Chapter 1

Introduction

1.1 Preface

Already in 1949 it was suggested by Ashkin and Marshak [Ash49] that in high-energy heavy-ion collisions, in addition to statistical photons and photons originating from meson decay, a third source of photons might be present associated with nucleon-nucleon collisions. Only in 1985 was nuclear bremsstrahlung in heavy-ion collisions actually observed by Grosse et al. [Gro85]. From this experiment it was not yet clear whether the origin of the bremsstrahlung was the nucleon-nucleon or the nucleus-nucleus system. Since this first experiment many experiments have been performed to determine the nature of the bremsstrahlung and it was found that the bremsstrahlung originates from incoherent nucleon-nucleon collisions.

In the present work nuclear bremsstrahlung has been used as a probe to study dissipation mechanisms in peripheral heavy-ion reactions at intermediate energies, i.e. energies close to the Fermi energy ($E_F \approx 35$ MeV). For energies well below the Fermi energy it has been established that the dissipation mechanism is governed by the nuclear mean field while at energies much higher than the Fermi energy individual nucleon-nucleon collisions govern the dissipation mechanism. It has been shown (see e.g. [Gue85]) that at intermediate energies both dissipation mechanisms, the mean field as well as the individual nucleon-nucleon collisions, play an important role.

It is the aim of the present work to determine the relative importance of both dissipation mechanisms by studying the reaction: $^{36}$Ar + $^{152}$Tb at a beam energy of 44 MeV/nucleon, where the nuclear bremsstrahlung is used as a signature for nucleon-nucleon collisions. In the following sections a brief description of both dissipation mechanisms will be given and the nature of the nuclear bremsstrahlung will be discussed.
1.2 Dissipation mechanisms

To transfer nucleons from projectile to target or vice versa the relative motion of the transferred nucleon has to be converted to an intrinsic motion, i.e. the energy difference with the receiving nucleus has to be dissipated. In intermediate-energy heavy-ion collisions the energy can be dissipated in two ways:

1. The transferred nucleon interacts with the mean field produced by the sum of all nucleons in the system. This is referred to as one-body dissipation.

2. The transferred nucleon interacts with one other nucleon in the system. This process is denoted two-body dissipation.

Thus in the case of one-body dissipation the mean field is the driving force for the nucleon transfer while in two-body dissipation incoherent nucleon-nucleon collisions are assumed to be the origin of the dissipation mechanism. In the next chapter both mechanisms will be explained more extensively with the help of transport-model calculations using the BUU-code.

![Figure 1.1: Fermi spheres for three beam energies: 5 MeV/nucleon, 44 MeV/nucleon and 100 MeV/nucleon. The full circles correspond to the momentum space uniformly filled up to the Fermi momentum and the dashed circles show the highest bound states.](image)

The influence of both mechanisms can be qualitatively understood by very simple pictures showing two Fermi spheres, i.e. spheres in momentum space corresponding to the target and the projectile nucleons with radii equal to the respective Fermi momenta, displaced by the Coulomb-corrected beam.
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momentum. In figure 1.1 the decreasing importance of one-body dissipation is shown. The full circles in figure 1.1 correspond to the momentum space of projectile and target uniformly filled up to the Fermi momentum while the dashed circles show the highest bound states. A one-body transfer within the hatched region is only possible if two nucleons in the same state are exchanged which has no net effect. The only possible one-body transfer that can take place, under the condition that the nucleons after the transfer are bound, is into the grey region. It is clear from the figure that this region decreases with increasing beam energy going to zero for energies larger than 100 MeV/nucleon.

![Figure 1.2: Fermi spheres for three beam energies: 5 MeV/nucleon, 44 MeV/nucleon and 100 MeV/nucleon. The full circles correspond to the momentum space uniformly filled up to the Fermi momentum, the dashed circles show the highest bound states and the dotted circles show the maximum momentum a nucleon can obtain due to a collision.](image)

