Available online at www.sciencedirect.com

ScienceDirect

Physics Procedia 52 (2014) 46 – 51

Physics Procedia

Physics and Applications of High Brightness Beams Workshop, HBEB 2013

Inverse Compton cross section revisited

C. Curatolo[a,b,*], L. Lanza[a], V. Petrillo[a,b]

[a] University of Milan, via Celoria 16, 20133 Milan, Italy
[b] INFN - Milan, via Celoria 16, 20133 Milan, Italy

Abstract

The design of advanced machines working in the quantum regime (ELI-NP, IRIDE, $e-\gamma$ and $\gamma-\gamma$ colliders) requires to set the fundamentals needed to have an accurate prediction of the radiation qualities after the Compton scattering. Due to the high energy of the electron beam in the cases above mentioned, the quantum effects, referred as inverse Compton, which occur during the collision with the laser radiation, are not negligible.
We present a rigorous method to obtain the inverse Compton cross section in the general case of not null initial momentum of the electrons from a pure QED calculation, avoiding the usual approaches based on the derivation of this cross section either from the Klein and Nishina formula and the Lorentz transformations or throught Feynman diagrams and Mandelstam invariants. In the derivation of the cross section from the transition amplitude we pay particular attention to the long time behavior of the system evolution. Proceeding in this way we obtain the transition probability in the time unit, which integrated over the solid angle of emission defines spectrum and number of the scattered photons.

Peer-review under responsibility of the scientific committee of HBEB 2013

Keywords: Inverse Compton scattering; QED; differential cross section

1. Introduction

In the approach to Compton effect, which consists in the calculation of the cross section of the process $e^- + photon \rightarrow e^- + photon$, the electron at the initial stage is considered at rest, $\vec{q_i} = 0$, as it happens in case of a X or γ radiation heading a target. The target's particles have kinetic energy negligible compared to the one of the incident photons. Under this hypotesis, considering in the electromagnetic interaction between electron and photon the transition amplitude up to the second order, the standard Klein and Nishina formula [1, 2] is obtained. Since the theory is relativistic, the result expressed by the Klein and Nishina formula is supposed to be generalized to the case $\vec{q_i} \neq 0$ with an appropriate Lorentz transformation. This transformation acts on the photon's 4-momentum k and polarization vectors $e_\lambda(k)$ with $\lambda = 0, 1, 2, 3$, which are related to the plane waves of the electromagnetic field potential $A^\mu(x, t)$. In each inertial frame, only the vectors whose index is $\lambda = 1, 2$ correspond to physical photons. This choice is not covariant: for

* Corresponding author.
E-mail address: camilla.curatolo@unimi.it

Peer-review under responsibility of the scientific committee of HBEB 2013
doi:10.1016/j.phpro.2014.06.008

this reason the Lorentz transformation has to be associated to a gauge transformation or it is necessary to consider the transformation law of the electromagnetic field's tensor and to write the electric field through the vectors $e_\Lambda(k)$. Since in quantum mechanics the interaction is expressed utilizing the electromagnetic potential, the procedure [3, 4, 5, 6] described above could be not completly justified. On the other hand the use of Feynman diagrams and Mandelstam invariants as in [7, 8] leads to some controversial mathatical procedure such as Dirac δ functions squared and forces the actual physical setting into the formalism.
In the following is presented the alternative method we developed and the results obtained.

2. Preliminaries

We recall here some preliminaries and notations usefull in the following calculation, for a more detailed treatise see [9].
Concerning the electrons we introduce the field operator

$$\hat{\Psi}_\sigma(x) = \sqrt{m_e c}\sum_{r=1}^{2}\int \frac{d^3p}{p^0}\left(\frac{e^{-\frac{i}{\hbar}p\cdot x}}{(2\pi\hbar)^{\frac{3}{2}}}v^r_\sigma(\vec{p})\hat{b}^r(\vec{p}) + \frac{e^{\frac{i}{\hbar}p\cdot x}}{(2\pi\hbar)^{\frac{3}{2}}}w^r_\sigma(\vec{p})\hat{d}^{r\dagger}(\vec{p})\right) \tag{1}$$

