

Production of η_b in ultraperipheral Pb-Pb collisions

C. N. Azevedo[✉], F. C. Sobrinho, and F. S. Navarra[✉]

*Instituto de Física, Universidade de São Paulo, Rua do Matão 1371 - CEP 05508-090,
Cidade Universitária, São Paulo, SP, Brazil*



(Received 2 January 2025; revised 17 September 2025; accepted 27 October 2025; published 11 December 2025)

Very recently, the two-photon decay width of the η_b meson was computed with lattice QCD methods. This decay has not yet been measured. The knowledge of this width allows for the calculation of the η_b production cross section through photon-photon interactions in ultraperipheral (UPC) Pb - Pb collisions. In this work we present this calculation, which is the first application of the lattice result. Since UPCs are gaining increasing attention of the heavy-ion community, we take the opportunity to perform a comprehensive study of the different ways of defining ultraperipheral collisions and of the different ways to treat the equivalent photon flux.

DOI: [10.1103/196s-xxkw](https://doi.org/10.1103/196s-xxkw)

I. INTRODUCTION

Ultraperipheral collisions (UPCs) provide an opportunity to improve our understanding of the standard model as well as to search for new physics [1–6]. In these collisions the incident nuclei do not overlap, which implies the suppression of the strong interactions and the dominance of the electromagnetic interactions between them. Over recent years, the study of photon-induced processes in hadronic colliders has become a reality with a large number of experimental results on different final states obtained by the STAR [7], PHENIX [8], ALICE [9,10], CMS [11], and ATLAS [12] collaborations. These results demonstrate that the study of photon-induced processes in hadronic colliders is feasible and can be useful to improve our understanding of the QCD dynamics at high energies. New states are expected to be observed in the future [13–15]. In particular, ultraperipheral collisions have been attracting increasing interest from the heavy ion community, as they provide a unique laboratory to study photon-photon and photon-hadron interactions [16–19]. The essential feature of these processes is that relativistic heavy ions give rise to strong electromagnetic fields, so that in a hadron-hadron collision, photon-hadron and photon-photon interactions can occur and they may lead to the production of particles [20,21]. Moreover, processes involving photons can be exclusive, where the resulting final state is very clean. A typical example of exclusive process is the production of pseudoscalar mesons due to two-photon fusion. The resulting final state is very simple, consisting of a pseudoscalar meson with very small transverse momentum, two intact nuclei, and two rapidity gaps, i.e., empty regions in pseudorapidity that

separate the intact very-forward nuclei from the produced state. Such aspects can, in principle, be used to separate the events and to test predictions.

Very recently, the decay $\eta_b \rightarrow \gamma\gamma$ was studied in lattice QCD for the first time in Ref. [22], and the decay width was calculated, providing an accurate prediction of $\Gamma_{\eta_b \rightarrow \gamma\gamma}$. This prediction may be tested at Belle II. So far, as discussed in Ref. [23], the efforts were dedicated to improve the accuracy of the η_b mass measurement, which is observed in decays such as $h_b \rightarrow \gamma \eta_b$ by measuring the photon energy. This ability to accurately measure the photon energies is crucial and may allow for the determination of $\Gamma_{\eta_b \rightarrow \gamma\gamma}$.

The description of heavy quark bound states is a challenging problem from the point of view of QCD. A proper quantum-field treatment of the heavy quarkonium system, based on the Bethe-Salpeter equation, turned out to be rather difficult [24]. Two promising approaches to the bottomonium bound-state problem are effective field theories (EFTs) and lattice gauge theories. EFTs may be directly derived from QCD, like nonrelativistic QCD (NRQCD) or potential nonrelativistic QCD (pNRQCD) (for reviews see Ref. [25]). However, so far, neither with EFTs nor with lattice QCD was it possible to compute the properties of the bottomonium excited states. This has also been a problem for QCD sum rules (for reviews see Ref. [26]). A discussion on lattice QCD calculations and their complexities in the bottomonium sector can be found in Ref. [27]. Because of these difficulties, most of the estimates for heavy quarkonia still rely on potential models. These models have always been successful in describing the heavy quark-antiquark spectra. Moreover, their predictions of heavy quarkonium properties have been very valuable for experimental searches.

Previous theoretical estimates of $\Gamma_{\eta_b \rightarrow \gamma\gamma}$ are shown in Table I. In quark (potential) models, this quantity is obtained from the solution of the Schrödinger equation with the Cornell potential. Different works use slightly different assumptions. In Refs. [28,29] the authors use a constant potential term in their nonrelativistic framework and spin-dependent correc-

Published by the American Physical Society under the terms of the Creative Commons Attribution 4.0 International license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

TABLE I. Decay width $\Gamma_{\eta_b \rightarrow \gamma\gamma}$ calculated with lattice QCD, (first column) and with several quark models.

Reference	[22]	[28]	[30]	[31]	[32]	[33]	[29]	[34]	[35]
$\Gamma_{\eta_b \rightarrow \gamma\gamma}$ (keV)	0.557	0.617	0.94	1.05	0.69	0.527	0.387	0.378	0.336

tions corresponding to the spin-spin, spin-orbit, and tensor interactions are added perturbatively. In Refs. [30–32] the authors use a relativized quark model. In Ref. [33] the authors incorporated the color-screening effect due to light quark-pair creation into the heavy quark-antiquark potential and investigated the effects of the screened potential on the bottomonium spectrum. In Ref. [34] the authors computed the mass spectra of bottomonium states using the instanton-induced potential obtained from the instanton Liquid Model for the QCD vacuum together with a confinement term. In Ref. [35] the authors employed a variational method with a single Gaussian-trial wave function with one parameter both in position space as well as momentum space to solve the Schrödinger equation. As can be seen in Table I, the predictions vary by a factor of 3 and hence a determination from first principles is very welcome.

As a first application of the recent lattice result, the first purpose of this work is to compute the η_b -production cross section through γ - γ interactions in ultraperipheral Pb - Pb collisions at the LHC energy $\sqrt{s} = 5.02$ TeV.

A second purpose of our work is to perform a comprehensive comparison of the ingredients used in this type of calculation: the practical definition of UPC, the method applied to obtain the equivalent photon flux, and the form factor of the photon source.

II. FORMALISM

In UPCs, the intense electromagnetic fields of heavy ions can be described within the Weizsäcker-Williams approximation as a flux of equivalent photons. In such collisions, the equivalent photons from one nucleus can then interact either with photons from the other nucleus or directly with the opposing nucleus. Here, we focus exclusively on photon-photon interactions, specifically on the exclusive production of η_b via the $\gamma\gamma \rightarrow \eta_b$ process (see Fig. 1). Moreover, when the nucleus acts as a coherent photon source, the photon flux is proportional to the square of the nuclear charge Z and the associated cross section is enhanced by a Z^4 factor. This leads to a large cross section at LHC energies.

