Nuclear $\gamma$-radiation as a signature of ultra-peripheral ion collisions at the LHC

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Abstract. We study the peripheral ion collisions at LHC energies where a nucleus is excited to a discrete state and then emits $\gamma$-rays. Large nuclear Lorenz factors allow the observation of high-energy photons up to a few tens GeV and in the angular region of a few hundred micro-radians from the beam direction. These photons can be used to tagg events with particle production in the central rapidity region in ultra-peripheral collisions. To detect these photons it is necessary to have an electromagnetic detector in front of the zero-degree calorimeter in LHC experiments.

PACS. 25.75.-q Relativistic heavy-ion collisions – 25.75.Dw Particle and resonance production

Introduction

There are several reviews devoted to coherent $\gamma\gamma$ and $\gamma g$ interactions in very peripheral collisions at relativistic ion colliders [1–3]. The advantage of relativistic heavy-ion colliders is that the effective photon luminosity for two-photon physics is orders of magnitude higher than that available in $e^+e^-$ machines. There have been many suggestions to use the electromagnetic interactions of nuclei to study production of meson resonances, Higgs bosons, Radions or exotic mesons. These interactions also probe fermion, vector meson or boson pair production, as well as investigate some new physics regions (see list in ref. [3]). The $\gamma g$ interactions will open a new area of nuclear physics such as the study of nuclear gluon distribution. It is also important for the knowledge of the details of medium effects in nuclear matter at the formation of the quark-gluon plasma [4]. These effects may be studied by photo-production of heavy quarks in virtual photon-gluon interactions [4–6].

For these investigations it is necessary to select processes with large impact parameters $b$ of the colliding nuclei, $b > (R_1 + R_2)$, to exclude background from strong interactions. Note that some processes, like $\gamma\gamma$-fusion to Higgs bosons or Radions, are free from any problems caused by strong interactions of the initial state [7]. Therefore, we need an efficient trigger to distinguish $\gamma\gamma$ and $\gamma g$ interactions from others. G. Baur et al. [8] suggested to detect intact nuclei after the interaction. Evidently this is impossible in LHC experiments since nuclei fly into the beam pipe.

It is interesting to consider $\gamma$-rays emitted by the relativistic nuclei at LHC energies. This process was used for the possible explanation of the high-energy ($E_\gamma \geq 10^{12}$ eV) cosmic photon spectrum [9].

We had considered [10] the process $A + A \to A^* + A + e^+e^-, A^* \to A + \gamma'$, where a nucleus is excited by the electron (positron) $e^+ + A \to e^+ + A^*$, and suggested to detect a nuclear $\gamma$ radiation after the excitation of discrete nuclear levels [10]. These secondary photons have the energy of a few GeV and a narrow angular distribution close the beam direction due to a large Lorentz boost. The angular width is large enough for them to be detected in the electromagnetic zero-degree calorimeters (ZDC) of the future LHC experiments CMS or ALICE.

Now we calculate the production process of some system $X_f$ in $\gamma\gamma$ fusion with simultaneous excitation of the discrete nuclear level. The nucleus retains its charge $Z$ and mass $A$ in this process. So we have a clear electromagnetic interaction of nuclei at any impact parameter. The nuclear $\gamma$ radiation may be used as “event-by-event” criteria in these collisions.

In this work we consider the processes

\begin{align*}
16\text{O} + 16\text{O} &\to 16\text{O} + 16\text{O}^*(2^+, 6.92 \text{ MeV}) + X_f, \\
16\text{O}^* &\to 16\text{O} + \gamma, \\
208\text{Pb} + 208\text{Pb} &\to 208\text{Pb} + 208\text{Pb}^*(3^-, 2.62 \text{ MeV}) + X_f, \\
208\text{Pb}^* &\to 208\text{Pb} + \gamma,
\end{align*}

where the $16\text{O}$ and $208\text{Pb}$ were taken since they are the lightest and heaviest ions in the LHC program. The trigger requirements will include a signal in the central rapidity region.
region of particles from \( X_f \) decay, a signal of photons in the electromagnetic detector in front of the zero-degree calorimeter and a veto signal of neutrons in the ZDC. We suggest to use the veto signal of neutrons in order to avoid the processes with nuclear decay into nucleon fragments.