where $\hat{b}^r(\vec{p})$ and $\hat{d}^{r\dagger}(\vec{p})$ are respectively the annihilation operator for particles and the creation operator for the antiparticles [1], satisfying the Dirac equations

$$\left(i\sum_\mu \gamma^\mu \frac{\partial}{\partial x^\mu} - \frac{m_e c}{\hbar}\right)\hat{\Psi}_\sigma(x) = 0 \Leftrightarrow p\cdot p = m_e^2 c^2 \tag{2}$$

$$\frac{1}{2}(\gamma^\nu\gamma^\mu + \gamma^\mu\gamma^\nu) = g^{\mu\nu}I \qquad \text{where} \qquad \gamma^0 = \begin{pmatrix} I & 0 \\ 0 & -I \end{pmatrix} \qquad \gamma^i = \begin{pmatrix} 0 & \sigma^i \\ -\sigma^i & 0 \end{pmatrix}$$

and the adjoint field operator $\hat{\bar{\Psi}}_\sigma(x) = \hat{\Psi}^\dagger(x)\gamma^0$. The operator acting on the photons is the potential vector

$$\hat{A}^\mu(x) = \hbar\frac{\sqrt{c}}{\sqrt{2}}\int \frac{d^3p}{p^0}\left(\frac{e^{-\frac{i}{\hbar}p\cdot x}}{(2\pi\hbar)^{\frac{3}{2}}}\sum_{\lambda=0}^{3} e^\mu_\lambda(\vec{p})g^{\lambda\nu}\hat{c}_\nu(\vec{p}) + \frac{e^{\frac{i}{\hbar}p\cdot x}}{(2\pi\hbar)^{\frac{3}{2}}}\sum_{\lambda=0}^{3} e^\mu_\lambda(\vec{p})g^{\lambda\nu}\hat{c}^\dagger_\nu(\vec{p})\right) \tag{3}$$

with $\hat{c}_\nu(\vec{p})$ and $\hat{c}^\dagger_\nu(\vec{p})$ the annihilation and the creation operator respectively [2] and $g^{\mu\nu}\frac{\partial^2}{\partial x^\mu \partial x^\nu}\hat{A}^\mu(x) = 0$. The physical states are characterized by photons having $\lambda = 1, 2$, the one with transverse polarization, i.e. the states $|f\rangle$ and $|g\rangle$ solutions of

$$\langle g|\frac{\partial \hat{A}^\mu(x)}{\partial x^\mu}|f\rangle = 0$$

Quantum Electrodynamics describes the time evolution of the system constituted by an electron and a photon colliding at a certain instant non exactly localized in a macroscopic time interval $[t_1, t_2]$:

$$|\Psi^{inter}_{t_2}\rangle = \hat{U}^{inter}(t_2, t_1)|\Psi^{inter}_{t_1}\rangle \qquad \text{where} \qquad \hat{U}^{inter}(t_2, t_1) = e^{\frac{i}{\hbar}\hat{H}_0 t_2}e^{-\frac{i}{\hbar}\hat{H}(t_2-t_1)}e^{-\frac{i}{\hbar}\hat{H}_0 t_1} \tag{4}$$

---

[1] For these operators anticommutators are null for all cases except

$$\left[\hat{b}^r(\vec{p}), \hat{b}^{r'\dagger}(\vec{p}')\right]_+ = p^0\delta_{rr'}\delta_3(\vec{p}-\vec{p}') \qquad \left[\hat{d}^r(\vec{p}), \hat{d}^{r'\dagger}(\vec{p}')\right]_+ = p^0\delta_{rr'}\delta_3(\vec{p}-\vec{p}')$$

[2] For these operators commutators are null for all cases except

$$\left[\hat{c}_0(\vec{p}), \hat{c}^\dagger_0(\vec{p}')\right]_- = -p^0\delta_3(\vec{p}-\vec{p}') \qquad \left[\hat{c}_l(\vec{p}), \hat{c}^\dagger_{l'}(\vec{p}')\right]_- = p^0\delta_{ll'}\delta_3(\vec{p}-\vec{p}') \qquad \text{for l=1,2,3}$$

