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Fotoexcitação Nuclear

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Resumo

Este trabalho apresenta uma descrição breve dos processos de excitação do núcleo atômico pela radiação eletromagnética e do tipo de informação que se consegue obter por esses métodos. O processo de monocromação de fótons de bremsstrahlung, conhecido por marcação de fótons, é descrito e discutido, bem como o cálculo e a integração da seção de choque diferencial de bremsstrahlung, necessários para determinar as características do equipamento. Por fim apresentamos alguns exemplos de experimentos que poderão ser realizados no LAL com o uso dos novos equipamentos em desenvolvimento.

Abstract

This work presents a brief description of the photoexcitation mechanisms of the nucleus, and of the information that can be derived through them. The photon tagging technique is presented and discussed, as well as the differential bremsstrahlung cross section, which must be calculated and integrated in order to simulate the performance of the equipment. Some examples of the new experimental perspectives of the lab are presented at the end.

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1. Introdução

O uso de sondas eletromagnéticas (elétrons ou fótons) no estudo das propriedades do núcleo atômico apresenta algumas características positivas fundamentais. Inicialmente, pelo processo de excitação, que é devido à bem conhecida interação eletromagnética dos elétrons ou fótons com as densidades de carga, corrente e magnetização do núcleo. Além disso a interação é fraca, permitindo que o processo seja tratado em teoria de perturbação. Dessa forma, informações precisas e detalhadas, muitas vezes independentes de modelos, podem ser obtidas sobre a estrutura do núcleo atômico através do processo de fotoexcitação, seja com fótons reais ou virtuais.

O presente trabalho restringe-se a tratar a parte da fotoexcitação nuclear em baixas energias, ou seja, a excitação de estados ligados e a região das ressonâncias gigantes, que constituem a essência de meu trabalho de pesquisa. Apesar da limitação em energia, trata-se de região muito rica em fenômenos interessantes, que espelham propriedades e comportamentos fundamentais do núcleo.

A construção de aceleradores lineares de elétrons estimulou o desenvolvimento de investigações de estrutura nuclear e o estudo da dinâmica de reações nucleares induzidas por fótons e elétrons. No entanto, nos aceleradores lineares tradicionais (com feixe pulsado), apenas os experimentos fora de linha, ou os em linha baseados na detecção de partículas no plano focal de espectrômetros magnéticos tiveram sucesso significativo. Toda uma categoria de experimentos, baseada na detecção de partículas em condições geométricas de visão direta do alvo, não foi assimilada por causa do caráter pulsado do feixe e do conseqüente pulso de ruído (eletrônico e radioativo) que acompanha o pulso de corrente. O desenvolvimento de aceleradores de onda contínua (*cw*) está mudando essa situação, permitindo a realização de experimentos em coincidência, com cinemática completa e o uso de detectores com alto ângulo sólido, por serem posicionados próximos do alvo.

Neste trabalho são apresentados resultados de experimentos inclusivos de eletro e fotodesintegração nuclear realizados no LAL e em outros laboratórios, com aceleradores lineares pulsados. Apresentamos também resultados obtidos no desenvolvimento de

equipamentos e processos visando a instalação, no LAL, de um sistema apto a realizar medidas exclusivas, em linha, de espalhamento elástico e inelástico de fótons e fotodesintegração, utilizando fótons monocromáticos e polarizados.

No Cap. 2 apresento uma descrição geral das principais características da fotodesintegração nuclear e sua relação com os processos de eletrodesintegração e de espalhamento de elétrons, exemplificando os métodos de medida utilizados e os resultados típicos obtidos com esses métodos, enfatizando aqueles com que trabalhei diretamente ou que apresentavam grande interseção com meu trabalho.

O Cap. 3 traz uma descrição detalhada do método de monocromação de fótons de bremsstrahlung conhecido por marcação de fótons, ou *photon tagging*. Apresento suas principais características e potencialidades, uma vez que esse processo é fundamental para a implementação das novas linhas de pesquisa do LAL. Discuto também a produção de fótons linearmente polarizados.