A. The production cross section

Initially, let us present a brief review of the formalism needed to describe the pseudoscalar meson production in γ - γ interactions in ultraperipheral Pb - Pb collisions. In the equivalent photon approximation, the total cross section for the production of η_b can be factorized in terms of the equivalent photon flux of each of the nuclei and the photoproduction cross section as follows:

$$\sigma(PbPb \rightarrow Pb \otimes \eta_b \otimes Pb; s)$$

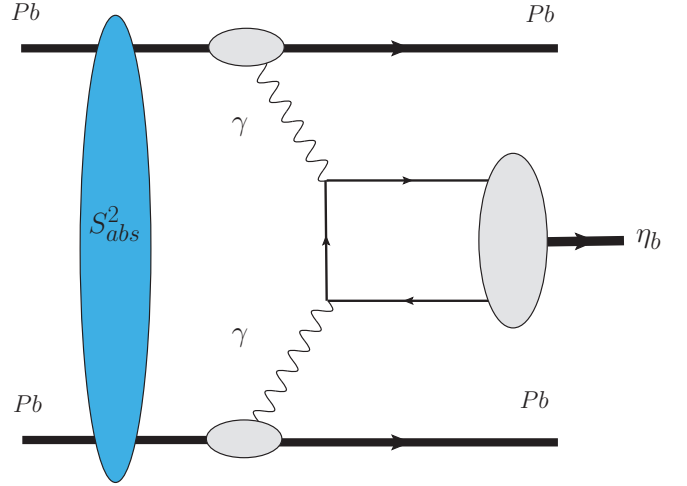


FIG. 1. Production of η_b by γ - γ interaction in ultraperipheral Pb - Pb collisions.

$$= \int \hat{\sigma}(\gamma\gamma \rightarrow \eta_b; W) N(\omega_1, \mathbf{b}_1) N(\omega_2, \mathbf{b}_2) \times S_{abs}^2(\mathbf{b}) d^2\mathbf{b}_1 d^2\mathbf{b}_2 d\omega_1 d\omega_2, \quad (1)$$

where $\sqrt{s}$ is the center-of-mass energy for the Pb - Pb collision, $\otimes$ characterizes a rapidity gap in the final state, $W = \sqrt{4\omega_1\omega_2}$ is the invariant mass of the γ - γ system, and $\hat{\sigma}(\gamma\gamma \rightarrow \eta_b)$ is the photoproduction cross section of η_b due to the fusion of two photons. Moreover, ω_i is the energy of the photon emitted by nucleus A_i at an impact parameter, or distance, b_i from A_i . The photons interact at the point P shown in Fig. 2. Note that, assuming the collinear approximation, the photon energies ω_1 and ω_2 are related to W and the rapidity Y of the produced

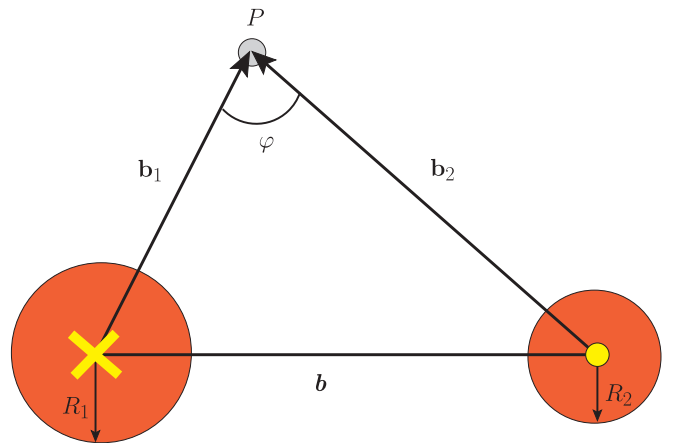


FIG. 2. View in the plane perpendicular to the direction of motion of the two ions. In a semiclassical picture the two equivalent photons of energy ω_1 and ω_2 collide in a point P with distance b_1 from ion 1 and b_2 from ion 2. The impact parameter $\mathbf{b}$ is the distance between the colliding nuclei with radius R_1 and R_2 .

state, through

$$\omega_1 = \frac{W}{2} e^{+Y} \quad \text{and} \quad \omega_2 = \frac{W}{2} e^{-Y}. \quad (2)$$

As a consequence, the total cross section can be rewritten as (for details see, e.g., Ref. [36])

$$\begin{aligned} \sigma(PbPb \rightarrow Pb \otimes \eta_b \otimes Pb; s) \\ = \int \hat{\sigma}(\gamma\gamma \rightarrow \eta_b; W) N(\omega_1, \mathbf{b}_1) N(\omega_2, \mathbf{b}_2) S_{abs}^2(\mathbf{b}) \\ \times \frac{W}{2} d^2\mathbf{b}_1 d^2\mathbf{b}_2 dW dY. \end{aligned} \quad (3)$$

Using the Low formula [37], we can express the cross section for the photon-photon interaction producing the pseudoscalar meson in terms of the two-photon decay width as follows:

$$\hat{\sigma}_{\gamma\gamma \rightarrow \eta_b}(\omega_1, \omega_2) = 8\pi^2(2J+1) \frac{\Gamma_{\eta_b \rightarrow \gamma\gamma}}{M_{\eta_b}} \delta(4\omega_1\omega_2 - M_{\eta_b}^2), \quad (4)$$

where M_{η_b} and J are the mass and spin of the produced η_b , respectively. The decay rate for $\eta_b \rightarrow \gamma\gamma$ was calculated in Ref. [22] using lattice QCD and was found to be $\Gamma(\eta_b \rightarrow \gamma\gamma) = 0.557(32)(1) \text{ keV}$.

In Eq. (3) the symbol $\otimes$ represents a rapidity gap. It is understood that, apart from the Pb projectiles which remain intact, nothing is observed except η_b , which is produced exclusively. However, the projectiles may become excited because of the photon emission. Nuclear excitation in UPCs has been discussed in several works. It is indeed possible to excite the nuclei in an UPC, with the consequent emission of fragments, mostly neutrons and extra photon production. The fragments would be produced at large rapidities. For recent works on the subject see Refs. [38,39] and references therein. This would destroy the exclusivity of the reaction and the photon production would generate additional background for the η_b signal. In the quoted works the contribution of these processes was found to be moderate, but a more detailed study remains to be done and would be quite relevant for η_b production.