The formalism of the considered process is presented in sect. 1. The nuclear form factors are calculated in sect. 2. The angular and energy distributions of secondary photons are in sect. 3. The cross-sections of \( \eta_c \) (2.979 GeV) production are presented in sect. 4 with and without nuclear excitation. Section 5 is our conclusion.

1 Formulae of nuclear excitation cross-section and photon luminosity in peripheral interactions

Let us consider the peripheral ion collision

\[
A_1 + A_2 \rightarrow A_1^*(\lambda^P, E_0) + A_2 + X_f, \quad (1)
\]

where \( X_f \) is the produced system in \( \gamma^*\gamma^* \) fusion and \( A_1^* \) is an excited nucleus in a discrete nuclear state with spin-parity \( \lambda^P \) and energy \( E_0 \) (see fig. 1). Here the nuclei \( A_1 \) and \( A_2 \) have equal mass \( A \) and charge \( Z \), only the nucleus \( A_1 \) is excited. We suppose that the reaction product \( X_f \) decay can be detected in the central rapidity region. The nuclear \( \gamma \) radiation \( A_1^* \rightarrow A_1 + \gamma \) will be measured in the forward detectors as ZDC.

We use the quantum-mechanical plane-wave formalism [3,11] and the derivation of the equivalent photon approximation. This allows us to introduce the elastic and inelastic nuclear form factors for process (1). We take the formulae (19) and (21) in [3]:

\[
\frac{d\sigma_{A_1A_2\rightarrow A_1^*A_2X_f}}{dw_1} = \frac{d\sigma_{\gamma\gamma\rightarrow X_f}}{w_1} \frac{1}{q_1^2} \cdot \frac{1}{q_2^2}, \quad (2)
\]

where \( n_i(w_i) = \frac{d\sigma_{\gamma\gamma\rightarrow X_f}}{w_i} \) and \( W_{i,1} \) and \( W_{i,2} \) are the Lorentz scalar functions. All kinematic variables have the same definitions as in [3].

For the “elastic” photon process \( A_1A_2 \rightarrow A_1A_2X_f \) we have

\[
W_1 = 0, \quad W_2(\nu, q^2) = Z^2 F_{el}^2(-q^2) \delta(\nu + q^2/2m). \quad (4)
\]

So that [3]

\[
n(w) = \frac{Z^2 a}{\pi} \int dq_1 \frac{q_1^2}{(q_1^2)^2} F_{el}^2(-q^2), \quad (5)
\]

where \( F_{el}(q) \) is the nuclear form-factor with \( F_{el}(0) = 1 \).

For the excitation of the nucleus to a discrete state with a spin \( \lambda \) and an energy \( E_0 \) (“inelastic” photon process \( A_1A_2 \rightarrow A_1^*(\lambda^P, E_0)A_2X_f \))

\[
W_{1,2}(\nu, q^2) = \hat{W}_{1,2}(q^2) \delta(\nu - E_0),
\]

\[
\hat{W}_1 = 2\pi [(\bar{E}^c)^2 + |T^m|^2],
\]

\[
\hat{W}_2 = 2\pi \frac{q^4}{(E_0^2 - q^2)^2} \times [2M^C]^2 - E_0^2 - q^2 [(\bar{E}^c)^2 + |T^m|^2]. \quad (6)
\]

See notations again in [3].