Figure 1.2 shows the increasing effect of two-body dissipation. The full and dashed circles have the same meaning as in figure 1.1 and the dotted circles correspond to the maximum momentum a nucleon can obtain due to a nucleon-nucleon collision. For the two-body transfer again the hatched area is forbidden, since Pauli blocking allows only a collision if the two nucleons go back to their original state which has no net effect. Therefore, with the same condition as for the one-body transfer, that the nucleons after the collision have to be bound, the grey area is the only available area for the nucleons to transfer to after the collision. Here, we find the opposite behaviour as compared to the one-body case since for the two-
body transfer the number of available states increases with increasing beam energy. From figure 1.2 it might seem that already at low energies the two-body mechanism is dominant. Therefore, it needs to be noted that the momenta of the two nucleons after the collision have to match precisely to the momenta of the available states (the gray region); and although there are already a large number of states available in low-energy collisions, there are only few nucleon-nucleon collisions possible for which the relative momentum matches to that of the available states. There are also many collisions leading to unbound states. These collisions would lead to pre-equilibrium emission if no further interaction of the collided nucleons with the projectile or the target takes place.

Thus from this simple picture it is seen that both mechanisms play a role at intermediate-energy heavy-ion collisions and that the relative importance changes with bombarding energy.

1.3 Nuclear bremsstrahlung

As mentioned in the preface the first experiment studying nuclear bremsstrahlung in a heavy-ion reaction was performed by Grosse et al. in 1985 [Gro85]. Although from this measurement it was not clear whether these hard photons originated from nucleon-nucleon collisions or nucleus-nucleus collisions, the hard-photon measurement was assumed to give an undistorted view of the early phase of the reaction in contrast to nucleons or nuclear fragments which suffer from final-state interactions and for which it is not clear whether they originate from the participant or the spectator part of the colliding nuclei. After this pioneering work many hard-photon measurements were performed [Bea85, Gro86, Ala86, Kwa86, Ste86, Ber87, Bre89, Cla89, Mur89, Tam89, Cla90, Gos90, Hof91, May93] in which the inclusive photon production cross section was measured as a function of energy and angle. Two important characteristics of the hard photons were found in these experiments:

1. The photon spectra showed an exponential behaviour (see 1.3.1).

2. The slopes of the photon spectra were found to depend on the angle of observation (see 1.3.2).

The second characteristic suggests that the photons are emitted from a moving source and on that basis a moving source fit to the angular distribution can be done as suggested by Bertholet et al. [Ber87]. By fitting with
this parameterization the angular distributions from different experiments it was found that the velocity of the emitting source was approximately half the beam velocity, which is equal to the nucleon-nucleon center-of-mass velocity, suggesting that the origin of the hard photons is nucleon-nucleon collisions.

After these first experiments, exclusive hard-photon measurements were performed measuring the photon characteristics as a function of the mass of the outgoing fragment [Hin87, Rie92, Mar94a, Pol95, Pol96] and as a function of the charged-particle multiplicity [Her88, Kwa88, Rep92, Mig93]. The scaling of the hard photons with the projectile mass led to the conclusion that the hard photons originate from incoherent, first-chance, nucleon-nucleon collisions.

In the following the different aspects of the hard-photon production will be discussed and its systematics will be shown.

1.3.1 Energy spectra

As mentioned above the energy spectra of the hard photons in the nucleon-nucleon center of mass exhibit an exponential shape which can be characterized by an inverse slope parameter $E_0$.

$$\frac{d\sigma}{dE_\gamma} \propto \exp\left(-\frac{E_\gamma}{E_0}\right)$$  \hspace{1cm} (1.1)

Already after the first experiments a correlation between the inverse slope parameter and the beam energy was suggested [Ber87]. This description was modified by Metag [Met91] correcting the beam energy for the Coulomb repulsion between projectile and target ($\varepsilon_{Cc}$).

$$\varepsilon_{Cc} = T_{beam} - \varepsilon_C$$  \hspace{1cm} (1.2)

where $T_{beam}$ is the kinetic beam energy per nucleon, $\varepsilon_C = \frac{Z_p Z_t e^2}{\mu r}$, $\mu = \frac{A_p A_t}{A_p + A_t}$ and $Z_p$, $Z_t$ and $A_p$, $A_t$ are the atomic numbers and masses of projectile and target, respectively. In the following, when we refer to the beam energy the Coulomb-corrected beam energy is meant. In figure 1.3 the systematic behaviour of the slope parameter as a function of beam energy is shown. The curve drawn in the figure is a fit to the data of the function

$$E_0 = a \cdot (\varepsilon_{Cc})^b$$  \hspace{1cm} (1.3)
where $a = 0.48 \pm 0.06$ and $b = 0.91 \pm 0.03$. From this behaviour one can conclude that the slope of the photon spectrum mainly depends on the beam energy and does not, in first order, depend on the specific projectile and target combination.