B. The photon flux

The equivalent photon flux $N(\omega_i, b_i)$ of photons with energy ω_i at a transverse distance b_i from the center of the nucleus, defined in the plane transverse to the trajectory, can be expressed in terms of the charge form factor $F(q)$, where q is the four-momentum of the quasireal photon. It reads

$$\begin{aligned} N(\omega, b) = \frac{Z^2 \alpha}{\pi^2 \omega b^2} \left[\int_0^\infty u^2 J_1(u) F\left(\sqrt{\frac{u^2 + (b\omega/\gamma)^2}{b^2}}\right) \right. \\ \left. \times \frac{1}{u^2 + (b\omega/\gamma)^2} du \right]^2, \end{aligned} \quad (5)$$

where $\alpha = e^2/(4\pi)$, J_1 is the Bessel function of the first kind and γ is the Lorentz factor of the photon source ($\gamma = \sqrt{s}/2m_p$ and m_p is the proton mass). The form factor F is the Fourier transform of the spatial charge distribution. In the case of a nucleus-nucleus collision, it provides a description of the

charge distribution in the momentum transfer space of the nucleus. The nuclear charge distribution can be inferred from elastic-scattering data of electrons on nuclei and is often parametrized using the two-parameter Fermi model [40]. For spherical nuclei, the charge density is well described by the Woods-Saxon distribution

$$\rho(r) = \frac{\rho_0}{1 + \exp\left(\frac{r-R}{a}\right)}, \quad (6)$$

where ρ_0 is the central density, R the nuclear radius, and a the diffuseness parameter. However, an alternative description considers the convolution of a hard sphere with a Yukawa potential [41]. In this case, the form factor can be determined analytically as

$$F(q^2) = \frac{4\pi}{A q^3} \rho_0 [\sin(qR) - qR \cos(qR)] \left[\frac{1}{1 + q^2 a^2} \right], \quad (7)$$

with the parameters $a = 0.549 \text{ fm}$ and $\rho_0 = 0.1604/A \text{ fm}^{-3}$ obtained for the lead nucleus [42,43]. A simpler form factor, often used in the literature, is of the monopole type, given by

$$F(q^2) = \frac{\Lambda^2}{\Lambda^2 + q^2}, \quad \text{with} \quad \Lambda_{Pb} = 0.088 \text{ GeV}, \quad (8)$$

where Λ is a constant adjusted to reproduce the root-mean-square (rms) radius of a nucleus [36]. Introducing the monopole form factor into Eq. (5), we obtain the following expression for the photon flux

$$\begin{aligned} N(\omega, b) = \frac{Z^2 \alpha}{\pi^2 \omega} \left[\frac{\omega}{\gamma} K_1\left(\frac{b\omega}{\gamma}\right) - \sqrt{\left(\frac{\omega^2}{\gamma^2} + \Lambda^2\right)} K_1 \right. \\ \left. \times \left(b \sqrt{\frac{\omega^2}{\gamma^2} + \Lambda^2} \right) \right]^2. \end{aligned} \quad (9)$$

On the other hand, assuming a pointlike form factor ($F = 1$), the photon flux takes the following form:

$$N(\omega, b) = \frac{Z^2 \alpha}{\pi^2 \omega b^2} \left(\frac{\omega b}{\gamma} \right)^2 \left[K_1^2\left(\frac{\omega b}{\gamma}\right) + \frac{1}{\gamma^2} K_0^2\left(\frac{\omega b}{\gamma}\right) \right], \quad (10)$$

where K_0 and K_1 are the modified Bessel functions of the second kind. It is important to emphasize that this flux diverges at $b \rightarrow 0$. In this case, we need to take a lower-limit cut for the integrals over b . Usually, the integration is performed from a minimum distance $b_{\min} = R$ [15]. As demonstrated in Ref. [36], the realistic and monopole form factors are similar within a limited range of q and differ at large q . Additionally, the pointlike form factor is an unrealistic approximation, as it disregards the internal structure of the nucleus.

C. On the practical definition of an ultraperipheral collision

1. Pure geometry

The factor $S_{abs}^2(b)$ in Eq. (1) depends on the impact parameter of the collision and is called survival factor, which excludes the overlap between the colliding nuclei and selects only ultraperipheral collisions. A widely used procedure to exclude the strong interactions between the incident nuclei was proposed by Baur and Ferreira-Filho [44]. They assumed

that

$$S_{abs}^2(\mathbf{b}) = \Theta(|\mathbf{b}| - R_1 - R_2) = \Theta(|\mathbf{b}_1 - \mathbf{b}_2| - R_1 - R_2), \quad (11)$$

where R_i is the nuclear radius. In this model, the probability of having a hadronic interaction when $b > R_1 + R_2$ is zero. Though intuitive, this procedure introduces a sharp cut in the calculation and tends to overemphasize the role of geometry. It is based on an overly classical picture of nuclear collisions. In addition, it does not contain any energy dependence.

2. Geometry + dynamics

A more realistic treatment can be obtained using the survival factor $P_{NH}(b)$ which describes the probability of no additional hadronic interaction between the nuclei, which is usually estimated using the Glauber formalism. In this way we take into account the fact that, even when $b \gtrsim 2R$, the probability of having strong interactions is finite. In this formalism, $S_{abs}^2(b)$ can be expressed in terms of the probability of interaction between the nuclei at a given impact parameter $P_H(b)$, given by [45]

$$S_{abs}^2(\mathbf{b}) = P_{NH}(\mathbf{b}) = 1 - P_H(\mathbf{b}), \quad (12)$$

where

$$P_H(\mathbf{b}) = 1 - \exp \left[-\sigma_{nn} \int d^2\mathbf{r} T_A(\mathbf{r}) T_A(\mathbf{r} - \mathbf{b}) \right], \quad (13)$$

with σ_{nn} being the total cross section of the hadronic interaction and T_A the nuclear-thickness function. As in Ref. [45], we will assume that $\sigma_{nn} = 88$ mb at the LHC. This approach provides a semiclassical estimate of survival probability. It is interesting to observe that, since σ_{nn} grows with the energy, the interaction probability also grows with the energy for a fixed impact parameter. In principle, this might lead to big differences between the results obtained with Eqs. (11) and (12). However, σ_{nn} grows with the energy as $\propto \ln s$ or at most as $\propto \ln^2 s$ and the differences between the two approaches turn out to be modest.

In Fig. 3 we show the survival factor $S_{abs}^2(b)$ as a function of the impact parameter obtained with the geometric model (dashed line) and with the Glauber model (solid line) given by Eqs. (11) and (12), respectively. We also include the case of no absorption (dot-dashed line). Both absorption models use a nuclear radius of $R = 7.11$ fm, derived from the parametrization $R = 1.2A^{1/3}$ for the lead nuclei (^{208}Pb), where A is the mass number. As can be seen (see also Ref. [46]), the main difference between the absorption models is that the introduction of the survival factor given by Eq. (12) implies a smooth transition between the behavior of the small ($b < 2R$) and large ($b > 2R$) impact parameter. For comparison purposes, in what follows we will also consider the case without absorption effects, where $S_{abs}^2(b) = 1$.