No Cap. 4 descrevo a seção de choque diferencial de bremsstrahlung, apresentando suas características mais relevantes ao processo de monocromação. Descrevo o programa de simulação por Monte Carlo desenvolvido para facilitar o estudo das propriedades da seção de choque de bremsstrahlung e otimizar o uso do monocromador.

No Cap. 5 descrevo as perspectivas experimentais que estão se abrindo no LAL com a introdução do monocromador de fótons e de periféricos associados que estão sendo desenvolvidos simultaneamente.

As conclusões são apresentadas no Cap. 6.

Os trabalhos que formam o corpo da tese são reunidos no anexo, numerados sequencialmente conforme aparecem no texto. Além dessa numeração, apresento ao lado, entre parênteses, a numeração em que eles aparecem no Memorial, para facilitar a correlação.

2. Fotodesintegração nuclear

2.1. Introdução

Nesse capítulo é apresentada uma descrição básica dos processos de excitação eletromagnética do núcleo por fótons reais e virtuais, com ênfase em algumas características do espalhamento de elétrons e fótons pelo núcleo, que são relevantes ao trabalho desenvolvido por mim no LAL.

No item 2.2 apresento uma descrição do espalhamento elástico e inelástico de fótons pelo núcleo. Em seguida discuto o espalhamento de elétrons pelo núcleo e mostro a relação entre os processos de foto e eletroexcitação, que constitui a base do método dos fótons virtuais.

2.2. Espalhamento elástico e inelástico de fótons

O espalhamento ressonante de fótons é uma ferramenta poderosa para extrair informações precisas sobre algumas propriedades do núcleo atômico. Como nesse caso os processos de excitação e desexcitação nuclear ocorrem via interação eletromagnética, esse método permite a determinação, independente de modelos, de informações sobre estados nucleares, tais como:

- energia de excitação, via espectro em energia dos fótons espalhados;
- spin, através da distribuição angular;
- largura radiativa, através da seção de choque; e
- paridade, através da determinação do grau de polarização linear da radiação emitida.

Uma das características do espalhamento de fótons pelo núcleo é sua alta sensibilidade à excitação de estados de baixo spin. Ocorrem essencialmente processos dipolares ($E1$ e $M1$), com uma modesta contribuição quadrupolar, devido ao baixo momento transferido ao núcleo pelo fóton. Somando-se a isso as técnicas que a moderna

espectroscopia gama oferece, como detectores de Ge de alta eficiência e excelente resolução em energia, as medidas de espalhamento de fótons podem alcançar níveis de sensibilidade muito altos.

Consideremos a situação esquematizada no diagrama da Fig. 1, em que um núcleo de spin J_0 , irradiado por um espectro contínuo de fótons, é excitado para um nível de energia E_x , de spin J , largura radiativa total Γ e largura radiativa para o estado fundamental Γ_0 . A seção de choque de espalhamento, integrada sobre um nível isolado é dada por:

$$\frac{d^2\sigma(\vec{\gamma}, \gamma')}{d\Omega} = \frac{2J+1}{2J_0+1} \frac{\pi^2}{k^2} \frac{\Gamma_0 \Gamma_f}{\Gamma} \frac{W(\theta, \phi)}{4\pi}, \quad (1)$$

onde k é o número de onda do fóton absorvido; $W(\theta, \phi)$ representa a distribuição angular normalizada e Γ_f a largura radiativa de decaimento para o estado final E_f, J_f . No caso da excitação por fótons não polarizados a dependência de $W(\theta, \phi)$ com o ângulo ϕ desaparece, $W(\theta, \phi) = W(\theta)$.

Figura 1 — Esquema do processo de excitação e desexcitação: (1) e (2) indicam os canais de entrada e saída, respectivamente. Γ representa a largura radiativa, λ a ordem multipolar e J o spin.

Através da largura total Γ , a seção de choque depende das propriedades de todos os estados existentes entre E_x e o estado fundamental. Além disso, no primeiro estado excitado, $\Gamma = \Gamma_0$, uma vez que esse nível só pode decair para o fundamental. À medida que a energia de excitação aumenta, cada novo estado que se abre pode decair para todos os inferiores. Dessa forma, a largura radiativa total de cada nível, que é a soma de todas suas larguras radiativas parciais, vai aumentando com o aumento da energia de excitação.