1.3.2 Angular distribution

The angular distribution for nuclear bremsstrahlung can be explained by looking at the bremsstrahlung amplitudes for the simplest two systems, proton-neutron and proton-proton collisions (see e.g. [Nif85, Nif89]). We will limit ourselves to the angular dependence of the bremsstrahlung amplitudes since the full expression is very sensitive to the approximations (e.g. the soft-photon approximation) made in the derivation. For the proton-
neutron case the bremsstrahlung amplitude can be written as
\[ |A_{pn}(\Theta, \omega)|^2 \propto \left( \sin^2 \Theta + \frac{2}{3} \right) \]
where \( \omega \) is the photon energy and \( \Theta \) is the photon angle. From this equation it can be seen that the angular distribution consists of a dipole \( \left( \sin^2 \Theta \right) \) and an isotropic part \( \left( \frac{2}{3} \right) \). This can be understood in terms of an oscillating point charge: the proton suddenly slows down \( (\Delta \tau \leq \frac{\Delta r}{v_{beam}} \approx 1 \text{ fm/c}, \) where \( \Delta r \) is the range of the nucleon-nucleon force) in the collision and then accelerates again. The radiation pattern for such a point charge has a dipole character. In the process of slowing down, the proton has a well defined direction (beam direction) and thus this part of the bremsstrahlung will exhibit a dipole nature in the laboratory frame. After the collision the direction of the proton is isotropic and therefore the total bremsstrahlung distribution will be a superposition of an isotropic and a dipole term.

For the proton-proton case the bremsstrahlung amplitude is
\[ |A_{pp}(\Theta, \omega)|^2 \propto \left( \sin^2 \Theta \cos^2 \Theta + \frac{2}{15} \right) \]
In this case we find for the angular distribution a superposition of a quadrupolar \( \left( \sin^2 \Theta \cos^2 \Theta \right) \) and an isotropic term \( \frac{2}{15} \). Comparing both amplitudes one finds that the proton-proton bremsstrahlung is suppressed relative to the proton-neutron bremsstrahlung by a factor [Nif89]
\[ \frac{|A_{pp}(\Theta, \omega)|^2}{|A_{pn}(\Theta, \omega)|^2} = \beta^2 + \left( \frac{E_\gamma}{197} \right)^2 \Delta r^2 \]
where \( E_\gamma \) is the photon energy and \( \beta \) is the velocity of the incoming nucleon. For photon energies less than 100 MeV the dominating term is \( \beta^2 \) which is in the order of 0.1. This means that proton-proton bremsstrahlung is suppressed by a factor 10 and therefore we will only consider proton-neutron bremsstrahlung in this work.

In heavy-ion experiments the situation is more complex in the sense that the motion of the nucleons of projectile and target is smeared out in all directions due to the Fermi motion, which would suggest that the angular distribution is more isotropic. This is indeed found in the experiments and it was suggested [Ber87] to parametrize the angular distribution by an isotropic term and an anisotropic dipole term.
\[ \left( \frac{d^2\sigma}{dE d\Omega} \right)_{cm} = K \left( 1 - \alpha + \alpha \sin^2 \Theta_{cm} \right) e^{-E_{cm}/E_0} \]

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where $E_{cm}$ and $\Theta_{cm}$ are the photon energy and angle in the source frame, $\alpha$ is a measure for the anisotropy and $K$ a normalization constant. One can, furthermore, use the angular distribution to extract the velocity of the emitting source. For this a transformation to the laboratory system is needed, resulting in equation 1.8.