In UPCs nuclei are sources of approximately real photons colliding with each other. Photon fusion can result in the production of particles in very clean events. These events are clean because the nuclei stay intact during the collision. When this is the case, we observe the particle produced in the central rapidity region and two rapidity gaps separating

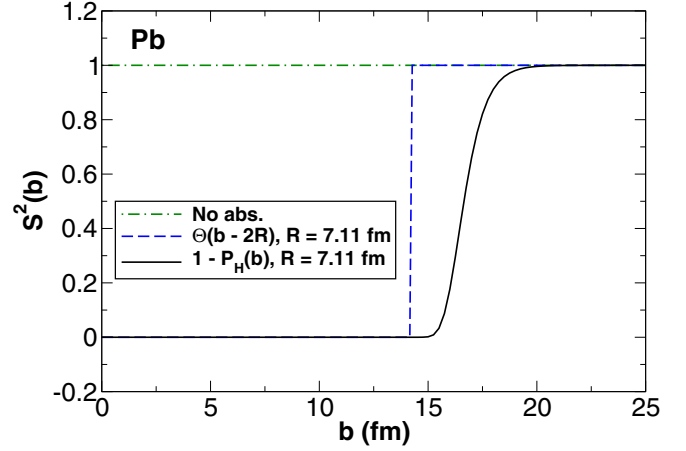


FIG. 3. Survival factor $S_{abs}^2(b)$ as a function of the impact parameter obtained with the geometric model of Eq. (11) (dashed line) and with the Glauber model of Eq. (12) (solid line). We also show the case of no absorption (dot-dashed line).

this particle from the beams. The finite size of the nuclei reduces the probability of this final state compared to pointlike particles. In order to address real measurements of the photon-photon fusion we have to include, in addition, the gap-survival probability. Using simple geometry to prevent the nuclei from overlapping, as implied by our Eq. (11), is a rough way to do this. Using our Eq. (12) is more realistic because we include the tails of the matter distribution and, in addition, we include the reaction cross section. This procedure deserves the name of semiclassical approach, since it takes the probabilistic nature of the interaction into account. However, these procedures do not fully take into account the structure of the nuclei and of the protons. The latter would be especially important in the production of large mass systems. As shown in Refs. [47,48], the inclusion of the hadron structure would lead to a more sophisticated theory of the survival factor. The treatment of absorption effects in photon-photon and photon-hadron interactions remains a subject of intense debate [45,49,50].

3. Kinematics

The absorptive factors, Eqs. (11) and (12), emphasize the geometrical and dynamical aspects involved in the operational definition of an UPC. There is yet a third way to define a UPC in terms of a kinematical constraint, as proposed in Ref. [51]. An ultraperipheral collision can also be defined in terms of the momentum of the photons involved in the interaction. The distribution of equivalent photons generated by a moving particle with the charge Ze is [1,2,4]

$$N(\mathbf{q}) = \frac{Z^2 \alpha}{\pi^2} \frac{(\mathbf{q}_\perp)^2}{\omega q^4} = \frac{Z^2 \alpha}{\pi^2 \omega} \frac{(\mathbf{q}_\perp)^2}{((\mathbf{q}_\perp)^2 + (\omega/\gamma)^2)^2}, \quad (14)$$

where q is the photon four-momentum, $\mathbf{q}_\perp$ is its transverse component, ω is the photon energy. To obtain the equivalent photon spectrum, one has to integrate this expression over the transverse momentum up to some value $\hat{q}$. After integrating over the photon transverse momentum, the equivalent

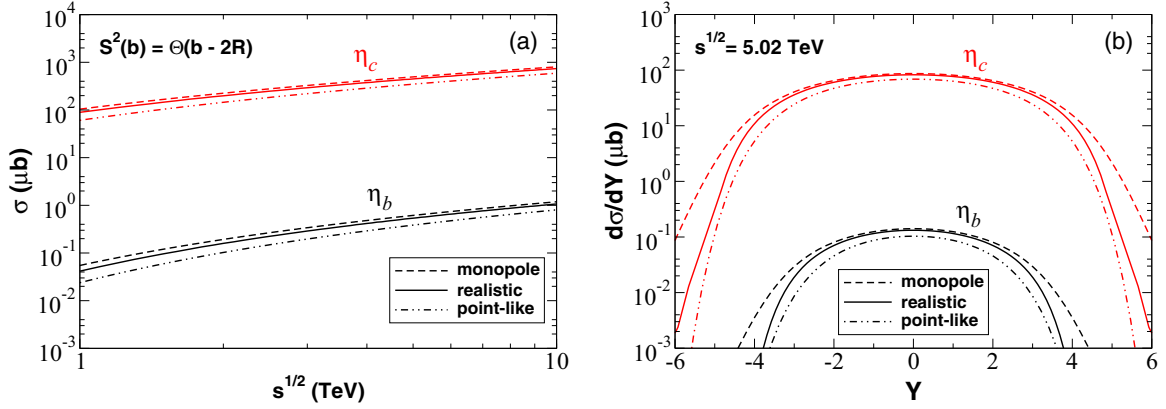


FIG. 4. (a) Energy dependence of the total cross section and (b) rapidity distribution of η_b produced in photon-photon collisions in ultraperipheral Pb - Pb collisions. The predictions for the η_c in the final state are also presented (red lines) for comparison.

photon-energy spectrum is given by

$$N(\omega) = \frac{2Z^2\alpha}{\pi} \ln\left(\frac{\hat{q}\gamma}{\omega}\right) \frac{1}{\omega}, \quad (15)$$

where the value of $\hat{q}$ must be chosen such that the particle emitting the photon does not break apart when emitting a photon with that momentum [51]. In an UPC, photon emission can be a coherent process, i.e., the photon is emitted by the whole source with a characteristic radius R . Incoherent emission can also take place. In this work, we focus exclusively on coherently emitted photons, for which the wavelength is at least of order R and we can interpret $\hat{q}$ as its maximum virtuality [52]. This gives us, in a first approximation, an estimate of $\hat{q} = \hbar c/R$. For Pb , $R \approx 7$ fm, and hence $\hat{q} \approx 0.028$ GeV. If $\hat{q}$ is larger than that, the photon starts to “resolve” the source and it might be emitted by a part of the source. In order to have an idea of the sensitivity of the results to this choice, we will take $\hat{q}$ to be in the range $0.014 < \hat{q} < 0.028$ GeV. These numbers are of the order of the binding energy of a nucleon in the nucleus. If $\hat{q}$ was larger, during the interaction the photon emission might induce the recoil of the emitting part of source (a nucleon) and cause its expulsion from the nucleus. In this case the final state would contain fragments and would not be equal to the initial state, contradicting the definition of an UPC. These arguments are very qualitative, but they set a scale for $\hat{q}$. As shown in Ref. [53], the dependence of the results on the choice of $\hat{q}$ is very strong and hence the predictions should be always given with a corresponding uncertainty. On the other hand, the value of $\hat{q}$ depends only on the photon source and, in this sense, it is universal, i.e., the same for a wide variety of photon-photon reactions with different final states. Therefore it can be determined studying measured reactions and then used to predict cross sections of yet-unobserved processes.

