Dibaryon resonance and two-photon bremsstrahlung in *pp* scattering

O. Scholten* and Ojwang J. G. O.
Kernfysisch Versneller Instituut, University of Groningen, 9747 AA Groningen, The Netherlands

S. Tamenaga
Research Center for Nuclear Physics (RCNP), Osaka University, Ibaraki, Osaka 567-0041, Japan
(Received 11 January 2005; published 23 March 2005)

The signature of a possible dibaryon state is studied in two-photon production in proton-proton scattering, that is, in the process $pp \to D\gamma \to 2\gamma NN$. We have investigated the effects on the cross section for cases in which the dibaryon has spin $0^{\pm}$. Special attention is paid to interference of the dibaryon signal with that of the two-photon bremsstrahlung process.

DOI: 10.1103/PhysRevC.71.034005 PACS number(s): 13.75.Cs, 13.30.Ce, 13.40.Hq, 25.10.+s

I. INTRODUCTION

Various QCD-inspired models have predicted the existence of a dibaryon state. The most well known is the H dibaryon with strangeness $S = -2$. In this work we will concentrate, however, on nonstrange dibaryons for which there are predictions based in quark models (see, for example, [1–3]) as well as in a potential model approach [4]. For such a dibaryon state (D) to exist and not be observed in proton-proton (pp) or proton-neutron (pn) scattering requires that the quantum numbers differ from the pp and pn partial waves; that is, they must be so-called decoupled states. As emphasized by Gerasimov *et al.* [5,6] the most promising way to observe a dibaryon with a mass below the pion-production threshold is through electromagnetic decay. The JINR phasotron collaboration claims to have seen the dibaryon resonance [7], using a proton beam of 216 MeV in the double-bremsstrahlung reaction $pp \to pp\gamma\gamma$. Its mass was determined at $m_D = 1956 \pm 6$ MeV and its width $W_D \leqslant 8$ MeV. There also exists some evidence that a dibaryon state has been excited in pion scattering off a nucleus [8]. It should be noted that in Ref. [9] it is contested that in this reaction the dibaryon is excited with an appreciable cross section. An important process contributing to the $pp \to pp\gamma\gamma$ reaction is two-photon bremsstrahlung. In previous estimates of the effects of the dibaryon, the two-photon bremsstrahlung contribution has not, however, been taken into account.

To investigate the effect of a possible dibaryon resonance on the cross section of the $pp \to pp\gamma\gamma$ reaction we have added coherently the amplitude for bremsstrahlung to that for the dibaryon process in the present work. A model for the bremsstrahlung process, which obeys the low-energy constraints, has been proposed in [10]. The dibaryon contribution is included assuming it to be a $(J^P, T) = (0^{\pm}, 2)$ state with mass $M_D = 1956$ MeV. A strong interference is observed where the pattern clearly distinguishes the quantum numbers of the state.

II. TWO-PHOTON EMISSION

The process of two-photon bremsstrahlung gives an important contribution to the cross section for $pp \to pp\gamma\gamma$. To a large extent the matrix element for this process is governed by the low- (photon) energy theorem (or soft-photon theorem). Reference [10] presents a derivation of a $pp\gamma\gamma$ amplitude that obeys the low-energy theorem (i.e., is gauge invariant, free of spurious poles, and based on a realistic pp-scattering amplitude). The Feynman diagrams included in the matrix element are given in Fig. 1. As described in [10] a Taylor series expansion around a point of average kinematics is used for the momentum dependence of the T matrix, which guarantees that current conservation is obeyed for proton-proton bremsstrahlung. The photon vertex includes contributions from the anomalous magnetic moment, which are very important in pp bremsstrahlung. The model is covariant. This model, when applied to one-photon bremsstrahlung, gives a good reproduction of the cross section obtained in a recent high-statistics experiment at 190 MeV [11]. This also shows that terms that are beyond those determined by the low-energy theorem are adequately accounted for.

A. The dibaryon mechanism

The possible quantum numbers for a decoupled dibaryon that can be excited in pp scattering through coupling with a photon are $(J^P, T) = (1^+, 1)$ or any $T = 2$ state. For the present investigation we will assume a spin-zero, $T = 2$ state and consider both positive and negative parity.