**Figure 1.4:** Systematics of the source velocities versus the beam velocity for asymmetric systems ($|\frac{A_p - A_t}{A_p + A_t}| \geq \frac{1}{3}$). For the drawn line the source velocity is half the beam velocity i.e. the nucleon-nucleon center-of-mass velocity. The data point from the present work is indicated by $\bigoplus$.

\[
\left( \frac{d^2\sigma}{dE d\Omega} \right)_{lab} = \frac{K}{X} \left( 1 - \alpha + \alpha \frac{\sin^2 \Theta_{lab}}{X^2} \right) e^{-X E_{lab}/E_0} \tag{1.8}
\]

with $X = \gamma(1 - \beta_s \cos \Theta_{lab})$ and $\beta_s$ is the source velocity. In figure 1.4 the
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systematics of the source velocity versus the beam velocity are shown. The line shown in the figure is \( \beta_s = \frac{1}{2} \beta_{\text{beam}} \) confirming that the origin of the hard photons is the nucleon-nucleon center of mass. The systematics hold over a large range of beam energies and systems. However in recent work [Sch94] the source velocity has been studied as a function of the system size, photon energy and impact parameter. Two asymmetric systems were studied, \(^{36}\text{Ar} + ^{12}\text{C}\) and \(^{36}\text{Ar} + ^{197}\text{Au}\), at a beam energy of 95 MeV/nucleon. For the \(^{12}\text{C}\)-target the source velocity was similar to the nucleon-nucleon center-of-mass velocity. For the \(^{197}\text{Au}\) target, a clear deviation from the systematics was found at small impact parameters showing that these hard photons also reflect the later, stopping-stage of the reaction. For the larger impact parameters the source velocity still agrees with the systematics as shown in figure 1.4.

1.3.3 Inclusive photon production probability

Having established the origin of the hard photons it is of interest to compare the photon production cross section \( (\sigma_\gamma) \) found in the different experiments. A commonly used variable to do this is the photon production probability per proton-neutron collision \( (P_{\gamma}^{pn}) \), which is defined as follows:

\[
P_{\gamma}^{pn} = \frac{\sigma_\gamma}{\sigma_R \cdot <N_{pn}>_b} \tag{1.9}
\]

where \( \sigma_R \) is the reaction cross section and \( <N_{pn}>_b \) is the number of pn-collisions averaged over the impact parameter. The latter is calculated in the framework of the equal participant model proposed by Nifenecker and Pinston [Nif89]. In this model the geometrical overlap of the two colliding nuclei \( (A_F(b)) \) is calculated as a function of the impact parameter \( (b) \) assuming hard sphere radii.

\[
A_F(b) = \frac{1}{2} A_p (2 - 3 \cos \Theta_p + \cos^3 \Theta_p) + \frac{1}{2} A_t (2 - 3 \cos \Theta_t + \cos^3 \Theta_t) \tag{1.10}
\]

with \( \cos \Theta_p = \frac{b^2 + r_p^2 - r^2}{2br_p} \), \( \cos \Theta_t = \frac{b^2 + r_t^2 - r_p^2}{2br_t} \) and \( r_p \) and \( r_t \) are the projectile and target radius, respectively. Furthermore, it is assumed that all nucleons in the overlap zone collide only once. The impact parameter averaged number of pn-collisions then becomes
\begin{equation}
<N_{\gamma n}>_b = A_p \frac{5 A_t^{2/3} - A_t^{2/3}}{5(A_t^{1/3} + A_t^{1/3})^2} \frac{Z_p N_t + Z_t N_p}{A_p A_t}
\tag{1.11}
\end{equation}

The derivation of this expression is shown in Appendix B.1. The reaction cross section used to calculate the systematics has been derived in Appendix B.2. The resulting equation becomes

\begin{align}
\sigma_R &= \pi R^2 \left( 1 - \frac{E_t}{T_{beam}} \right) \\
R(\text{fm}) &= 1.16 \left( A_p^{1/3} + A_t^{1/3} + 2.0 \right)
\tag{1.12}
\tag{1.13}
\end{align}

For the radius (R) an expression suggested by Metag [Met89] has been used. The systematics of the inclusive photon production probability are shown in figure 1.5 (in appendix E the values are listed in a table). The full line fitted to the systematics assumes that the photon production probability only depends on the inverse slope parameter and the threshold energy.

\begin{equation}
P_{\gamma n}^{\text{ppm}} = P_{0}^{\text{nm}} \cdot \exp \left( - \frac{E_t}{E_0} \right)
\tag{1.14}
\end{equation}

where \(E_t\) is the lower energy threshold for integrating the photon spectrum \((E_t = 30 \text{ MeV})\) and \(E_0\) is taken from equation 1.3. The fitted value for \(P_{0}^{\text{nm}}\) is \(6.3 \pm 0.1 \times 10^{-4}\).