It is difficult to make rigorous predictions for $\hat{q}$. At the level of qualitative estimates, from what was said above, we should have $1/2R < \hat{q} < 1/R$. From an empirical point of view, in Ref. [51] the authors obtained the values of $\hat{q} \in [0.018, 0.030]$ GeV by fitting experimental data to the form factor of the ^{208}Pb nucleus. These values are consistent with previous

estimates and lie between $1/2R$ and $1/R$. Therefore, we will consider $\hat{q}$ in this range to make our predictions.

Since the photon flux (15) does not depend on the impact parameter, Eq. (1) simplifies to

$$\begin{aligned} \sigma(PbPb \rightarrow PbPb \eta_b) \\ = \int_{m_\eta^2/\hat{q}\gamma}^{\hat{q}\gamma} d\omega_1 \int_{m_\eta^2/\omega_1}^{\hat{q}\gamma} d\omega_2 \hat{\sigma}_{\gamma\gamma \rightarrow \eta_b}(\omega_1, \omega_2) N(\omega_1) N(\omega_2), \end{aligned} \quad (16)$$

where $m_\eta = m_{\eta_b}$.

In principle, the three definitions mentioned above are related to each other. They all refer, in one way or another, to the size of the source. We could even say that pure geometry and kinematics are related through a Fourier transform. In practice, however, these prescriptions are employed independently and it is useful to check whether they lead to equivalent results.

Before presenting the results, we would like to make a remark. The use of the coherence condition together with the uncertainty relation leads to the conclusion that the transverse momentum of the emitted photons (and hence of the produced meson) is small. It is possible to go beyond the qualitative estimates, compute the p_T distribution of the produced meson, and find that the average transverse momentum is indeed small. This was done, for example, in Refs. [54–56]. The knowledge of the η_b rapidity and p_T distributions will be crucial for its measurement.

III. RESULTS AND CONCLUSIONS

In Fig. 4(a) we present our predictions for the energy dependence of the total cross section for η_b production in ultraperipheral collisions Pb - Pb , obtained using the geometric absorption factor (11) and different models for the form factor. For comparison, we also include the predictions for the production of η_c .

In Fig. 4(b) we show the rapidity distribution of η_b produced in Pb - Pb collisions at $\sqrt{s} = 5.02$ TeV. As expected, the maximum of the distribution occurs at central rapidities, with the monopole (pointlike) predicting larger (smaller) values in

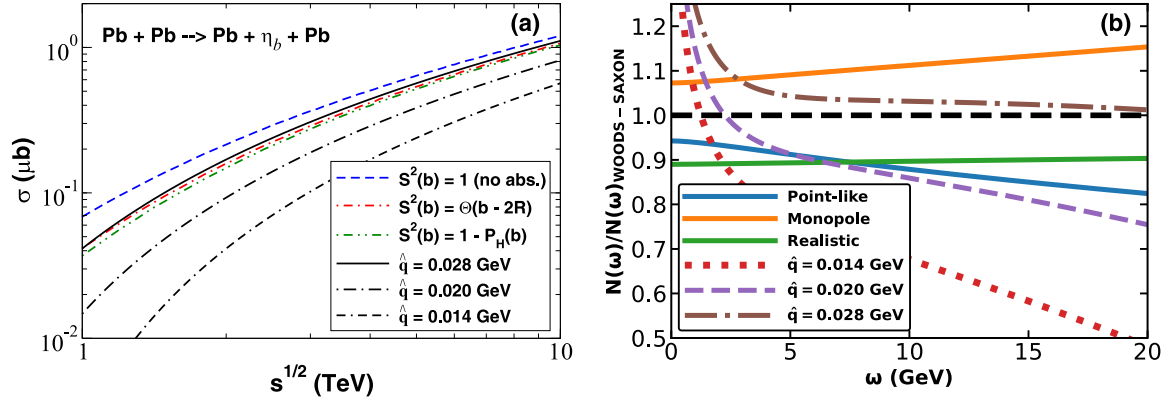


FIG. 5. (a) Energy dependence of the total cross section (in μb) for η_b production in Pb - Pb collisions calculated with Eq. (16) for three different values of $\hat{q}$. The blue, red, and green lines were obtained with Eq. (1) and with the realistic form factor (7). (b) Comparison between the photon flux ratios.

comparison to the more precise prediction derived using the realistic form factor. Moreover, the predictions for the η_b meson are characterized by smaller normalizations and narrower rapidity distributions than those for the η_c meson.

In Fig. 5(a) we show the energy dependence of the total cross section for η_b production in ultraperipheral collisions Pb - Pb calculated with Eq. (16) for three different values of $\hat{q}$. We also show the result obtained with Eq. (1), using the realistic form factor (7), combined with the geometric absorption factor (11) and the survival factor $P_{NH}(b)$ (12). For completeness, the prediction obtained disregarding the absorptive effects [$S^2_{abs}(b) = 1$] is also presented. Beyond the comparison of the final cross sections obtained with different approaches to UPCs, it is interesting to study the differences in photon flux, comparing the more usual definitions with the photon flux written in the momentum space for different values of $\hat{q}$. We would like to compare directly Eq. (5) with Eq. (15). Since these two fluxes have different dimensions, we will integrate Eq. (5) over the impact parameter. In order to make the differences more explicit, we will normalize all the fluxes by a standard one, which is obtained by integrating Eq. (5) over the impact parameter and choosing the form

factor F as the Fourier transform of the Woods-Saxon density. In Fig. 5(b) we show this ratio. In the energy range of interest the fluxes obtained with Eq. (5) differ at most by 15%. The fluxes obtained with Eq. (15) have a strong dependence on the choice of $\hat{q}$. On the other hand, once the “right” choice is made, this flux is the closest to the one obtained with the Woods-Saxon density. The singular behavior observed at low energies does not have a big impact on the calculation of the cross sections because of the production-threshold condition in Eq. (16).

As can be seen in Eqs. (3) and (4), the decay width enters as a multiplication factor in the expression for the cross section. The uncertainty in the lattice QCD determination of $\Gamma_{\eta_b \rightarrow \gamma\gamma}$ adds an uncertainty of about 5.8% to the cross section. This affects only the normalization of the predictions. On the other hand, the systematic error coming from the choice of the form factor is much larger than that. As can be seen in Table II, it can be of the order of 10–30%.

Having shown the energy dependence, let us focus on $\sqrt{s} = 5.02$ TeV, which is the most interesting case and where we may have data. In Table II we show all the possible ways to compute the cross section with the different ingredients.

TABLE II. Total cross sections (in μb) for η_b and η_c production in γ - γ interactions in ultraperipheral Pb - Pb collisions at $\sqrt{s} = 5.02$ TeV (with $\gamma = 2675.9$).

