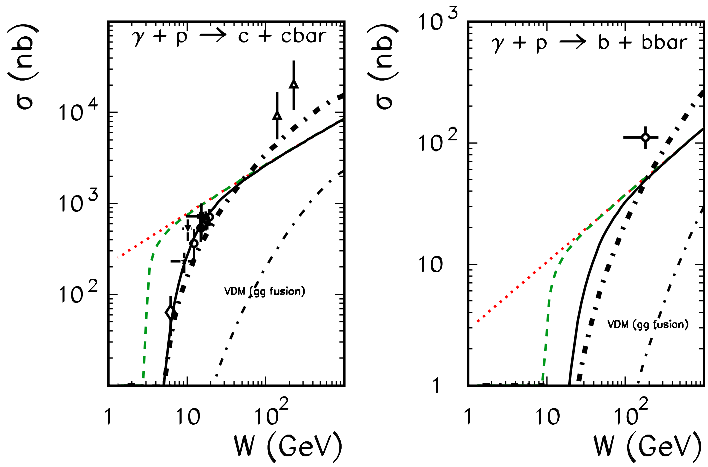

Fig. 1. The cross section for $\gamma + p \to Q\bar{Q}X$. The dotted line: standard SM, the dashed line: includes kinematical threshold, the solid line includes in addition a suppression by $(1 - x_c)^7$, the thick dash-dotted line: collinear approximation and the thin dash-dotted line: the LO VDM contribution.

Here the f_V constants describe the transition of the photon into vector mesons (ρ, ω, ϕ). The gluon distributions in vector mesons are taken as that for the pion [7].

The dash-dotted line in Fig. 1(a) shows the VDM contribution calculated in the Leading Order (LO) approximation for $\sigma_{gg \to Q\bar{Q}}$. The so-calculated VDM contribution cannot be neglected at high energies.

The situation for bottom photoproduction seems similar. In Fig. 1(b) we compare the SM predictions with the data from the H1 collaboration [6]. Here the threshold effects may survive up to very high energy $W \sim 50$ GeV. Again the predictions of SM are slightly below the H1 experimental data point. The relative magnitude of the VDM component is similar as for the charm production.

3. Heavy quark production in photon–photon scattering

In the dipole–dipole approach

$$\begin{aligned}
&\sigma^{\mathrm{dd}}_{\gamma\gamma \to Q\bar{Q}}(W) \\
&= \sum_{f_2 \neq Q} \int |\Phi^{Q\bar{Q}}(\rho_1, z_1)|^2 |\Phi^{f_2 \bar{f}_2}(\rho_2, z_2)|^2 \sigma_{\mathrm{dd}}(\rho_1, \rho_2, x_{Qf}) d^2\rho_1 dz_1 d^2\rho_2 dz_2 \\
&+ \sum_{f_1 \neq Q} \int |\Phi^{f_1 \bar{f}_1}(\rho_1, z_1)|^2 |\Phi^{Q\bar{Q}}(\rho_2, z_2)|^2 \sigma_{\mathrm{dd}}(\rho_1, \rho_2, x_{fQ}) d^2\rho_1 dz_1 d^2\rho_2 dz_2, \quad (3)
\end{aligned}$$

where σ_{dd} is the dipole–dipole cross section.

There are two problems associated with direct use of (3). First of all, it is not completely clear how to generalise $\sigma_{\rm dd}$ from σ_{dN} parametrised in [1]. Secondly, formula (3) is correct only at $W \gg 2m_Q$. At lower energies one should worry about proximity of the kinematical threshold.

In a very recent paper [2] a new phenomenological parametrisation for $\sigma_{\rm dd}$ has been proposed. The phenomenological threshold factor in [2] does not guarantee automatic vanishing of the cross section exactly below the true kinematical threshold $W = 2m_a + 2m_b$. Therefore, instead of the phenomenological factor we rather impose an extra kinematical constraint: $M_{f\bar{f}} + M_{Q\bar{Q}} < W$ on the integration in (3).

It is not completely clear how to generalise the energy dependence of σ_{dN} in photon–nucleon scattering to the energy dependence of σ_{dd} in photon–photon scattering. In [3] I have defined the parameter which controls the SM energy dependence of σ_{dd} in a symmetric way with respect to both photons. In comparison to the prescription in [2], our prescription leads to a small reduction of the cross section far from the threshold [3].

Up to now we have calculated the contribution when photons fluctuate into quark-antiquark pairs. The dipole approach must be supplemented to include the contribution when either of the photons fluctuates into vector mesons. If the first photon fluctuates into the vector mesons, the so-defined single-resolved contribution to the heavy quark-antiquark production can be calculated analogously to the photon–nucleon case as

$$\sigma^{\rm SR,1}_{\gamma\gamma\to Q\bar{Q}}(W) = \sum_{V_1} \frac{4\pi}{f^2_{V_1}} \int |\Phi_2^{Q\bar{Q}}(\rho_2, z_2)|^2 \sigma_{V_1 d}(\rho_2, x_1)\; d^2\rho_2 dz_2\,, \qquad (4)$$

where $\sigma_{V_1 d}$ is vector meson–dipole total cross section. In the spirit of SM, we parametrise the latter exactly as for the photon–nucleon case [1] with a simple rescaling of the normalisation factor $\sigma_0^{dV} = 2/3 \cdot \sigma_0^{dN}$. In the present calculation σ_0^{dN} as well as the other parameters of SM are taken from [1]. Analogously, if the second photon fluctuates into vector mesons we obtain

$$\sigma^{\rm SR,2}_{\gamma\gamma\to Q\bar{Q}}(W) = \sum_{V_2} \frac{4\pi}{f^2_{V_2}} \int |\Phi_1^{Q\bar{Q}}(\rho_1, z_1)|^2 \sigma_{dV_2}(\rho_1, x_2)\; d^2\rho_1 dz_1\,. \qquad (5)$$

This clearly doubles the first contribution (4) to the total cross section.