The Hamiltonian is

$$\hat{H} = \hat{H}_0 + \hat{H}_1^{inter} \qquad \text{with} \qquad \hat{H}_1^{inter} = e\int d^3x : \hat{\overline{\Psi}}(x)\gamma^\mu\hat{\Psi}(x)\hat{A}_\mu(x) : \tag{5}$$

and, being T the chronological ordering operator, the interaction operator has the form

$$\hat{U}^{inter}(t,t_1) = T\left(e^{-\frac{i}{\hbar}\int_{t_1}^{t} dt' \int_{\mathbb{R}^3} d^3x' e:\hat{\overline{\Psi}}(x')\gamma^\mu\hat{\Psi}(x')\hat{A}_\mu(x'):}\right) \tag{6}$$

3. Calculation

The electron-photon system is described by an appropriate mixture of the initial states

$$|\Psi_{t_1}\rangle = \iint \frac{d^3q'}{q'^0}\frac{d^3k'}{k'^0} e^{\frac{i}{\hbar}(q'^0+k'^0)ct_1} \sum_{r'\lambda'=1}^{2} \Psi_{1e}(q',r')\Psi_{1\gamma}(k',\lambda')\hat{b}^{r'\dagger}(q')\hat{c}^{\lambda'\dagger}(k')|0\rangle \tag{7}$$

with the wave functions peaked respectively at ($\vec{q_i}$ and $\vec{k_i}$) and displaced inside the beams. We are interested in the transition from $(\vec{q_i},\vec{k_i})$ at t_1 to $(\vec{q_f},\vec{k_f})$ at the final time t_2 such that $\Psi_{1e}(q_f,r)=0$, $\Psi_{1\gamma}(k_f,\lambda)=0$, i.e. not belonging to the initial state. The statistics of these events is obtained by using

$$\hat{E}^{r\lambda}(M) = \iint_M \frac{d^3q}{q^0}\frac{d^3k}{k^0}\hat{b}^{r\dagger}(\vec{q})\hat{c}^\dagger_\lambda(\vec{k})|0\rangle\langle 0|\hat{c}_\lambda(\vec{k})\hat{b}^r(\vec{q}) \qquad \text{normalized in the sense that} \qquad \sum_{r\lambda}\hat{E}^{r\lambda}(\mathbb{R}^6) = \hat{P}_{e,\gamma} \tag{8}$$

with $\hat{P}_{e,\gamma}$ projector on the Fock space of one electron and one photon.
In order to calculate the probability $\|\hat{E}^{r\lambda}(M)\hat{U}^{inter}(t_2,t_1)\Psi_{t_1}\|^2$ we consider

$$\Pi^{r\lambda}(\vec{q},\vec{k};t_2,t_1) = |\langle\Phi_{t_2}|T\left(e^{-\frac{ie}{\hbar c}\int_{ct_1}^{ct_2}dx^0\int_{\mathbb{R}^3}d^3x:\hat{\overline{\Psi}}(x)\gamma^\mu\hat{\Psi}(x)\hat{A}_\mu(x):}\right)|\Psi_{t_1}\rangle|^2\frac{1}{k^0q^0} \tag{9}$$

choosing

$$|\Phi_{t_2}\rangle = \hat{b}^{r_f\dagger}(\vec{q_f})\hat{c}^\dagger_{\lambda_f}(\vec{k_f})|0\rangle$$

From a second order perturbative approach expansion of the time evolution operator we obtain

$$\begin{aligned}\langle\Phi_{t_2}| - i\left(\frac{e}{\hbar c}\right)^2 \iint_{ct_1}^{ct_2} dx_1^0 dx_2^0 \iint_{\mathbb{R}^3} d^3x_1 d^3x_2 \hat{\overline{\Psi}}^{(-)}(x_2)\gamma^{\mu_2}\Big(\hat{A}^{(-)}_{\mu_2}(x_2)S_F(x_2-x_1)\hat{A}^{(+)}_{\mu_1}(x_1)\\ +\hat{A}^{(-)}_{\mu_1}(x_1)S_F(x_2-x_1)\hat{A}^{(+)}_{\mu_2}(x_2)\Big)\gamma^{\mu_1}\hat{\Psi}^{(+)}(x_1)|\Psi_{t_1}\rangle\end{aligned} \tag{10}$$

where (+) e (−) are respectively the annihilation and creation part of the operators and S_F describes the normal order.
To proceed with the calculation of this expression, the most important steps are the following:

- To put in evidence the energy conservation we consider the long time behavior of the system evolution.
- Since

$$S_F(x) = \lim_{\epsilon\to 0^+}\int dk^0 d^3k \frac{e^{-ikx}}{(2\pi)^4}\frac{k\gamma + \frac{m_e c}{\hbar}I}{k^2 - \frac{m_e^2c^2}{\hbar^2} + i\epsilon} \tag{11}$$

  we integrate on k and develop the calculation of the integral on k^0 by using the complex analysis methods.

- We take into accont that the initial wave function are peaked around initial momentum $\vec{q_i}$ for the electron and $\vec{k_i}$ for the photon, so that the initial momentum of the particles is in good approximation $\vec{q_i}$ and $\vec{k_i}$.
- The wave function of a single particle describes the momentum of the particle through its modulus and is related to the position through the phase. The position of the photon is hardly determined, so it is necessary to take an average over the infinite possible choices of $\vec{x}_0$ in a macroscopic space region ω with volume V, inside the bunch, symmetric around the origin, where the density of the photons is constant.

4. Differential cross section

The calculation provides the transition probability, defined as the probability that a photon characterized by initial polarization and momentum $\lambda_i, \vec{k_i}$, after the scattering has polarization and momentum $\lambda_f, \vec{k_f}$:

$$P(\lambda_i,\lambda_f,\vec{k_f}\in d\Omega)=\frac{d\Omega(t_2-t_1)e^4m_e^2c}{V(4\pi\hbar)^2}\frac{1}{|\vec{k_i}|}\frac{1}{\sqrt{m_e^2c^2+|\vec{q_i}|^2}}\frac{1}{\sqrt{m_e^2c^2+|\vec{q_i}+\vec{k_i}-\vec{k_f}|^2}}\frac{|\vec{k_f}|\left(\sqrt{m_e^2c^2+|\vec{q_i}|^2}+|\vec{k_i}|-|\vec{k_f}|\right)}{\sqrt{m_e^2c^2+|\vec{q_i}|^2}+k_i-|\vec{q_i}+\vec{k_i}|\cos\theta}$$
$$\cdot Tr_{\mathbb{C}^4}\left\{\left[\left(\not{e}_{\lambda_f}(\vec{k_f})\not{k}_i\not{e}_{\lambda_i}(\vec{k_i})+\not{e}_{\lambda_f}(\vec{k_f})q_i\cdot e_{\lambda_i}(\vec{k_i})\right)\frac{\hbar}{2q_i\cdot k_i}+\left(\not{e}_{\lambda_i}(\vec{k_i})(-\not{k}_f)\not{e}_{\lambda_f}(\vec{k_f})+\not{e}_{\lambda_i}(\vec{k_i})q_i\cdot e_{\lambda_f}(\vec{k_f})\right)\frac{\hbar}{-2q_i\cdot k_f}\right]\left(\frac{m_ecI+\not{q}_i}{4m_ec}\right)\right.$$
$$\left.\cdot\left[\left(\not{e}_{\lambda_i}(\vec{k_i})\not{k}_i\not{e}_{\lambda_f}(\vec{k_f})+\not{e}_{\lambda_f}(\vec{k_f})q_i\cdot e_{\lambda_i}(\vec{k_i})\right)\frac{\hbar}{q_i\cdot k_i}-\left(\not{e}_{\lambda_f}(\vec{k_f})(-\not{k}_f)\not{e}_{\lambda_i}(\vec{k_i})+\not{e}_{\lambda_i}(\vec{k_i})q_i\cdot e_{\lambda_f}(\vec{k_f})\right)\frac{\hbar}{q_i\cdot k_f}\right]\frac{m_ecI+\not{q}_i+\not{k}_i-\not{k}_f}{4m_ec}\right\} \tag{12}$$

where $\not{a}=a\cdot\gamma$, $d\Omega=d\Omega(\theta,\phi)$ is the solid angle and θ is the angle between the directions of the incoming electron and the emitted photon.
From the transition probability we can extract the geometrical quantity called differential cross section:

$$\sigma_{\lambda_i,\lambda_f}(\theta,\phi)=\frac{P(\lambda_f,\lambda_i,\vec{k_f}\in d\Omega(\theta,\phi))V}{d\Omega(t_2-t_1)} \tag{13}$$