We neglect the transverse electric \( T^c \) and transverse magnetic \( T^m \) matrix elements compared to the Coulomb one \( M^C \equiv M_\lambda \) for \( 0^+ \rightarrow \lambda^P \) nuclear transitions. Then for the “inelastic” photon process with a nuclear discrete state excitation we get

\[
n_{1,(\lambda)}(w) = \frac{4\alpha}{\pi} \int d^2q_\perp \frac{q_\perp^2}{(E_0^2 - q_\perp^2)^2} |M_{\lambda}(\nu q^2)|^2, \quad (7)
\]

where \( M_{\lambda}(q) \) is the inelastic nuclear form factor and \( -q^2 = q_\perp^2(w) + q_\perp^2 \).

The equivalent photon number (7) can be represented as function of \( q_\perp \) for inelastic photon emission:

\[
\frac{dn_{1,(\lambda)}}{dq_\perp^2}(w_1, q_\perp) = \frac{4\alpha}{\pi} \frac{q_\perp^2}{(E_0^2 - q_\perp^2)^2} |M_{\lambda}(\nu q^2)|^2 = \frac{4\alpha}{\pi} \left| \frac{q_\perp}{(E_0^2 - q_\perp^2)} M_{\lambda}(\nu q^2) e^{i\varphi_\perp} \right|^2, \quad (8)
\]

where \( q_\perp e^{i\varphi_\perp} = q_\perp \) (see [12]).

Let us do the inverse transformation to the impact parameter \( b \) presentation:

\[
f(b) = \frac{1}{\pi} \int d^2q_\perp e^{-i\nu q_\perp b} f(q_\perp). \quad (9)
\]

For the function under the module in eq. (8) we get

\[
f(b) = \frac{1}{\pi} \int d^2q_\perp \frac{q_\perp^2}{(E_0^2 - q_\perp^2)} M_{\lambda}(\nu q^2) e^{i\varphi_\perp} e^{-i\nu q_\perp b} = \int dJ_{1}(q_\perp b) = \int d\nu \frac{\nu^2}{(E_0^2 - \nu^2 - q_\perp^2) b^2} \times M_{\lambda} \left( \frac{\nu^2}{b^2} \right) J_1(\nu). \quad (10)
\]
Here $x = q_1 b = w b / \gamma_A$ and $u = q_\perp b$.

If we take $M_{el}$ instead of the inelastic $M_\lambda$ as

$$|M_{el}(-q^2)|^2 = \frac{Z^2}{4\pi} F_{el}^2(-q^2)$$  \hspace{1cm} (11)

and put $E_0 = 0$, then we get a well-known formula of the impact parameter-dependent equivalent photon number of the $A_2$ nucleus (see (4) in [12]):

$$N_{2el}^{(e)}(w, b) = \frac{Z^2 \alpha}{\pi^2} \frac{1}{b^2} \quad \int \frac{u^2}{x^2 + u^2} J_1(u) F_{el}[-(x^2 + u^2)/b^2] \, du.$$

(12)

For a point charge, $F_{el}(q) \equiv 1$, we readily obtain

$$N_{2el}^{(e)}(w, b) = \frac{Z^2 \alpha}{\pi^2} \frac{1}{b^2} \quad K_1^2(x),$$

(13)

in agreement with [3] at very large $\gamma_A$.

We write the form factors of the elastic and inelastic nuclear process in the same forms:

$$F_{el}^2(q) = \left( \frac{Z^2 \alpha}{\pi^2} \right) \frac{1}{b^2} F_{el}^2(q),$$

(14)

$$F_0^2(q) = \left( \frac{4\pi}{q^2} \right) \int \frac{\sin(qr)\rho_0(r)\,dr}{\cos^2(qr)} \bigg|_{q=0} \rightarrow 1,$$

(15)

$$F_{\lambda}^2(q) = (2\lambda + 1) \frac{4\pi}{q^2} \int f_\lambda(qr) \rho_\lambda(r, Z) r^2 \, dr \bigg|_{q=0} \rightarrow (4\pi)^2 B(E\lambda) \frac{1}{e^2Z^2(2\lambda + 1)!} q^{2\lambda},$$