The structure of the four-point vertex for the coupling of the dibaryon to the proton-proton-photon channel can be obtained in several ways. The simplest structure that obeys the necessary symmetries is

$$\Gamma^{\mu}_{pp\gamma D} = \frac{eg}{m^2}\sigma^{\mu\rho}q_{\rho}\Gamma C, \tag{1}$$

where the photon momentum is given by q, C is the charge conjugation operator, and $\Gamma = 1$ ($\Gamma = \gamma^5$) for a $J^{\pi} = 0^-$ ($J^{\pi} = 0^+$) dibaryon, respectively. One may also obtain an expression for the vertex assuming a sequential process in which the photon first couples the proton to form a spin-1/2

*Electronic address: scholten@kvi.nl.

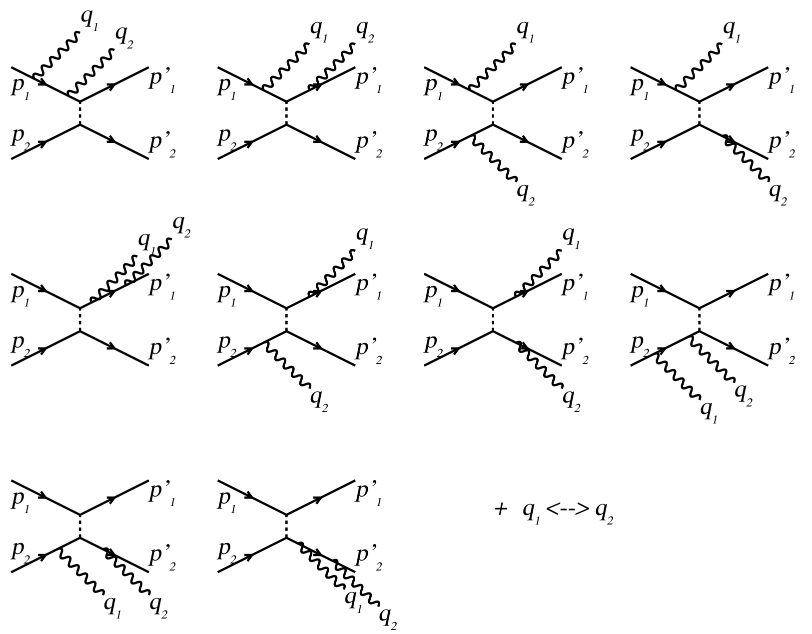

FIG. 1. The Feynman diagrams included in the calculation of two-photon bremsstrahlung.

baryon, which subsequently couples with the other proton to form the dibaryon. One can show that in the limit of an infinitely heavy intermediate baryon the effective vertex reduces to the expression given in Eq. (1). The same result also holds if the intermediate propagator is taken to be spin $3/2$ instead of $1/2$. Since the vertex given in Eq. (1) is magnetic in structure, gauge invariance is automatically obeyed and it is sufficient to include the Feynman diagrams given in Fig. 2 to account for the contribution of the dibaryon. In particular it should be noticed that the diagrams where both photons are emitted before or after the formation of the dibaryon do not contribute because the quantum numbers prevent the dibaryon from coupling directly to the initial or final two-proton state through the strong interaction.

The matrix element for the contribution of an intermediate dibaryon state to two-photon production in proton-proton scattering can now be written as

$$\mathcal{M}_D^{\mu\nu} = e^2 g^2 \bar{u}_1' \Gamma^{\mu}_{pp\gamma D} \bar{u}_2'^T D(p_D) u_1^T \Gamma^{\nu}_{pp\gamma D} u_2, \tag{2}$$

where the spinors of the ingoing (outgoing) nucleons are denoted by u_i ($\bar{u}_i$), $i = 1, 2$, respectively. The propagator of the dibaryon with momentum p_D and mass m_D is

$$D(p_D) = \frac{1}{p_D^2 - m_D^2 + i m_D W_D}\ . \tag{3}$$

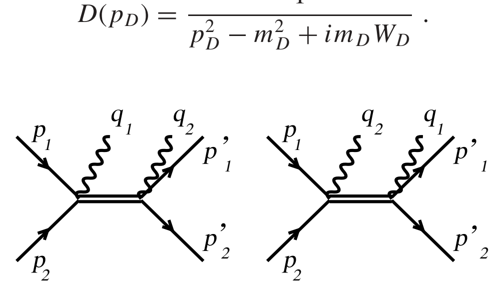

FIG. 2. Feynman diagrams included in the calculation of the dibaryon mechanism.