It should be noted that this systematics is based on the assumption that all participants collide at least once. We will show in this work that this assumption is not correct.

### 1.3.4 Photon multiplicity

In the previous section the inclusive photon production probability was defined and compared to the systematics. More exclusively one can measure the photon multiplicity \(M_{\gamma}(R)\) associated with a given reaction channel \(\langle R \rangle\) which is defined as the photon cross section \(\langle \sigma_{\gamma} \rangle\) divided by the reaction cross section \(\langle \sigma_R \rangle\).

\begin{equation}
M_{\gamma}(R) = \frac{\langle \sigma_{\gamma}(R) \rangle}{\langle \sigma_R(R) \rangle}
\tag{1.15}
\end{equation}

The approach chosen in previous work is to correlate the photon multiplicity...
with the impact parameter. This procedure is illustrated in figure 1.6; the data are taken from reference [Rie92]. The first step is to determine from the photon multiplicity the number of pn-collisions. This is done by combining equations 1.15 and 1.9 which results in

$$N_{pn}(b) = \frac{M_{\gamma}(b)}{P_{\gamma}^{pn}}$$  \hspace{1cm} (1.16)
where $P_{\gamma}^{\text{pn}}$ is the inclusive photon production probability per pn-collision taken from the inclusive systematics shown in figure 1.5. It should be noted that the inclusive $P_{\gamma}^{\text{pn}}$ already relies on the validity of the geometrical overlap model as described in the previous section.

The second step is to correlate the number of pn-collisions with the impact parameter for which equation 1.10 is used. This is shown in the right part of figure 1.6. The result is a unique correlation between the photon multiplicity and the impact parameter. However, this result depends very strongly on the geometrical overlap model.

In the present work a different approach will be chosen since we want to measure the photon production probability per removed unit of mass without relying on the geometrical overlap model. To do this the photon multiplicity is measured as a function of the mass of the primary projectile-like fragment (in chapter 4 it will be explained how the primary projectile-like fragment mass is obtained). Now one can define the exclusive photon
production probability per removed nucleon as follows:

$$p_{\gamma}^{exc,NN} = \frac{dM_\gamma}{dA}$$

(1.17)

In chapter 5 the results of this new approach will be discussed and compared to the previously used methods.

1.4 Motivation

It is the aim of this work to study the dissipation mechanism for peripheral heavy-ion reactions at intermediate energies using the hard-photon probe. The improvement compared to previous experiments performed along these lines [Hin87, Rie92] is twofold.

1. Fragments are generally produced in a particle-unstable state and will therefore subsequently decay via the evaporation of light particles. This evaporation step washes out part of the dissipation mechanism since the measurements do not allow to disentangle the actual reaction from the evaporation stage.

2. To determine exclusive values for the photon production probability ($P_\gamma$) from the photon production cross section, which was measured in these experiments, it is necessary to know the average number of participants. To calculate this number the authors relied on the participant spectator model which is based on the assumption that only nucleons in the overlap region will contribute to the nucleon-nucleon collisions, thus neglecting the role of the mean field.

In the experiment described in this thesis the evaporated light charged particles are measured in addition to the projectile-like fragment making it possible to determine the projectile-like fragment before particle evaporation which is a more direct probe of the dissipation mechanism. Also the participant spectator model is no longer needed because the fragment before the evaporation stage is a direct measure for the mass difference between projectile and projectile-like fragment. In this way the influence of the N/Z ratio of the target on the produced fragments can be measured, since the evaporation will generally populate the strongest bound fragments.
1.5 Outline

In chapter 2 the transport model, BUU, and the most important features of both dissipation mechanisms will be discussed. In chapter 3 the experimental setup will be described and the performance of the individual detectors will be shown. In chapter 4 the reconstruction method of the fragment before the evaporation stage will be examined in great detail. The results as well as a comparison to previous work will be presented in chapter 5. Finally in chapter 6 the main conclusions will be drawn and an outlook for future work will be given.