(16)

where $\rho_\lambda(r, Z)$ is the nuclear transition density and $B(E\lambda)$ is the reduced transition probability. These form factors are equivalent for the inelastic process with $A_1$ nuclear transition $0 \rightarrow \lambda$ will be

$$N_{1el}^{(e)}(w, b) = \frac{Z^2 \alpha}{\pi^2} \frac{1}{b^2} \left[ \int \frac{u^2}{x^2 + u^2} J_1(u) F_\lambda[-(x^2 + u^2)/b^2] \, du \right]^2,$$

(17)

The equivalent photon number for the inelastic process with $A_1$ nuclear transition $0 \rightarrow \lambda$ will be

$$N_{1el}^{(e)}(w, b) = \frac{Z^2 \alpha}{\pi^2} \frac{1}{b^2} \left[ \int \frac{u^2}{x^2 + u^2} J_1(u) F_\lambda[-(x^2 + u^2)/b^2] \, du \right]^2,$$

(18)

as the generalization of (12). Here $x^2_{\text{in}} = (E_0^2 + \frac{u^2}{\gamma^2} + 2\frac{wE_0}{\gamma} + \frac{K_0^2}{\gamma}) \gamma^2$.

The value $q_0^2$ is close to $q_{\text{in}}^2$ at a large $\gamma_A$ factor at LHC energies.

We take the inelastic form-factor from inelastic electron scattering off nuclei. A good parameterization of the inelastic form-factor is

$$F_{\lambda}^2(q) = 4\pi \beta_{\lambda}^2 \frac{1}{b^2} F_{el}^2(qR) e^{-q^2 g^2}$$

(21)

in Helm’s model [13]. The squared transition radius is equal to $R_2^2 = R^2 + (2\lambda + 3)q^2$, where $R$ and $g$ are the model parameters. According to (19) the reduced transition probability in this case is equal to

$$B(E\lambda) = \frac{\beta_{\lambda}^2}{4\pi} Z^2 e^2 R_{2\lambda}^2.$$ 

(22)

So, the formulae for the process (1) are

$$d\sigma_{A_1 A_2 \rightarrow A_1^* A_2 X_f} = \frac{\int \frac{d\omega_1}{\omega_1} \int \frac{d\omega_2}{\omega_2} n_1^{(e)}(w_1)n_2^{(e)}(w_2)}{E_0^2 - q_{\text{in}}^2} \cdot d\sigma_{\gamma \gamma \rightarrow X_f}(w_1, w_2);$$

(23)

$$n_1^{(e)}(w_1) = \left( \frac{Z^2 \alpha}{\pi^2} \right) \frac{1}{b^2} F_{el}^2(q),$$

(24)

$$-q_{\text{in}}^2 = \frac{w^2}{\gamma_A^2} + \frac{2wE_0}{\gamma_A} + \frac{E_0^2}{\gamma_A^2} + q_0^2;$$

(25)

$$n_2^{(e)}(w_2) = \left( \frac{Z^2 \alpha}{\pi^2} \right) \frac{1}{b^2} F_{el}^2(q),$$

(26)

$$-q_{\text{in}}^2 = \left( \frac{w}{\gamma_A} \right)^2 + q_0^2.$$ 

(27)

The effective two-photon luminosity can be expressed as

$$L(\omega_1, \omega_2) = 2\pi \int \frac{b_1 db_1}{R_1} \int \frac{b_2 db_2}{R_2} \int \frac{d\phi}{0} \cdot$$