Meson	$\Gamma_{\gamma\gamma}$ [KeV]	Absorptive factor	σ (μb)		
$\eta_b(1S)$	0.557	$S^2_{abs}(\mathbf{b}) = 1 \rightarrow \text{no abs.}$ $S^2_{abs}(\mathbf{b}) = \Theta(\mathbf{b} - 2R)$ $S^2_{abs}(\mathbf{b}) = 1 - P_H(\mathbf{b})$	pointlike	monopole	realistic
			0.43	0.76	0.64
			0.39	0.62	0.54
		$\hat{q}$	0.38	0.59	0.52
			$\hat{q} = 0.014$ GeV	$\hat{q} = 0.020$ GeV	$\hat{q} = 0.028$ GeV
			0.24	0.39	0.57
$\eta_c(1S)$	5.152	$S^2_{abs}(\mathbf{b}) = 1 \rightarrow \text{no abs.}$ $S^2_{abs}(\mathbf{b}) = \Theta(\mathbf{b} - 2R)$ $S^2_{abs}(\mathbf{b}) = 1 - P_H(\mathbf{b})$	pointlike	monopole	realistic
			368.62	557.39	497.03
			348.66	495.25	451.22
		$\hat{q}$	343.72	485.04	441.94
			$\hat{q} = 0.014$ GeV	$\hat{q} = 0.020$ GeV	$\hat{q} = 0.028$ GeV
			259.65	355.74	465.77

TABLE III. Total cross sections (in μb) for η_b and η_c production in $\gamma\text{-}\gamma$ interactions obtained in previous works.

Meson	$\Gamma_{\gamma\gamma}$ [KeV]	σ (μb)		
		$Z = 82, \gamma = 2750$	$Z = 82, \gamma = 3000$	$Z = 82, \gamma = 4000$
$\eta_b(1S)$	0.40	0.322 [57] 0.466 [58]	0.445 [59]	
	0.41			0.46 [44] 0.50 [57] 0.90 [60]
	0.43	0.346 [57] 0.37 [52]	0.32 [61]	
	6.30	464 [57] 598 [58]	603 [59]	590 [44] 644 [57] 800 [60]
	7.50	552 [57] 590 [52]	540 [61] 610 [43]	
$\eta_c(1S)$				

As expected, the upper limit of the cross sections is set by $S_{abs}^2(b) = 1$, i.e., when there is no condition enforcing the ultraperipheral nature of the collision. For all absorption factors, the pointlike and monopole form factors set, respectively, the lower and upper limit of the cross section. The results obtained with the realistic form factor (7) lie in between. Since the predictions obtained with the geometric and geometric + dynamic absorption factor nearly coincide, we can say that the most likely value of the cross section is 0.52–0.54 microbarns. The error can be estimated from Table II, since the different choices of the form factors give a good representation of the uncertainty. These values are compatible with those found using the $\hat{q}$ prescription. The latter present a larger uncertainty. Taken together, these numbers show a nice convergence to the value of the cross section mentioned above, which is, in principle, large enough to be observed.

In Table III we present a compilation of previous estimates of the same quantities. These cross sections have been calculated by several authors. In particular, the formalism used in Refs. [44] and [57] is quite similar to the one used here.

Comparing the results shown in Tables II and III we observe a general consistency within the theoretical uncertainties. Small discrepancies in the cross-section values can be attributed to the different levels of precision used in the treatment of the nuclear form factor and the models for the absorption effects. On the other hand, the least-known ingredient was $\Gamma(\eta_b \rightarrow \gamma\gamma)$ and in the quoted previous works it was assumed to be $\simeq 0.41$ keV. The lattice QCD value is $\simeq 0.55$ keV, i.e., 34% larger and approximately so are the corresponding cross sections. It is important to emphasize that when we redo our calculations using $\Gamma(\eta_b \rightarrow \gamma\gamma) = 0.41$ keV and the appropriate γ factor, we faithfully reproduce the values of the previous predictions. For completeness, in Table IV we present a brief study of how the cross section of η_b production changes when we vary the Woods-Saxon parameters. As can be seen, the uncertainties in these parameters introduce an additional uncertainty of the order of 5% in the cross section.

In Tables II and III we have included the cross sections for η_c production. This is interesting for a comparison with η_b and also because very recently, the BESIII collaboration

performed a direct measurement of the decay of η_c into two photons in the process $J/\psi \rightarrow \gamma\eta_c$. They reported a determination of the product branching fraction $\mathcal{B}(J/\psi \rightarrow \gamma\eta_c) \times \mathcal{B}(\eta_c \rightarrow \gamma\gamma) = (5.23)(40) \times 10^{-6}$ [62]. Using $\mathcal{B}(J/\psi \rightarrow \gamma\eta_c) = (1.41)(14)\%$ and the total decay width $\Gamma_{\eta_c}^{\text{total}} = (30.5)(5)$ MeV from the PDG [63], the partial decay width of $\eta_c \rightarrow \gamma\gamma$ is calculated as $\Gamma(\eta_c \rightarrow \gamma\gamma) = (11.30)(1.43)$ keV, which is more than twice as large as the current reported world average for this decay. Using this new width in our calculations, combined with the monopole form factor and the geometric absorption factor, we obtain a total cross section of $\sigma = 1086.24 \mu\text{b}$ for η_c production in ultraperipheral $Pb\text{-}Pb$ collisions at $\sqrt{s} = 5.02$ TeV. This result highlights the strong sensitivity of the cross section to the η_c two-photon decay width.

The exclusive η_b production studied here has a very clean signature, which is the presence of two rapidity gaps (regions with no particle production) in the final state separating the $b\bar{b}$ state from the intact outgoing nuclei. These events can be separated at the LHC by measuring the outgoing nuclei with the help of the forward proton detectors (FPD) such as the ATLAS Forward Proton detector (AFP) [64,65] and the Precision Proton Spectrometer (CT-PPS) [66]. To know how often η_b is produced, it is crucial to know the number of useful interactions (events). The quantity that measures the ability of an accelerator to produce these interactions is the luminosity, which is the proportionality factor between the number of

TABLE IV. Dependence of the η_b production cross section on the Woods-Saxon parameters R and a . For details, see Ref. [42], page 16.

	R [fm]	a [fm]	σ_{η_b} [μb]
(R, a)	6.624(35)	0.549(8)	$0.541^{+0.027}_{-0.035}$
$(R + \Delta R, a)$	6.659	0.549	0.568
$(R - \Delta R, a)$	6.589	0.549	0.506
$(R, a + \Delta a)$	6.624	0.557	0.540
$(R, a - \Delta a)$	6.624	0.541	0.540

TABLE V. Cross sections, luminosity and number of η_b for Pb - Pb collisions at different energies.