A largura radiativa para o estado fundamental de um nível de spin J , carrega informações sobre características do núcleo, através de sua dependência com as probabilidades de transição reduzida, $B(E\lambda, E)$ e $B(M\lambda, E)$:

$$\Gamma_0 = 8\pi\alpha \sum_{\lambda=1}^{\infty} \frac{(\lambda+1)E^{2\lambda+1}}{\lambda[(2\lambda+1)!!]^2} \frac{2J_0+1}{2J+1} [B(E\lambda, E) + B(M\lambda, E)], \quad (2)$$

onde λ é a multipolaridade da transição. Devido ao baixo momento transferido ao núcleo pelo fóton, usualmente apenas o primeiro termo da somatória é relevante. A largura radiativa para o estado fundamental está também relacionada com a seção de choque de absorção de fótons pelo núcleo. A seção de choque de absorção, integrada sobre a largura de um nível, é dada por:

$$\int \sigma(E) dE = (2\pi)^3 \alpha \sum_{\lambda=1}^{\infty} \frac{(\lambda+1)E^{2\lambda+1}}{\lambda[(2\lambda+1)!!]^2} [B(E\lambda, E) + B(M\lambda, E)] = \frac{2J+1}{2J_0+1} \pi^2 \Gamma_0. \quad (3)$$

Sofisticando um pouco mais nossa descrição do esquema da Fig. 1, podemos considerar que haja uma mistura multipolar na radiação tanto no canal de entrada quanto no de saída. Nessas condições, podemos separar as contribuições das diferentes multipolaridades através da distribuição angular $W(\theta, \phi)$. O ângulo θ é medido entre a direção do feixe incidente e a do fóton espalhado, enquanto ϕ é o ângulo entre os planos de espalhamento e de polarização do fóton incidente, conforme ilustrado na Fig. 2.

Figura 2 — Definição dos ângulos de espalhamento, θ , e azimutal, ϕ , relevantes nas experiências de espalhamento de fótons polarizados.

A assimetria azimutal na distribuição dos fótons espalhados pode ser usada para determinar a paridade da transição. A distribuição angular dos fótons espalhados é escrita, nesse caso:

$$W(\theta, \phi) = W(\theta) + (\pm)_{\lambda'} \frac{\cos 2\phi}{(1+\delta_1^2)(1+\delta_2^2)} \sum_{\nu} K_{\nu}(1) A_{\nu}(2) P_{\nu}^{(2)}(\cos \theta), \quad (4)$$

onde $(\pm)_{\lambda'}$ é igual a +1 ou -1 conforme as transições sejam respectivamente elétricas ou magnéticas, δ_1 e δ_2 correspondem às razões de mistura multipolar nos canais de entrada e

saída, respectivamente e $W(\theta)$ é a distribuição angular para o espalhamento de fótons não polarizados. Valores numéricos para os coeficientes de correlação angular A_v e K_v e para a função de Legendre associada $P_v^{(2)}(\cos\theta)$ podem ser encontrados, junto a uma descrição detalhada dos aspectos teóricos e práticos da medida da polarização de gamas espalhados pelo núcleo atômico, na Ref. [1].

Costuma-se, nesses casos, trabalhar com a função de assimetria, definida como

$$\Sigma(\theta) = \frac{\sigma_{\perp}(\theta) - \sigma_{\parallel}(\theta)}{\sigma_{\perp}(\theta) + \sigma_{\parallel}(\theta)}, \quad (5)$$

onde $\sigma_{\perp}(\theta) = \sigma(\theta, \phi=90^\circ)$ e $\sigma_{\parallel}(\theta) = \sigma(\theta, \phi=0^\circ)$. É importante notar que a função de assimetria medida depende do grau de polarização linear do feixe P_γ , que é sempre menor que 1: $\Sigma^{\text{exp}}(\theta) = \frac{N_{\perp} - N_{\parallel}}{N_{\perp} + N_{\parallel}} = P_\gamma \Sigma(\theta)$, sendo que $N_{\perp}$ e $N_{\parallel}$ correspondem às taxas de contagens obtidas nas direções perpendicular e paralela ao plano de polarização do feixe incidente, respectivamente.