In the usual Compton scattering the electron at the initial stete is considered at rest, $\vec{q_i}=0$. In this case, we get the Klein and Nishina formula:

$$(\sigma_{\lambda_i,\lambda_f}(\theta,\phi))_{\vec{q_i}=0}=\frac{1}{4}\left(\frac{e^2}{4\pi m_ec^2}\right)^2\left(\frac{m_ec}{m_ec+|\vec{k_i}|(1-\cos\theta)}\right)^2\left[4(\vec{e}_{\lambda_f}\cdot\vec{e}_{\lambda_i})^2+\frac{|\vec{k_i}|^2(1-\cos\theta)^2}{m_ec(m_ec+|\vec{k_i}|(1-\cos\theta))}\right] \tag{14}$$

Since $|\vec{k_i}|\ll m_ec$, for not polarized photon beam and not observed polarization of the scattered photons

$$(\sigma_{\lambda_i,\lambda_f}(\theta,\phi))_{\vec{q_i}=0}=\left(\frac{e^2}{4\pi m_ec^2}\right)^2\frac{1}{2}(1+\cos^2\theta)=r_0^2\frac{(1+\cos^2\theta)}{2} \tag{15}$$

Considering the inverse Compton scattering we have to analize the general case of $\vec{q_i}\neq 0$ and the result is:

$$\sigma_{\lambda_i,\lambda_f}(\theta,\phi)=A\cdot B \tag{16}$$

where

$$A=\left(\frac{e^2}{4\pi m_ec^2}\right)^2\frac{m_ec}{\sqrt{m_e^2c^2+|\vec{q_i}|^2}}\frac{m_ec}{\sqrt{m_e^2c^2+|\vec{q_i}+\vec{k_i}-\vec{k_f}(\theta,\phi)|^2}}\frac{|\vec{k_f}(\theta)|\left(\sqrt{m_e^2c^2+|\vec{q_i}|^2}+|\vec{k_i}|-|\vec{k_f}(\theta)|\right)}{\sqrt{m_e^2c^2+|\vec{q_i}|^2}+|\vec{k_i}|-|\vec{q_i}+\vec{k_i}|\cos\theta} \tag{17}$$