Energy	Cross section [μb]	Luminosity [$\text{cm}^{-2}\text{s}^{-1}$]	Number of η_b [s^{-1}]
5.00 TeV (ATLAS 2018)	0.539	6.2×10^{27} [69]	3.34×10^{-3}
5.20 TeV (ALICE 2017)	0.559	6×10^{27} [70]	3.35×10^{-3}
2.76 TeV (ALICE 2014)	0.258	1×10^{26} [71]	2.58×10^{-5}

events per second and the cross section [67]:

$$\frac{dR}{dt} = \mathcal{L} \cdot \sigma. \quad (17)$$

The event rate can be estimated from the above equation. By integrating this expression over time, we obtain the total number of η_b mesons produced as a function of the integrated luminosity. Table V presents the integrated luminosity values provided by the ATLAS and ALICE collaborations for different collision energies, from which we estimate the expected number of η_b mesons. Taking the current Pb - Pb luminosity at the LHC [68] to be $\mathcal{L} = 6.4 \times 10^{27} \text{ cm}^{-2} \text{ s}^{-1}$, we predict that the associated number of events per second will be $3.4 \times 10^{-3} \text{ s}^{-1}$ for η_b production. This gives 10 000 events in approximately 34 days. This is the number of events needed to achieve a precise detection of a particle such as the η_c [72].

One of the cleanest ways to measure η_b is through its decay into two photons with low transverse momentum in the central rapidity region. The electromagnetic calorimetry of ALICE, CMS, and LHCb provides the resolution necessary to separate the two-photon decay of the η_b from background processes. Given the large number of events per second mentioned in the last paragraph, the experimental analysis of η_b production and decay into two photons is, in principle, feasible.

The formalism described above can be easily extended to ultraperipheral proton-nucleus (pA) collisions by replacing one of the nuclear-photon fluxes by the flux originating from the proton. Since the coherent flux scales as Z^2 , the resulting cross sections in p - Pb collisions are expected to be smaller than in Pb - Pb . In this case, the elastic form factor of the

proton must be used and the absorptive factor adapted to the asymmetric system, typically by imposing the condition $b > R_p + R_A$. Moreover, the rapidity distribution becomes asymmetric, reflecting the very different photon fluxes of the proton and the nucleus.

IV. SUMMARY

To summarize, we have updated the estimate of the η_b production cross section in ultraperipheral Pb - Pb collisions at $\sqrt{s} = 5.02 \text{ TeV}$ per nucleon pair. The two main improvements were the introduction of a very recently calculated value of $\Gamma(\eta_b \rightarrow \gamma\gamma)$ and also the systematic comparison of different prescriptions to treat the absorptive effects. The obtained values of the η_b production cross section are larger than the previous ones but within the same order of magnitude. Finally, it is reassuring to observe the three central lines in Fig. 5(a) and realize that the three ways to define an UPC lead to very similar results.

ACKNOWLEDGMENTS

This work was partially financed by the Brazilian funding agencies CNPq, CAPES, FAPESP, and INCT-FNA (Contract No. 464898/2014-5). The authors gratefully acknowledge the support from the Fundação de Amparo à Pesquisa do Estado de São Paulo (FAPESP), Contracts No. 2023/14456-8 and No. 2024/19103-9.

DATA AVAILABILITY

The data that support the findings of this article are openly available [15].

-
- [1] C. A. Bertulani and G. Baur, *Phys. Rep.* **163**, 299 (1988).
 - [2] C. A. Bertulani, S. R. Klein, and J. Nystrand, *Ann. Rev. Nucl. Part. Sci.* **55**, 271 (2005).
 - [3] V. P. Goncalves and M. V. T. Machado, *J. Phys. G* **32**, 295 (2006).
 - [4] A. J. Baltz, G. Baur, D. d'Enterria, L. Frankfurt, F. Gelis, V. Guzey, K. Hencken, Y. Kharlov, M. Klasen, S. R. Klein *et al.*, *Phys. Rep.* **458**, 1 (2008).
 - [5] J. G. Contreras and J. D. Tapia Takaki, *Int. J. Mod. Phys. A* **30**, 1542012 (2015).
 - [6] S. Klein and P. Steinberg, *Annu. Rev. Nucl. Part. Sci.* **70**, 323 (2020).
 - [7] C. Adler *et al.* (STAR Collaboration), *Phys. Rev. Lett.* **89**, 272302 (2002).
 - [8] S. Afanasiev *et al.* (PHENIX Collaboration), *Phys. Lett. B* **679**, 321 (2009).
 - [9] B. Abelev *et al.* (The ALICE Collaboration), *Phys. Lett. B* **718**, 1273 (2013).
 - [10] E. Abbas *et al.* (The ALICE Collaboration), *Eur. Phys. J. C* **73**, 2617 (2013).
 - [11] V. Khachatryan *et al.* (The CMS Collaboration), *Phys. Lett. B* **772**, 489 (2017).
 - [12] M. Dyndal *et al.* (ATLAS Collaboration), *Nucl. Phys. A* **967**, 281 (2017).
 - [13] B. D. Moreira, C. A. Bertulani, V. P. Goncalves, and F. S. Navarra, *Phys. Rev. D* **94**, 094024 (2016).
 - [14] C. Azevedo, V. P. Gonçalves, and B. D. Moreira, *Phys. Rev. C* **101**, 024914 (2020).
 - [15] C. A. Bertulani, *Phys. Rev. C* **79**, 047901 (2009); R. Fariello, D. Bhandari, C. A. Bertulani, and F. S. Navarra, *ibid.* **108**, 044901 (2023).
 - [16] S. Klein and J. Nystrand, *Phys. Today* **70**, 40 (2017).