The decay width of the dibaryon due to the coupling to the $pp\gamma$ channel is W_D and is calculated in the next section.

B. Dibaryon decay width

In general the decay width of a particle with mass m_A can be expressed as

$$dW_{(A\to 1\cdots n)} = g^2 |\mathcal{M}|^2 \frac{dQ}{2m_A}, \tag{4}$$

where the multiparticle phase space is $dQ = (2\pi)^4 \delta^4(P) \prod \frac{dp_i^3}{(2\pi)^3 2E_i}$.[1] For a dibaryon decaying into two protons (p_a, p_b) and a photon (q), this phase space can be written as

$$\begin{aligned} dQ &= \frac{d^3 p_a}{2E_a} \frac{d^3 p_b}{2E_b} \frac{d^3 q}{2\omega} \frac{1}{2(2\pi)^5} \\ &\quad \times \delta(M_D - E_a - E_b - \omega) \delta^3(\vec{p}_a + \vec{p}_b + \vec{q}) \\ &= \int \frac{4}{(2\pi)^5} \frac{1}{8 E_a E_b \omega} x^2 dx d\Omega_x y^2 dy d\Omega_y \\ &\quad \times \delta(M - E_a - E_b - \omega), \end{aligned} \tag{5}$$

where $\vec{x} = (\vec{p}_a + \vec{p}_b)/2$, $\vec{y} = (\vec{p}_a - \vec{p}_b)/2$, $E_a = \sqrt{m^2 + \vec{p}_a^{\,2}}$, $E_b = \sqrt{m^2 + \vec{p}_b^{\,2}}$, $\omega = |\vec{q}|$, $x = |x|$, and $y = |y|$. Two identical protons give rise to a factor $1/2$. Integration over Ω_x gives a factor 4π and $\int d\phi_y = 2\pi$ since for a spin-zero dibaryon the integrand is independent of these variables.

The δ function in energy can be used to cancel the y integral, giving

$$y^2 = (M_D - 2x)^2 \frac{(M_D - 2x)^2 - 4(m^2 + x^2)}{4(M_D - 2x)^2 - 16x^2 \cos\theta_y^2}. \tag{6}$$

[1] This assumes we take $u\bar{u} = \not{p} + m$ for fermions in Eq. (10).

The upper limit (x_u) of x integration is obtained when $y = 0$,

$$x_u = \left(M_D^2 - 4m^2\right)/4M_D. \tag{7}$$

For $M_D \leqslant 4m$ the boundaries for the $\cos\theta_y$ integration are simply -1 and $+1$.

The phase space now reads

$$\begin{aligned} dQ &= dx d\cos\theta_y \frac{2(4\pi)(2\pi)}{(2\pi)^5} \frac{x^2 y^2}{4E_a E_b (2x)} \\ &\quad \times \left(\frac{y + x\cos\theta_y}{E_a} + \frac{y - x\cos\theta_y}{E_b}\right)^{-1} \\ &= \frac{1}{2} dx^2 d\cos\theta_y\, J, \end{aligned} \tag{8}$$

where

$$J = \frac{y^2}{2(2\pi)^3 (E_b(y + x\cos\theta_y) + E_a(y - x\cos\theta_y))}. \tag{9}$$

The square of the matrix elements appearing in Eq. (4) for dibaryon decay can be expressed as

$$|\mathcal{M}|^2 = e^2 g^2 \mathrm{Tr}[\sigma^{\mu\nu} q_\nu (\not{p}_a + m) \sigma^{\mu\rho} q_\rho (\not{p}_b - m)]/m^4, \tag{10}$$

where the dibaryon $pp\gamma$ vertex Eq. (1) is used. Evaluation of this trace gives

$$|\mathcal{M}|^2 = 16\, e^2 g^2 (p_a \cdot q)(p_b \cdot q)/m^4. \tag{11}$$

In terms of the notation just introduced, $(p_a \cdot q)/\omega = E_a - x + y\cos\theta_y$ and $(p_b \cdot q)/\omega = E_b - x - y\cos\theta_y$. The decay width can now be expressed as

$$\begin{aligned} W_D &= \frac{e^2 g^2}{4 M_D m^4} \int 16 \times 4x^2 \frac{(p_a \cdot q)(p_b \cdot q)}{\omega^2} J dx^2 d\cos\theta_y \\ &= e^2 g^2 \frac{8}{M_D m^4} \int \frac{(p_a \cdot q)(p_b \cdot q)}{\omega^2} J dx^4 d\cos\theta_y, \end{aligned} \tag{12}$$

where J is given by Eq. (9). The decay width is independent of the parity of the dibaryon state.