The integrations in (4) and (5) are not free of kinematical constraints. When calculating both single-resolved contributions, it should be checked additionally if the heavy quark-antiquark invariant mass $M_{Q\bar{Q}}$ is smaller than the total photon–photon energy W (see [3]).

In Fig. 2 we show different contributions to the inclusive $c/\bar{c}$ (left panel) and $b/\bar{b}$ (right panel) production in photon–photon scattering. The thick solid line represents the sum of all contributions.

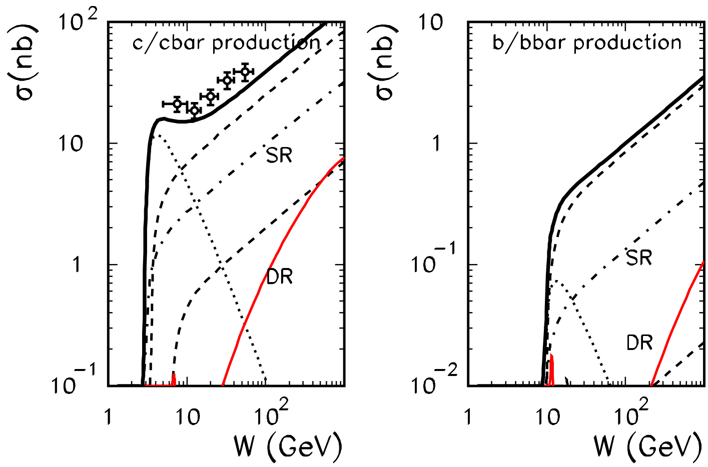

Fig. 2. Different contributions to the inclusive charm (left panel) and bottom (right panel) production. The long-dashed line: the dipole–dipole contribution, the dash-dotted line: the single-resolved contribution, the lower dashed line: the $2Q2\bar{Q}$ contribution, the dotted line: the direct contribution, the solid line: double-resolved contribution. The experimental data for inclusive $c/\bar{c}$ production are from Ref. [9].

Let us start from the discussion of the inclusive charm production. The experimental data of the L3 collaboration [9] are shown for comparison. The modifications discussed above lead to a small damping of the cross section in comparison to [2]. The corresponding result (long-dashed line) stays below the recent experimental data of the L3 collaboration [9]. The hadronic single-resolved contribution constitutes about 30–40 % of the main SM contribution. At high energies the cross section for the $2c2\bar{c}$ component is about 8 % of that for the single $c\bar{c}$ pair component. In the inclusive cross section its contribution should be doubled because each of the heavy quarks/antiquarks can be potentially identified experimentally.

At higher energies the direct contribution is practically negligible. In contrast, the hadronic double-resolved contribution, when each of the two photons fluctuates into a vector meson [3] is shown by the thin solid line in the figure becomes important only at very high energies relevant for TESLA. Here we have consistently taken $g_V(x_V,\mu_F^2) = g_\pi(x_V,\mu_F^2)$.

The situation for bottom production (see right panel) is somewhat different. Here the main SM component is dominant. Due to smaller charge of the bottom quark than that for the charm quark, the direct component is effectively reduced with respect to the dominant SM component by the corresponding ratio of quark/antiquark charges: $(1/9)^2 : (4/9)^2 = 1/16$. The same is true for the $2b2\bar{b}$ component. Here, in addition, there are threshold effects which play a role up to relatively high energy.

4. Conclusions

There is no common consensus in the literature on detailed understanding of the dynamics of photon–nucleon and photon–photon collisions. In this presentation I have limited the discussion to the production of heavy quarks simultaneously in photon–nucleon and photon–photon collisions at high energies. The sizeable mass of charm or bottom quarks sets a natural low energy limit on a naive application of SM. Here a careful treatment of the kinematical threshold is required.

We have started the analysis from (real) photon–nucleon scattering, which is very close to the domain of SM as formulated in [1]. If the kinematical threshold corrections are included, SM gives similar results as the standard collinear approach for both charm and bottom production. We have estimated the VDM contribution to the heavy quark production.

The second part of the present analysis has been devoted to real photon–real photon collisions. For the first time in the literature we have estimated the cross section for the production of $2c2\bar{c}$ final state. We have found that this component constitutes up to 10–15 % of the inclusive charm production at high energies and is negligible for the bottom production. We have shown how to generalise SM to the case when one of the photons fluctuates into light vector mesons. It was found that this component yields a significant correction of about 30–40 % for inclusive charm production and 15–20 % for bottom production. We have shown that the double resolved component, when both photons fluctuate into light vector mesons, is non-negligible only at very high energies, both for the charm and bottom production.

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