No Cap. 5, quando discutiremos as perspectivas experimentais no LAL com os novos equipamentos em instalação, voltaremos a esse assunto.

2.3. Eletrodesintegração nuclear

No espalhamento de elétrons obtém-se informação sobre os parâmetros do elétron espalhado, cujo espectro é medido a determinado ângulo, integrando-se sobre todos os estados finais do núcleo espalhador. O espalhamento inelástico de um elétron por um núcleo pode ser interpretado como a troca de um fóton virtual de energia E e momento $q = \hbar\kappa$, conforme esquematizado na Fig. 3. A seção de choque fotonuclear corresponde ao limite de baixo momento transferido, $\kappa \rightarrow k$, da seção de choque de espalhamento de elétrons, onde k é o número de onda do fóton real de energia E . Na absorção de fótons reais apenas os multipolos de ordem mais baixa são envolvidos, enquanto que no espalhamento de elétrons multipolos mais altos participam, tornando a seção de choque

de espalhamento mais complexa com o aumento do momento transferido.

Figura 3 — Cinemática do processo de espalhamento inelástico de elétrons na aproximação PWBA. $|i\rangle$ e $|f\rangle$ representam os estados inicial e final do núcleo espalhador, respectivamente. E e q indicam a energia e o momento transferido pelo fóton virtual, e k_1 e k_2 os números de onda dos elétrons incidente e espalhado, respectivamente.

A energia do fóton virtual corresponde à energia de excitação transferida ao núcleo e é determinada pela perda de energia do elétron espalhado. O momento transferido ao núcleo, q , para uma mesma energia de excitação, E , pode ser variado mudando-se o ângulo de espalhamento e/ou a energia do elétron incidente, de acordo com:

$$\kappa = (k_1^2 + k_2^2 - 2k_1 k_2 \cos\theta)^{1/2} \quad (6)$$

$$E = E_1 - E_2,$$

onde E_1 , k_1 e E_2 , k_2 são as energias e os números de onda dos elétrons incidente e emergente, respectivamente.

Apesar de experimentos de espalhamento de elétrons, particularmente em núcleos de número atômico elevado, necessitarem de análise na aproximação de Born com ondas distorcidas (DWBA), na discussão a seguir as características relevantes do processo de espalhamento de elétrons são resumidas numa descrição feita na aproximação de Born com ondas planas (PWBA), uma vez que, nesse tratamento, os processos físicos subjacentes aparecem com mais nitidez, devido à sua maior simplicidade.

Em PWBA, a seção de choque de espalhamento inelástico de elétrons pode ser escrita como a soma de seções de choque para a excitação de diferentes multipolaridades elétricas e magnéticas [2-5]:

$$\frac{d^2\sigma}{d\Omega} = \sum_{\lambda=0}^{\infty} \frac{d^2\sigma^{C\lambda}}{d\Omega} + \sum_{\lambda=1}^{\infty} \left(\frac{d^2\sigma^{E\lambda}}{d\Omega} + \frac{d^2\sigma^{M\lambda}}{d\Omega} \right), \text{ sendo} \quad (7)$$

$$\frac{d^2\sigma^{C\lambda}}{d\Omega} = \left(\frac{e}{\hbar c} \right)^2 \frac{4\pi(\lambda+1)}{\lambda[(2\lambda+1)!!]^2} \frac{\kappa^{2\lambda}}{k_1^2} \frac{\lambda}{\lambda+1} V_C(\theta) B(C\lambda, \kappa); \quad (8)$$

$$\frac{d^2\sigma^{E\lambda}}{d\Omega} = \left(\frac{e}{\hbar c}\right)^2 \frac{4\pi(\lambda+1)}{\lambda[(2\lambda+1)!!]^2} \frac{\kappa^{2\lambda}}{k_1^2} \frac{\lambda}{\lambda+1} V_T(\theta) B(E\lambda, \kappa); \quad (9)$$

$$\frac{d^2\sigma^{M\lambda}}{d\Omega} = \left(\frac{e}{\hbar c}\right)^2 \frac{4\pi(\lambda+1)}{\lambda[(2\lambda+1)!!]^2} \frac{\kappa^{2\lambda}}{k_1^2} \frac{\lambda}{\lambda+1} V_T(\theta) B(M\lambda, \kappa). \quad (10)$$