A detailed analysis of Eq. (12) in the limit of $\omega \ll m$ shows that the width can be approximated by

$$W_D = \frac{\alpha g^2 m}{2\sqrt{2}(4\pi)^2} \left(\frac{M_D - 2m}{m}\right)^{4.5}. \tag{13}$$

As shown in Fig. 3 this indeed gives a good approximation to the exact expression over the full energy range of interest.

Closely related to the decay width is the formation cross section of the dibaryon resonance, given by

$$d\sigma = \frac{1}{64\pi^2 s} \frac{1}{4} \frac{p_f}{p_i} |\mathcal{M}|^2 d\Omega, \tag{14}$$

where $p_i = p(p_f = \omega)$ are the initial (final) momenta in the two-body center-of-mass system and the factor $1/4$ comes from averaging over initial spin states. Integrating Eq. (14) gives for the total formation cross section

$$\begin{aligned} \sigma_f &= \frac{1}{64\pi^2 s} \frac{\omega}{p} \frac{16 e^2 g^2}{4m^4} \int (p_a \cdot q)(p_b \cdot q) 2\pi\, d\cos\theta \\ &= \frac{\alpha g^2 \omega^3}{2p\, s\, m^4} (2E^2 - 2\vec{p}^{\,2}/3), \end{aligned} \tag{15}$$

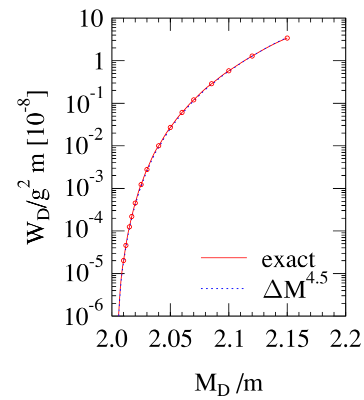

FIG. 3. (Color online) Width of a dibaryon as a function of its mass calculated from Eq. (12) compared to the approximation of Eq. (13).

where the square matrix element from Eq. (11) has been used and $E = \sqrt{m^2 + \vec{p}^{\,2}}$.

III. RESULTS

Following the suggestions of the Dubna group [7] we have taken in the numerical examples the dibaryon at a mass of $M_D = 1956$ MeV. As quantum numbers we have considered $(J^P, T) = (0^{\pm}, 2)$. The coupling constant has been taken equal to $g = 0.2$ for both of the assumed parities for the dibaryon. According to Eq. (15) this corresponds to a formation cross section (at a beam energy of 216 MeV for which the calculations are presented) of $\sigma_f = 2$ pb, well within the bounds given by [9]. The width of the resonance calculated from Eq. (13) is $W_D = 4$ MeV, an extremely narrow width. A small formation cross section, however, also implies a small decay width.

For a narrow dibaryon state, as is the case for the present investigation, we expect to see two sharp peaks in the cross section as a function of the energy of one of the two emitted photons. One peak corresponds to the formation of the dibaryon ($p + p \to \gamma^F + D$), where the resonant photon energy in the formation process (E_γ^F) can be calculated from

$$E_\gamma^F = \frac{\left(W^2 - M_D^2\right)}{2W}, \tag{16}$$

where W^2 is the square of the center-of-mass energy of the colliding nucleons. For a resonance mass of $M_D = 1956$ MeV and an incoming beam energy of $E_p = 216$ MeV the resonance energy is $E_\gamma^F = 25.4$ MeV in the center-of-mass system. Another peak in the spectrum will correspond to the decay of the dibaryon into two protons and a photon ($D \to p + p + \gamma^D$). The energy of the decay photon (E_γ^D) will depend strongly on the relative energy of the pp system in the final state.