- [17] V. P. Gonçalves and B. D. Moreira, *Phys. Rev. D* **97**, 094009 (2018).
- [18] C. A. Bertulani, *J. Phys.: Conf. Ser.* **2619**, 012003 (2023).
- [19] C. Bertulani, *Phys. Proc. UPC* **1**, 1 (2024).
- [20] J. D. Jackson, *Classical Electrodynamics* (Wiley, New York, (1975).
- [21] V. M. Budnev, I. F. Ginzburg, G. V. Meledin, and V. G. Serbo, *Phys. Rep.* **15**, 181 (1975).
- [22] B. Colquhoun, C. T. H. Davies, and G. P. Lepage, *arXiv:2410.24041*.
- [23] A. Garmash, *Int. J. Mod. Phys. A* **36**, 2044018 (2021); R. Mizuk, *EPJ Web Conf.*, **212**, 02005 (2019).
- [24] T. Hilger, C. Popovici, M. Gomez-Rocha, and A. Krassnigg, *Phys. Rev. D* **91**, 034013 (2015); C. S. Fischer, S. Kubrak, and R. Williams, *Eur. Phys. J. A* **51**, 10 (2015).
- [25] N. Brambilla, A. Pineda, J. Soto, and A. Vairo, *Rev. Mod. Phys.* **77**, 1423 (2005); A. Pineda, *Prog. Part. Nucl. Phys.* **67**, 735 (2012).
- [26] Z. G. Wang, *Front. Phys.* **21**, 016300 (2026); R. M. Albuquerque, J. M. Dias, K. P. Khemchandani, A. Martínez Torres, F. S. Navarra, M. Nielsen, and C. M. Zanetti, *J. Phys. G* **46**, 093002 (2019).
- [27] T. Burch, C. DeTar, M. Di Pierro, A. X. El-Khadra, E. D. Freeland, S. Gottlieb, A. S. Kronfeld, L. Levkova, P. B. Mackenzie, and J. N. Simone, *Phys. Rev. D* **81**, 034508 (2010); A. Gray, I. Allison, C. T. H. Davies, E. Dalgic, G. P. Lepage, J. Shigemitsu, M. Wingate, HPQCD, and UKQCD, *ibid.* **72**, 094507 (2005).
- [28] C. M. A. and R. Dhir, *J. Phys. G* **51**, 115004 (2024).
- [29] N. R. Soni, B. R. Joshi, R. P. Shah, H. R. Chauhan, and J. N. Pandya, *Eur. Phys. J. C* **78**, 592 (2018).
- [30] S. Godfrey and K. Moats, *Phys. Rev. D* **92**, 054034 (2015).
- [31] J. Z. Wang, Z. F. Sun, X. Liu, and T. Matsuki, *Eur. Phys. J. C* **78**, 915 (2018).
- [32] J. Segovia, P. G. Ortega, D. R. Entem, and F. Fernandez, *Phys. Rev. D* **93**, 074027 (2016).
- [33] B. Q. Li and K. T. Chao, *Commun. Theor. Phys.* **52**, 653 (2009).
- [34] B. Pandya, M. Shah, and P. C. Vinodkumar, *Eur. Phys. J. C* **81**, 116 (2021).
- [35] V. Kher, R. Chaturvedi, N. Devlani, and A. K. Rai, *Eur. Phys. J. Plus* **137**, 357 (2022).
- [36] M. Klusek-Gawenda and A. Szczurek, *Phys. Rev. C* **82**, 014904 (2010).
- [37] F. E. Low, *Phys. Rev.* **120**, 582 (1960).
- [38] P. Jucha, M. Klusek-Gawenda, and A. Szczurek, *Phys. Rev. D* **112**, L051501 (2025).
- [39] M. Klusek-Gawenda, V. P. Goncalves, and A. Szczurek, *Phys. Lett. B* **867**, 139614 (2025).
- [40] H. De Vries, C. W. De Jager, and C. De Vries, *Atom. Data nucl. Data tabl.* **36**, 495 (1987).
- [41] K. T. R. Davies and J. R. Nix, *Phys. Rev. C* **14**, 1977 (1976).
- [42] C. W. De Jager, H. De vries, and C. De vries, *Atom. Data nucl. Data tabl.* **14**, 479 (1974); **16**, 580(E) (1975).
- [43] C. A. Bertulani and F. Navarra, *Nucl. Phys. A* **703**, 861 (2002).
- [44] G. Baur and L. G. Ferreira Filho, *Nucl. Phys. A* **518**, 786 (1990).
- [45] A. J. Baltz, Y. Gorbunov, S. R. Klein, and J. Nystrand, *Phys. Rev. C* **80**, 044902 (2009).
- [46] C. Azevedo, V. P. Gonçalves, and B. D. Moreira, *Eur. Phys. J. C* **79**, 432 (2019).
- [47] S. I. Godunov, V. A. Novikov, A. N. Rozanov, M. I. Vysotsky, and E. V. Zhemchugov, *J. High Energy Phys.* **10** (2021) 234.
- [48] M. Luszczak, W. Schaefer, and A. Szczurek, *J. High Energy Phys.* **05** (2018) 064; L. Forthomme, M. Luszczak, W. Schaefer, and A. Szczurek, *Phys. Lett. B* **789**, 300 (2019); M. Luszczak, L. Forthomme, W. Schaefer, and A. Szczurek, *J. High Energy Phys.* **02** (2019) 100.
- [49] M. Dyndal and L. Schoeffel, *Phys. Lett. B* **741**, 66 (2015).
- [50] V. A. Khoze, A. D. Martin, and M. G. Ryskin, *J. Phys. G* **45**, 053002 (2018).
- [51] M. I. Vysotskii and E. Zhemchugov, *Phys. Usp.* **62**, 910 (2019).
- [52] G. Baur, K. Hencken, and D. Trautmann, *J. Phys. G* **24**, 1657 (1998).
- [53] F. C. Sobrinho, L. M. Abreu, C. A. Bertulani, and F. S. Navarra, *Phys. Rev. D* **110**, 034037 (2024).
- [54] K. Hencken, G. Baur, and D. Trautmann, *Phys. Rev. Lett.* **96**, 012303 (2006).
- [55] S. R. Klein and J. Nystrand, *Phys. Rev. Lett.* **84**, 2330 (2000).
- [56] S. R. Klein and J. Nystrand, *Phys. Rev. C* **60**, 014903 (1999).
- [57] K. A. Chikin, V. L. Korotkih, A. P. Kryukov, L. I. Sarycheva, I. A. Pshenichnov, J. P. Bondorf, and I. N. Mishustin, *Eur. Phys. J. A* **8**, 537 (2000).
- [58] F. Krauss, M. Greiner, and G. Soff, *Prog. Part. Nucl. Phys.* **39**, 503 (1997).
- [59] M. Vidovic, M. Greiner, and G. Soff, *J. Phys. G* **21**, 545 (1995).
- [60] G. Baur and C. A. Bertulani, *Nucl. Phys. A* **505**, 835 (1989).
- [61] G. Baur, K. Hencken, D. Trautmann, S. Sadovsky, and Y. Kharlov, *Phys. Rept.* **364**, 359 (2002).
- [62] M. Ablikim *et al.* (BESIII Collaboration), *Phys. Rev. Lett.* **134**, 181901 (2025).
- [63] R. L. Workman *et al.* [Particle Data Group], *PTEP* **2022**, 083C01 (2022).
- [64] L. Adamczyk *et al.*, CERN-LHCC-2015-009, ATLAS-TDR-024.
- [65] M. Tasevsky *et al.* (ATLAS Collaboration), *AIP conf. Proc.* **1654**, 090001 (2015).
- [66] M. Albrow *et al.* (CMS and TOTEM collaborations), CERN-LHCC-2014-021, TOTEM-TDR-003, CMS-TDR-13.
- [67] W. Herr, *arXiv:1601.07311*.
- [68] R. Bruce, M. A. Jebramcik, J. M. Jowett, T. Mertens, and M. Schaumann, *Eur. Phys. J. Plus* **136**, 745 (2021).
- [69] ATLAS experiment, <http://twiki.cern.ch/twiki/bin/view/AtlasPublic/LuminosityPublicResultsRun2>.
- [70] Virginia Greco, <http://alicematters.Web.Cern.Ch/?Q=content/node/998>.
- [71] B. B. Abelev *et al.* (The ALICE Collaboration), *Int. J. Mod. Phys. A* **29**, 1430044 (2014).
- [72] R. Aaij *et al.* (LHCb collaboration), *Eur. Phys. J. C* **80**, 191 (2020).