V_C e V_T são fatores cinemáticos adimensionais, dados por:

$$V_C(\theta) = k_1 k_2 \frac{2k_1^2 + 2k_2^2 + 4m_0^2 c^4 / \hbar^2 - \varepsilon^2 - \kappa^2}{\kappa^4}; \quad (11)$$

$$V_T(\theta) = k_1 k_2 \frac{(k_1^2 + k_2^2 - \varepsilon^2)\kappa^2 - 2(\vec{k}_1 \cdot \vec{\kappa})(\vec{k}_2 \cdot \vec{\kappa})}{\kappa^2(\kappa^2 - \varepsilon^2)}, \quad (12)$$

onde m_0 é a massa de repouso do elétron e $\varepsilon = E/\hbar c = (E_1 - E_2)/\hbar c$. Desprezando o recuo do núcleo, $\varepsilon = k$, o número de onda do fóton que produz a mesma excitação E . A probabilidade de transição reduzida coulombiana, $B(C\lambda, \kappa)$, está associada ao operador de densidade de carga nuclear, ρ_n , enquanto que as probabilidades de transição reduzida transversais, $B(E\lambda, \kappa)$ e $B(M\lambda, \kappa)$, estão associadas aos operadores densidade de corrente, $\vec{j}_n$, e de magnetização, $\vec{\mu}_n$, nucleares. Nas experiências de espalhamento de elétrons os detalhes da estrutura nuclear são explorados através da dependência com o momento transferido dos operadores de transição. Neste ponto é importante notar que, no limite de baixo momento transferido ao núcleo, $\kappa \rightarrow k$, os operadores de transição são idênticos para excitações induzidas por elétrons e fótons, viabilizando o conceito de fótons virtuais [6-8] e permitindo a expressão da seção de choque induzida por elétrons em termos da seção de choque da reação correspondente induzida por fótons.

O espectro de fótons virtuais, para uma multipolaridade $X\lambda$, é definido pela razão:

$$\frac{d^2 N^{X\lambda}}{d\Omega} = \frac{E \frac{d^2\sigma^{X\lambda}}{d\Omega}}{\int \sigma_\gamma^{X\lambda}(E) dE}, \quad (13)$$

onde a integral no denominador é dada pela eq. (3), e X é E ou M , conforme a transição tenha caráter elétrico ou magnético, respectivamente. É importante lembrar que elétrons apresentam probabilidades de transição reduzidas coulombianas (ou longitudinais), que

não têm análogo para fótons. Assim, além do limite de baixos momentos transferidos, $\kappa \rightarrow k$, uma restrição adicional tem que ser imposta, o limite de comprimento de onda longo, $kr \ll 1$, no qual $B(C\lambda, E) = B(E\lambda, E)$. Essa relação pode ser demonstrada através do uso do teorema de Siegert [4]. O espectro de fótons virtuais, calculado em PWBA, entretanto, pode ser usado apenas na análise de experiências envolvendo núcleos muito leves, sendo necessário, na maioria das situações, considerar a distorção das funções de onda dos elétrons incidente e emergente no campo coulombiano do núcleo. O espectro de fótons virtuais foi calculado em DWBA, levando em consideração o tamanho do núcleo, por Onley e colaboradores [9-12].

Em uma experiência de eletrodesintegração, mede-se os fragmentos nucleares emitidos ou a radioatividade induzida por uma certa reação, integrando-se sobre todos os estados finais do elétron espalhado. Isso corresponde a integrar a seção de choque de espalhamento inelástico de elétrons sobre todos os ângulos de espalhamento e sobre a faixa de energia de excitação que produz a reação em estudo. Assim, cada ponto medido recebe contribuições de uma distribuição de momentos transferidos ao núcleo. A seção de choque de eletrodesintegração inclui todos os multipolos excitados no espalhamento de elétrons, dando ênfase, no entanto, às componentes transversais, uma vez que ela é dominada pelos ângulos dianteiros, não medidos em espalhamento de elétrons.