Figures 4 and 5 show the fully exclusive cross section $\frac{d^8\sigma}{d\Omega_1 d\Omega_2 d\theta_{\gamma\gamma} dM_{\gamma\gamma} d\phi_1 d\varepsilon_1}$ as a function of photon energy. The peaks at energies $E_\gamma^F = 25.4$ MeV and E_γ^D, resulting from the dibaryon mechanism, stand out from the "background" due

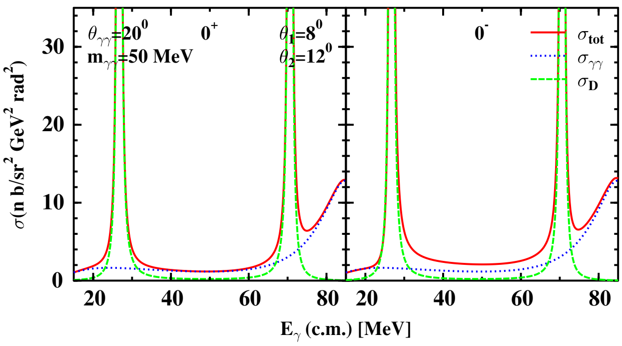

FIG. 4. (Color online) The fully exclusive cross section $\frac{d^8\sigma}{d\Omega_1 d\Omega_2 d\theta_{\gamma\gamma} dM_{\gamma\gamma} d\phi_1 d\varepsilon_1}$ as a function of photon energy for two different coplanar kinematics. The bremsstrahlung contribution is given by the dotted curve, the dibaryon contribution is given by the dashed curve, and the total cross section is given by the solid curve. The left and right panels show the results for $J^\pi = 0^-$ and $J^\pi = 0^+$ dibaryons, respectively.

to two-photon bremsstrahlung ($\sigma_{\gamma\gamma}$). This peak to background ratio depends strongly on the dibaryon coupling constant. The energy of the second peak (E_γ^D, moving from 70 to 55 MeV) clearly depends on the relative energy of the two outgoing protons, which, in turn, depends strongly on their relative angle. It is worth noting that these results closely parallel that obtained by the Dubna group [7]. The rise of the cross section at higher energies can be understood from the fact that the energy of the other photon (which is not plotted) becomes small and one approaches the soft-photon pole in the cross section.

It should be noted that the cross section due to the dibaryon mechanism (σ_D) is independent of the assumed parity of the dibaryon. The fact that the total cross section (σ_{tot}) does depend on the parity of the dibaryon shows that there is strong interference between the dibaryon and bremsstrahlung amplitudes. To express this more clearly we have plotted in Fig. 6 and 7 cross-section ratios

$$P_k = \frac{\sigma_k}{\sigma_{\text{tot}}}, \tag{17}$$

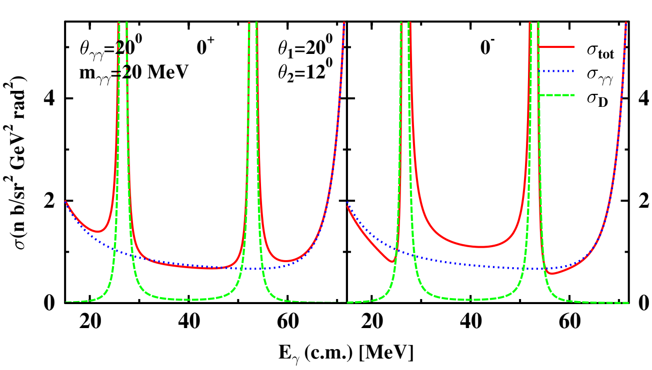

FIG. 5. (Color online) Same as Fig. 4 for different kinematics.

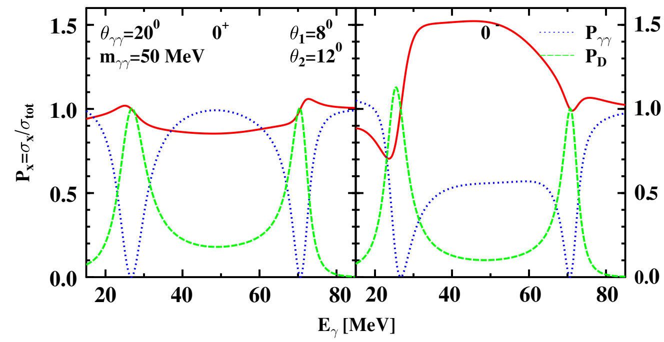

FIG. 6. (Color online) The cross-section ratios, defined in Eqs. (17) and (18), plotted as a function of photon energy. The dotted curve shows the resonance contribution, the dashed curve shows the bremsstrahlung contribution, and the solid curve shows the quantity P.

where the index $k = D$ or $\gamma\gamma$. Further, we also define

$$P = \frac{\sigma_{\text{tot}}}{\sigma_D + \sigma_{\gamma\gamma}}. \tag{18}$$

This last quantity shows most clearly the effects of interference.