(28)

$$\cdot N_{1el}^{(e)}(\omega_1, b_1) N_{2el}^{(e)}(\omega_2, b_2) \Theta(B^2),$$

where $R_1$ and $R_2$ are the nuclear radii, $\Theta(B^2)$ is the step function and $B^2 = b_1^2 + b_2^2 - 2b_1 b_2 \cos \phi - (R_1 + R_2)^2$ [3]. Then the final cross-section is

$$\sigma_{A_1 A_2 \rightarrow A_1^* A_2 X_f} = \int \frac{d\omega_1}{\omega_1} \int \frac{d\omega_2}{\omega_2} L(\omega_1, \omega_2) \sigma_{\gamma \gamma \rightarrow X_f}(w_1, w_2).$$

(29)

### 2 Nuclear levels and form factors

The elastic form factor of a light nucleus is

$$F_{el}(q) = \exp\left(-\frac{(qR)^2}{6} q^2\right)$$

(30)
$30\%$ of the levels decay to the first 3 levels below the neutron threshold. They correspond to

\begin{equation}
\rho(r) = \rho_0 \left\{ \frac{1}{1 + \exp\left(\frac{r-R}{g}\right)} + \frac{1}{1 + \exp\left(\frac{r-R}{g}\right)} - 1 \right\} = \rho_0 \frac{\sinh(R/g)}{\cosh(R/g) + \cosh(r/g)},
\end{equation}

and allows us to calculate the elastic form factor analytically:

\begin{equation}
F_{el}(q) = \frac{4\pi^2 Rg\rho_0}{q \sinh(\pi qg)} \left\{ \frac{\pi qg}{R} \sin(qR) \coth(\pi qg) - \cos(qR) \right\}.
\end{equation}

There are a few discrete levels of $^{16}$O below the $\alpha$, $p$ and $n$ thresholds $E_{th}(\alpha) = 7.16$ MeV, $E_{th}(p) = 12.1$ MeV, $E_{th}(n) = 15.7$ MeV [15]. The level $2^+$ at $E_0 = 6.92$ MeV is the strongest excited one in the electron scattering.

The parameters from the inelastic electron scattering fit on $^{16}$O with excitation of $2^+$ level ($E_0 = 6.92$ MeV) are [16]

$\beta_2 = 0.30$, $R = 2.98$ fm, $g = 0.93$ fm.

They correspond to

\begin{equation}
B(E_{02}) = (36.1 \pm 3.4)c^2 \text{ fm}^4.
\end{equation}

Fig. 2. The elastic form factor (1) of $^{16}$O and the inelastic form factor (2) of $^{16}$O ($2^+, 6.92$ MeV) from the electron scattering.

$E_0 = 2.615$ MeV. This level is well studied experimentally [18] and has a large excitation cross-section.

\begin{equation}
B(E_{03}) = (6.12 \times 10^5 \pm 2.2\%)c^2 \text{ fm}^6.
\end{equation}

We calculate the parameter $\beta_3$, using this $B(E_{03})$, and take $R$ and $g$ from the density of the $^{208}$Pb ground state:

$\beta_3 = 0.113$, $R = 6.69$ fm, $g = 0.545$ fm.

Note that there are many levels higher than $E_0 = 2.615$ MeV which decay to the first level of $^{208}$Pb. This fact increases the event rate of the process (1), but we do not know the excitation cross-section of these levels.

The elastic form factor (30) of $^{16}$O and the inelastic form-factor of $^{16}$O ($2^+, 6.92$ MeV) (21), corresponding to the electron scattering data, are shown in fig. 2. The same for a nucleus $^{208}$Pb and the excited state $^{208}$Pb ($3^-, 2.64$ MeV) are shown in fig. 3.

The squared inelastic form factor is less than the elastic form factor by more then two orders at small $q < q_0$ ($q_0 = 0.5$ fm$^{-1}$ for $^{16}$O and $q_0 = 0.4$ fm$^{-1}$ for $^{208}$Pb). In the region of $q > q_0$ they are comparable. The region of large $q > q_0$ will contribute to the small impact parameter $b$. We are able to calculate the photon luminosity (28) for all regions of $b$ to get the maximum electromagnetic cross-section of the process we are interested in. Then it should be possible to compare with experimental data in condition of clear selection of such process by the photon signal and the veto neutron or proton signal in the ZDC.