A seção de choque de eletrodesintegração pelo canal x , $\sigma_{e,x}(E_1)$, pode então ser relacionada com a correspondente seção de choque de fotodesintegração, $\sigma_{\gamma,x}(E)$, através do espectro de fótons virtuais:

$$\sigma_{e,x}(E_1) = \int_{\text{limiar}}^{E_1 - m_0 c^2} \sum_{X\lambda} \sigma_{\gamma,x}^{X\lambda}(E) N^{X\lambda}(E_1, E, Z) E^{-1} dE. \quad (14)$$

Como fótons virtuais transferem mais momento ao núcleo do que fótons reais, há a possibilidade de excitação de multipolos mais elevados e isso se reflete na forma dos espectros de fótons virtuais, como pode ser visto na Fig. 4, que compara espectros $E1$ e $E2$ calculados em PWBA e DWBA. Nota-se que ambos os espectros DWBA são intensificados em relação aos resultados em PWBA, mas o efeito é mais pronunciado para $E2$.

Figura 4 — Espectros de fótons virtuais calculados em PWBA e DWBA, para um elétron de 9,5 MeV espalhado por um núcleo de Urânio.

É interessante, neste ponto, chamar a atenção para a similaridade entre a expressão da seção de choque de eletrodesintegração, eq. (14), e a correspondente taxa de produção (*yield*) de fotodesintegração, obtida quando um feixe de fótons de bremsstrahlung incide sobre um alvo:

$$Y_{br,x}(E_1) = N_r \int_{\text{limiar}}^{E_1 - m_0 c^2} \sigma_{\gamma,x}(E) N_{br}(E_1, E, Z_r) E^{-1} dE, \quad (15)$$

sendo E_1 a energia do feixe de elétrons que produz o bremsstrahlung, N_r e Z_r o número de átomos por cm^2 e o número atômico do radiador, respectivamente, e N_{br} o espectro de bremsstrahlung. A seção de choque de fotodesintegração, $\sigma_{\gamma,x}(E) = \sum_{\lambda} \sigma_{\gamma,x}^{\lambda}(E)$, uma vez

que o espectro de bremsstrahlung apresenta todos os multipolos na mesma proporção. Dessa forma um experimentador, usando feixes de elétrons e de bremsstrahlung, pode alterar a composição multipolar dos fótons que incidem sobre o alvo. Essa característica foi bastante explorada no LAL e em outros laboratórios, para estudar as componentes multipolares da seção de choque fotonuclear.

O método utilizado, na maioria dos casos (inclusive no LAL), por uma questão de

simplicidade experimental, é o esquematizado na Fig. 5, em que são alternadas medidas de eletro e eletro+fotodesintegração.

Figura 5 — Descrição esquemática dos experimentos de eletro e eletro+fotodesintegração. E_1 é a energia dos elétrons incidentes e x representa a partícula ou radiação detectada para identificar a reação. A distância radiador-alvo, na parte inferior da figura, foi aumentada para facilitar a visualização.

No primeiro caso o alvo é colocado diretamente no feixe de elétrons. No segundo, um radiador fino é colocado imediatamente antes do alvo, para produzir os fótons de bremsstrahlung. O feixe de elétrons, que passa praticamente ileso pelo radiador (tipicamente 2 a 3 % dos elétrons são espalhados no radiador), incide também no alvo produzindo eletrodesintegração, além da fotodesintegração provocada pelo feixe de bremsstrahlung. Nessas condições as taxas obtidas são aproximadamente o triplo daquelas devidas apenas à eletrodesintegração, permitindo boa separação dos efeitos. Apesar de algumas correções necessárias, esse método evita a necessidade da colocação de um eletroímã entre o radiador e o alvo para desviar o feixe de elétrons. A proximidade entre o radiador e o alvo garante também que todo o feixe de bremsstrahlung incide sobre o alvo.

Grande parte de meu trabalho de pesquisa baseou-se nesse método, utilizando-o como ferramenta para o estudo dos canais de decaimento das ressonâncias gigantes, da

composição multipolar de seções de choque fotonucleares, ou ainda, na fase em que muitos laboratórios o utilizavam, como objeto de estudo, refinando-o ou avaliando seus limites de aplicação.