and

$$
\begin{aligned}
B = \frac{1}{2}\Bigg\{ & k_i \cdot k_f \left[\frac{\vec{q}_i \cdot \vec{e}_{\lambda_f} \vec{e}_{\lambda_f} \cdot \vec{k}_i - \vec{q}_i \cdot \vec{e}_{\lambda_i} \vec{e}_{\lambda_i} \cdot \vec{k}_f}{q_i \cdot k_i q_i \cdot k_f} + \frac{1}{2}\left(\frac{1}{q_i \cdot k_f} - \frac{1}{q_i \cdot k_i}\right)\left(1 + \frac{(\vec{q}_i \cdot \vec{e}_{\lambda_i})^2}{q_i \cdot k_i} - \frac{(\vec{q}_i \cdot \vec{e}_{\lambda_f})^2}{q_i \cdot k_f}\right)\right] \\
& + \frac{1}{2}\frac{(\vec{q}_i \cdot \vec{e}_{\lambda_i})^2 (\vec{q}_i \cdot \vec{e}_{\lambda_f})^2 (q_i \cdot k_i - q_i \cdot k_f)^2}{(q_i \cdot k_i)^2 (q_i \cdot k_f)^2} + 2(\vec{e}_{\lambda_i} \cdot \vec{e}_{\lambda_f})^2 + \frac{(\vec{q}_i \cdot \vec{k}_f \vec{q}_i \cdot \vec{e}_{\lambda_i} \vec{e}_{\lambda_f} \cdot \vec{k}_i + \vec{q}_i \cdot \vec{k}_i \vec{q}_i \cdot \vec{e}_{\lambda_f} \vec{e}_{\lambda_i} \cdot \vec{k}_f)^2}{(q_i \cdot k_i)^2 (q_i \cdot k_f)^2} \\
& + \frac{1}{2}\left(\vec{q}_i \cdot \vec{e}_{\lambda_f} \vec{e}_{\lambda_f} \cdot \vec{k}_i + \vec{q}_i \cdot \vec{e}_{\lambda_i} \vec{e}_{\lambda_i} \cdot \vec{k}_f + 5 \vec{q}_i \cdot \vec{e}_{\lambda_i} \vec{q}_i \cdot \vec{e}_{\lambda_f} \vec{e}_{\lambda_i} \cdot \vec{e}_{\lambda_f}\right)\left(\frac{1}{q_i \cdot k_i} - \frac{1}{q_i \cdot k_f}\right) \\
& + \frac{3}{2}\left[(\vec{q}_i \cdot \vec{e}_{\lambda_i})^2 \vec{q}_i \cdot \vec{e}_{\lambda_f} \vec{e}_{\lambda_f} \cdot \vec{k}_i \left(\frac{1}{(q_i \cdot k_i)^2} - \frac{1}{q_i \cdot k_i q_i \cdot k_f}\right) + \vec{q}_i \cdot \vec{e}_{\lambda_i} (\vec{q}_i \cdot \vec{e}_{\lambda_f})^2 \vec{e}_{\lambda_i} \cdot \vec{k}_f \left(\frac{1}{q_i \cdot k_i q_i \cdot k_f} - \frac{1}{(q_i \cdot k_f)^2}\right)\right] + \\
& + \vec{e}_{\lambda_i} \cdot \vec{e}_{\lambda_f} \left(\frac{\vec{q}_i \cdot \vec{e}_{\lambda_i} \vec{e}_{\lambda_f} \cdot \vec{k}_i}{q_i \cdot k_i} + \frac{\vec{q}_i \cdot \vec{e}_{\lambda_f} \vec{e}_{\lambda_i} \cdot \vec{k}_f}{q_i \cdot k_f}\right) + \frac{1}{4}\left(\frac{(\vec{q}_i \cdot \vec{e}_{\lambda_i})^2}{(q_i \cdot k_i)^2} + \frac{(\vec{q}_i \cdot \vec{e}_{\lambda_f})^2}{(q_i \cdot k_f)^2}\right)\left(q_i \cdot k_i - q_i \cdot k_f\right)\Bigg\}
\end{aligned}
\tag{18}
$$

This formula keeps into account all the possible polarization of the photons.

5. Results

We consider now the situation in which the electron and the photon are perfectly counterpropagating with $|\vec{q}_i| \gg m_e c \gg |\vec{k}_i|$ and $m_e^2 c^2 \gg |\vec{q}_i||\vec{k}_i|$. Since

$$
|\vec{k}_f(\theta)| = \frac{|\vec{k}_i| \sqrt{m_e^2 c^2 + |\vec{q}_i|^2} + |\vec{q}_i||\vec{k}_i|}{|\vec{k}_i| + \sqrt{m_e^2 c^2 + |\vec{q}_i|^2} - |\vec{q}_i + \vec{k}_i| \cos\theta}
\tag{19}
$$

from eq. (17) we obtain

$$
A \approx \left(\frac{e^2}{4\pi m_e c^2}\right)^2 \frac{8|\vec{q}_i|^2}{m_e^2 c^2} \frac{1}{\left(1 + \frac{2|\vec{q}_i|^2}{m_e^2 c^2}(1 - \cos\theta)\right)^2} = r_0^2 \frac{8\gamma^2}{(1 + 2\gamma^2(1 - \cos\theta))^2}
\tag{20}
$$

and by operating the sum and the average on the polarizations, eq. (18) becomes

$$
B \approx 1 + \cos^2\theta - \frac{\sin^2\theta\,(1 - 3\cos\theta)}{8\left(1 - \cos\theta + \frac{1}{2\gamma^2}\right)} + \frac{\sin^4\theta\left(1 - \cos\theta + \frac{1}{2\gamma^2}\right)^2}{8}
\tag{21}
$$

We notice that the dependence on the angle θ is critical for both A and B implying a strong decreasing of $\sigma_{\lambda_i,\lambda_f}(\theta, \phi)$ as shown in Fig.(1), where the cross section (16) is presented as function of the angle θ. The electron's energy increasing enhances the critical dependence of the differential cross section on the emission angle, reason why the acceptance angle is determined by the energy of the electron and is usually $\theta_{acc} = \frac{1}{\gamma}$.