3 Angular and energy distributions of secondary nuclear photons

We suppose that the nucleus $A_1^*(\lambda \mu)$ in process (1) is unpolarized. At this point now we do not know the relative excitation probability of $|\lambda \mu\rangle$ states, where $\mu$ is a projection of spin $\lambda$. This assumption needs further study in the
future. So we use a formula (27) in our work [10] for the angular distribution of secondary photons, which is valid for isotropic photon distribution in the rest system of $A_{\gamma}^1$ according to equal probabilities of excitation.

If we calculate the integral cross-section of reaction (1) using eq. (29), then the angular and energy distribution of photons are equal to [10]

$$d\sigma_{A_{\gamma}^1} d\Omega_{A_{\gamma}^1} = \frac{2\gamma_{A_{\gamma}^1}^2 \sin \theta_{\gamma}}{(1 + \gamma_{A_{\gamma}^1}^2 \tan^2 \theta_{\gamma})^2 \cdot \cos^3 \theta_{\gamma}},$$

$$d\sigma_{A_{\gamma}^1} dE_{\gamma} = \frac{\Theta(2\gamma_{A_{\gamma}^1} E_0 - E_{\gamma})}{2\gamma_{A_{\gamma}^1}^2 E_0},$$

where $\Theta(x)$ is the step function.

The angular distribution does not depend on the photon energy and the energy distribution is uniform.

The photon energy $E_{\gamma}$ and the polar angle $\theta_{\gamma}$ in the laboratory system are defined as

$$E_{\gamma} = \gamma_{A_{\gamma}^1} E_0 (1 + \cos \theta_{\gamma})$$

$$\cos \theta_{\gamma} = 1$$

$$\tan \theta_{\gamma} = \frac{1}{\gamma_{A_{\gamma}^1}^2 \tan^2 \theta_{\gamma}}$$

where $\theta_{\gamma}$ and $\theta_{\gamma}$ are the polar angles of the nuclear photon in the rest nuclear system and in the laboratory system with an axis $z||p_{A_{\gamma}^1}$. The photon energy $E_{\gamma}$ dependence on $\theta_{\gamma}$ is shown in fig. 4. Thus the energy $E_{\gamma}$ will depend on the position of photon hit.

Our calculations with the TPHIC event generator [19] show that a deflection of the direction $p_{A_{\gamma}^1}$ from $p_{beam}$ at LHC energies in the reaction (1) is very small at large $\gamma_{A_{\gamma}^1}$, $\langle\Delta \theta\rangle \approx 0.5$ $\mu$rad.

In the experiments CMS and ALICE, which are planned at LHC (CERN), the zero-degree calorimeter [20, 21] was suggested for the registration of nuclear neutrons after ion interaction. We demonstrate a schematic figure of the ZDC CMS at a distance $L = 140$ m in the plane transverse to the beam direction in fig. 5. The CMS group also plans to include the electromagnetic calorimeter in front of the ZDC.

As an example, we show the angular distributions (36) in arbitrary units and the energy dependence (38) on the $(x, y)$ coordinates of the ZDC CMS for the two nuclei $^{16}\text{O}$ and $^{208}\text{Pb}$ in fig. 6. The direction of the nucleus $A_{\gamma}^1$ coincides here with the beam direction. The point $(x, y) = (0, 0)$ is the center of the ZDC plane.