Without interference between the bremsstrahlung and the dibaryon contributions, P [see Eq. (18)] should be equal to unity. However, because of the interference its value can vary between the extremes of zero and two. The effect of the interference is more pronounced for the case of a 0^- dibaryon than for 0^+ (see Figs. 6 and 7), demonstrating that the interference can serve as a tool to determine the parity of the dibaryon resonance.

IV. CONCLUSION

We have investigated the effects of a possible dibaryon state in two-photon emission in *pp* scattering. For a coupling strength where the formation cross section is well within the experimental bounds, the width of the resonance is very narrow and stands out clearly over the continuum two-photon bremsstrahlung spectrum. In high-statistics experiments it thus should be clearly visible. The "background" due to two-photon

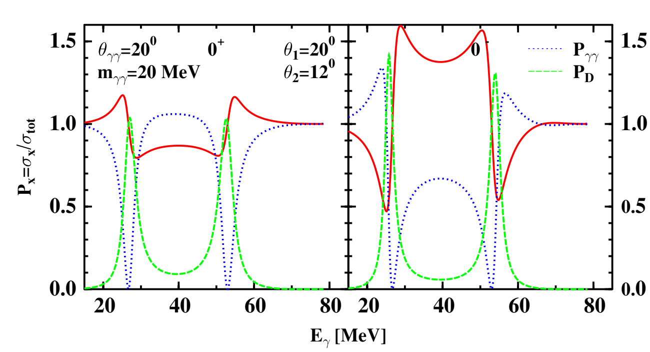

FIG. 7. (Color online) Same as Fig. 6, however for kinematics corresponding to larger proton angles and a smaller two-photon invariant mass.

bremsstrahlung, however, cannot be ignored and should be taken into account in a simulation of the experimental results. In an exclusive measurement the interference with the bremsstrahlung continuum allows for a determination of the quantum numbers of the dibaryon. In integrated spectra, especially when folded with experimental resolution of several MeV, the peak may reduce to a fractional enhancement of the cross section over that due to two-photon bremsstrahlung [12].

ACKNOWLEDGMENTS

We thank J. Bacelar and A. Yu. Korchin for discussions and a critical reading of the manuscript. Part of this work was performed as part of the research program of the Stichting voor Fundamenteel Onderzoek der Materie (FOM) with financial support from the Nederlandse Organisatie voor Wetenschappelijk Onderzoek (NWO). One of the authors (S.T.) acknowledges support from KVI and the 21st COE program (Towards a New Basic Science; Depth and Synthesis) of Osaka University.

APPENDIX A: KINEMATICS FOR TWO-PHOTON PRODUCTION

For the reaction $N + N \to N + N + \gamma + \gamma$ the momenta are denoted by p_1, p_2, p_1', p_2', q_1, and q_2 (see Fig. 1). Energy-momentum conservation reads $p_1 + p_2 = p_1' + p_2' + q_1 + q_2$. The cross section is

$$d\sigma = \frac{m_p^4}{j}\int |A|^2 (2\pi)^4 \delta^4(p_1 + p_2 - p_1' - p_2' - q_1 - q_2) \times \frac{d^3p_1'}{(2\pi)^3 E_1'}\frac{d^3p_2'}{(2\pi)^3 E_2'}\frac{d^3q_1}{(2\pi)^3 2\varepsilon_1}\frac{d^3q_2}{(2\pi)^3 2\varepsilon_2},$$

where $A = \mathcal{M}^{\mu\nu}\epsilon_{1\mu}^*\epsilon_{2\nu}^*$ is the invariant amplitude, ϵ_1 and ϵ_2 are the polarization vectors of the photons, $\varepsilon_1 = |\vec{q}_1|$, $\varepsilon_2 = |\vec{q}_2|$, and $j = \sqrt{(p_1 \cdot p_2)^2 - m_p^4} = m_p|\vec{p}_{\rm Lab}|$ in the laboratory frame where $p_2 = (m_p, \vec{0})$. Using the identity $\int \delta^4(q_1 + q_2 - k) d^4k = 1$, we can put the cross section in the form