6. Conclusions

In this work, we revise the derivation of the electron-photon cross section in the case of scattering between moving particles. The method we propose is based on a pure QED rigorous calculation. In the classical approximation we obtain the Thomson cross section, in the case of Compton scattering (with electrons at rest) we get the Klein and Nishina formula (14) and considering the inverse Compton scattering our result is in agreement with the ones reported in litterature [3, 4, 5, 6, 7, 8].

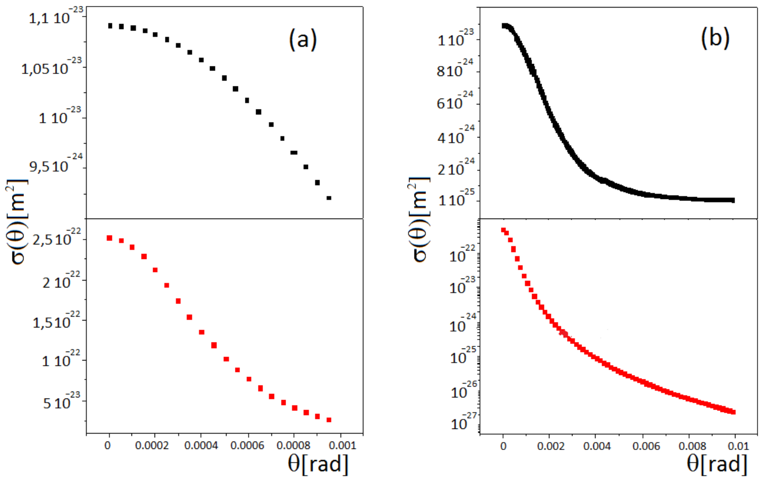

Fig. 1. (a) Differential cross section for inverse Compton scattering as function of the scattering angle $0 < \theta < 0,001$ for $|\vec{q}_i| = 150\,MeV$ (in black) and $|\vec{q}_i| = 720\,MeV$ (in red); (b) Differential cross section for inverse Compton scattering as function of the scattering angle $0 < \theta < 0,01$ for $|\vec{q}_i| = 150\,MeV$ (in black) and $|\vec{q}_i| = 720\,MeV$ (in red).

References

[1] O. Klein , Y. Nishina. *Uber di Streuung von Strahlung durch freie Elektronen nach der neuen relativistichen Quantendynamik von Dirac*. Z. Phys. 52, 11-12 (1929)

[2] A. Compton. *A Quantum Theory of the Scattering of X-rays by Light Elements*. Phys. Rev. 21, 483-502 (1923)

[3] W. Heitler. *The Quantum Theory of Radiation*. The International Series of Monographs in Physics, Oxford at the Clarendon Press (1057)

[4] J. M. Jauch, F. Roerlich. *The theory of photons and electrons*. Addison Wesley Publishing Company. Cambridge (1955)

[5] G. Bhatt, H. Grotch, E. Kazes. *Relativistic spin-dependent Compton scattering from electrons*. Phys. Rev. A 28 2195-2200 (1983)

[6] V. Petrillo et al., Nuclear Instruments and Methods in Physics Research A 693 (2012) 109.

[7] V. B. Berestetskii, E. M.Litshitz, L. P. Pitaevskii. *Quantum Electrodynamics. Landau and Lifshitz, Course of Theoretical Physics, Vol. 4* Pergamon Press (1982).

[8] I. Chaikovska, These de Doctorat *Polarized positron sources for the future linear colliders* (2012).

[9] L. Lanz. *Appunti del corso di Meccanica Quantistica* 2. Scientifica CUSL. Milano (2008)