4 Cross-section of the process with the nuclear $\gamma$ radiation

We demonstrate our results for the $\eta_c(2.979)$ production. The previous results [3] used old values of the widths and a point nuclear charge. Now we take resonance parameters.
Fig. 6. The photon angular distributions (upper row) and the energy dependence (lower row) for $^{16}$O$^\ast$(2+, 6.92 MeV) (left column) and $^{208}$Pb$^\ast$(3−, 2.62 MeV) (right column) radiation decay in the laboratory system on the ZDC plane $(x, y)$ at 140 m distance from point interaction. $x$, (cm) is the horizontal and $y$, (cm) is the vertical axis. The photon energy interval in the ZDC region is 19–48 GeV for $^{16}$O$^\ast$(2+) and 7–14 GeV for $^{208}$Pb$^\ast$(3−).

from the review of particle physics [22], $\Gamma_{\gamma \gamma \rightarrow \gamma} = 4.8$ keV, and a realistic charge distribution. The calculations was made with the help of TPHIC event generator [19].

We use a well-known formula [2] of the narrow resonance cross-section:

$$\sigma_{\gamma \gamma \rightarrow \chi}(w_1, w_2) = 8\pi^2(2\lambda_X + 1)\Gamma_{\gamma \gamma \rightarrow \gamma\chi}(W^2 - M_X^2)/M_X,$$

(40)

where $W^2 = 4w_1w_2$, $\lambda_X$ and $M_X$ is the spin and mass of the resonance. The LHC luminosity and our results according to (29) and (28) are in table 1 for the process (1) with $A_{\text{final}} = A_1$ or $A_1^\ast$.

Our results in table 1 show that though the cross-section of the process (1) for the nucleus $^{208}$Pb is larger than that for $^{16}$O, the event rate is smaller because of the lower LHC luminosity for $^{208}$Pb. The cross-section with a nuclear excitation is smaller by three orders of magnitude than that without the excitation, since the intensity of excitation is not large and the inelastic form factor is smaller than the elastic form factor (see figs. 2 and 3). Therefore for the accepted LHC luminosities it is possible to use secondary photons as a signature of clear electromagnetic nuclear processes only for the production $X_f$ with rather large cross-section $\sigma_{\gamma \gamma \rightarrow \chi}$. Light ions are more preferable than heavy ions to detect the nuclear $\gamma$ radiation.

5 Conclusion

In this work we suggest a new signature of the peripheral ion collisions.
The formalism of the process (1) is developed in the frame of the equivalent photon approximation. The new point is the introduction of the inelastic nuclear form factor. It allows to consider the excitation of discrete nuclear levels and their following decay. It is shown that the energy of this secondary photons are in the GeV domain. The angular distribution of the photons has a peculiar form as a function of polar angle in the beam direction. The majority of photons fly in the region of angles of a few hundred micro-radians, which are detectable in the ZDC CMS and ALICE experiments.

Thus the nuclear $\gamma$ radiation is a good signature of ultra-peripheral collisions at LHC energies when $A$ and $Z$ of the beam ion are conserved. The trigger requirements will include a signal in the central rapidity region of particles from $X_f$ decay, a signal of photons in the electromagnetic detector in front of the zero-degree calorimeter and a veto signal of neutrons in the ZDC. We suggest to use the veto signal of the neutron in order to avoid the processes with nuclear decay into nucleon fragments. The nuclear $\gamma$ radiation can be used for tagging the events with particle production in the central rapidity region in ultra-peripheral collisions.

Light nuclei are more preferable in comparison with heavy ions, since they have higher beam luminosity at LHC. The cross-sections of the process with the nuclear excitation are three orders of magnitude smaller than the one without excitation. The accepted nuclear luminosities enable us to use this signature for the large cross-section of the $X_f$ system production.

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<table>
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<th>$A_{\text{final}}$</th>
<th>$L$ (cm$^{-2}$ s$^{-1}$)</th>
<th>$L$ (pb$^{-1}$)</th>
<th>$\sigma$</th>
<th>event/10$^6$ s</th>
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<td>$^{208}\text{Pb}_{82}$</td>
<td>4.2 $\cdot$ 10$^{26}$</td>
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<td>356 $\mu$b</td>
<td>147000</td>
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<tr>
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