$$d\sigma = \frac{m_p^4}{(2\pi)^8 j}\int |A|^2 \delta^4(p_1 + p_2 - p_1' - p_2' - k) I_{\gamma\gamma} \times \frac{d^3p_1'}{E_1'}\frac{d^3p_2'}{E_2'} d^4k, \tag{A1}$$

where $I_{\gamma\gamma}$ is the two-photon phase-space integral defined as

$$I_{\gamma\gamma} = \int \delta^4(q_1 + q_2 - k)\frac{d^3q_1}{2\varepsilon_1}\frac{d^3q_2}{2\varepsilon_2}.$$

To calculate this integral we deviate from the steps followed in Ref. [10]. To emphasize the dependence on photon energy the δ function is used to eliminate the polar-angle dependence,

$$\begin{aligned} I_{\gamma\gamma} &= \int \delta(\varepsilon_1 + \varepsilon_2 - k_0)\frac{\varepsilon_1}{4\varepsilon_2} d\cos\theta_1\, d\phi_1\, d\varepsilon_1 \\ &= \frac{d\phi_1\, d\varepsilon_1}{4|\vec{k}|}, \end{aligned} \tag{A2}$$

with $\varepsilon_2 = \sqrt{|\vec{k}|^2 - 2|\vec{k}|\varepsilon_1\cos\theta_1 + \varepsilon_1^2}$, where we introduced the polar and azimuthal angles θ_1 and ϕ_1 between the 3-vectors $\vec{k}$ and $\vec{q}_1$.

As a last step the integration over k_0 in Eq. (A1) is replaced by an integration over the two-photon invariant mass using $k_0 dk_0 = m_{\gamma\gamma} dm_{\gamma\gamma}$. We obtain

$$d\sigma = \frac{2m_p^4}{(2\pi)^8 j}\int |A|^2 J(m_{\gamma\gamma}) I_{\gamma\gamma} m_{\gamma\gamma} dm_{\gamma\gamma}, \tag{A3}$$

where we introduced the three-particle phase-space integral

$$J(m_{\gamma\gamma}) = \int \delta^4(p_1 + p_2 - p_1' - p_2' - k) \times \frac{d^3p_1{}'}{E_1'}\frac{d^3p_2{}'}{E_2'}\frac{d^3k}{2k_0}. \tag{A4}$$

In the one-photon bremsstrahlung the similar integral is traditionally evaluated in polar coordinates (see, for example, [13]), leading to the cross section of the type $\frac{d^8\sigma}{d\Omega_1 d\Omega_2 d\theta_{\gamma\gamma} dM_{\gamma\gamma} d\phi_1 d\varepsilon_1}$ as shown in Figs. 4 and 5.

[1] G. Wagner, L. Y. Glozman, A. J. Buchmann, and A. Faessler, Nucl. Phys. **A594**, 263 (1995), nucl-th/9502028.
[2] Q. B. Li, P. N. Shen, Z. Y. Zhang, and Y. W. Yu, Nucl. Phys. **A683**, 487 (2001), nucl-th/0009038.
[3] P. J. Mulders, A. T. M. Aerts, and J. J. De Swart, Phys. Rev. D **21**, 2653 (1980).
[4] R. Mota, Phys. Rev. C **65**, 034006 (2002).
[5] S. B. Gerasimov and A. S. Khrykin, Mod. Phys. Lett. A **8**, 2457 (1993).
[6] S. B. Gerasimov, S. N. Ershov, and A. S. Khrykin, Phys. At. Nucl. **58**, 844 (1995).
[7] A. S. Khrykin *et al.*, Phys. Rev. C **64**, 034002 (2001), nucl-ex/0012011.
[8] R. Bilger, Nucl. Phys. **A629**, 141c (1998).
[9] H. Calen *et al.*, Phys. Lett. **B427**, 248 (1998).
[10] O. Scholten and A. Y. Korchin, Phys. Rev. C **65**, 054004 (2002), nucl-th/0203078.
[11] H. Huisman *et al.*, Phys. Rev. C **65**, 031001 (2002).
[12] I. Gasparic (private communication).
[13] A. Korchin and O. Scholten, Nucl. Phys. **A602**, 423 (1996).