O Anexo 1 (7.1.2.19), relata medidas de seções de choque das reações (e,p) e (e,α) nos isótopos ^{90}Zr e ^{92}Zr , realizadas durante meu estágio de pós-doc no NIST (então NBS, em Gaithersburg, Maryland, EUA). Nesse experimento, medindo as partículas emitidas com um espectrômetro magnético, pudemos determinar as componentes $E1$ e $E2$ das seções de choque (γ,p) e (γ,α) nesses dois isótopos do Zr. A experiência apresentou resultados interessantes em vários aspectos. O primeiro deles por confirmar resultado anterior obtido por Tamae e colaboradores [13], no qual a seção de choque fotonuclear de ambos os canais estudados apresentava crescimento quase monotônico acima do pico da ressonância gigante de dipolo elétrico (RGE1). Esse mesmo efeito já havia sido reportado em trabalho anterior sobre a seção de choque da reação $^{65}\text{Cu}(e,\alpha)$ [14], não confirmado quando medimos a mesma seção de choque por atividade residual (um dos assuntos de minha tese de doutorado), conforme artigo no Anexo 2 (7.1.2.22). Uma das explicações para o estranho comportamento observado seria a emissão múltipla de partículas, uma vez que os sistemas de detecção usados por Tamae e colaboradores e por nós no NIST não conseguiam separar a multiplicidade dos eventos observados. Essa questão ainda está em aberto, pois medidas da multiplicidade de emissão de partículas carregadas em outros núclídeos (veja discussão abaixo), apresentaram resultados incompatíveis com essa hipótese.

No experimento do Zr medíamos o espectro das partículas emitidas, e assim tivemos oportunidade de determinar a seção de choque do nível análogo isobárico I^- de 16,28 MeV do ^{90}Zr . Esse nível é bastante estreito, e a medida de sua seção de choque em função da energia do elétron, para fótons virtuais (eletrodesintegração) e fótons reais (fotodesintegração), permitiu a realização de um teste experimental do espectro de fótons virtuais de $E1$, para fótons de 16,28 MeV, mostrando que os espectros calculados em DWBA, corrigidos para o tamanho finito do núcleo, apresentavam bons resultados para elétrons na faixa de 20 a 100 MeV. Veja o Anexo 3 (7.1.1.3).

Um trabalho semelhante a este último, realizado um pouco antes, na linha de teste da validade dos espectros de fótons virtuais, foi feito usando resultados obtidos no LAL para a seção de choque $^{181}\text{Ta}(e,n)$, Anexo 4 (7.1.2.18). Fizemos, nesse caso, um estudo comparativo entre as várias correções de tamanho do núcleo disponíveis para os espectros de fótons virtuais calculados em DWBA. Mostramos que os espectros de fótons virtuais $E1$ e $E2$, calculados em DWBA considerando o núcleo finito por Onley e colaboradores apresentavam os melhores resultados. Mostramos também que essa correção não apresentava bons resultados no caso do espectro $M1$. Além disso determinamos a componente $E2$ da seção de choque (γ,n) do ^{181}Ta . Concluimos que a análise de experimentos de eletrodesintegração através do método dos fótons virtuais, que geralmente envolve uma seção de choque fotonuclear determinada em um experimento independente, necessita, para a obtenção de resultados confiáveis, da medida da taxa de produção de fotodesintegração, de forma a permitir que diferenças em valor absoluto possam ser consideradas.

O trabalho mostrado no Anexo 5 (7.1.2.16) relata os resultados obtidos na análise, pelo método dos fótons virtuais, da seção de choque de eletrodesintegração e da taxa de produção de eletro+fotodesintegração da reação $^{197}\text{Au}(e,n)$. Obtivemos resultados para as componentes $E2$ e $M1$ da seção de choque (γ,n) do ^{197}Au , mostrando ser possível determinar a intensidade $M1$ acima do limiar de emissão de partículas.

No Anexo 6 (7.1.1.4) utilizamos a seção de choque da reação $^{63}\text{Cu}(\gamma,2n)$ para investigar a ressonância $E2$ isovetorial, que tem pico em torno de 30 MeV. Mostramos que o canal $2n$ é responsável por 65 % do decaimento da RGE2 isovetorial no ^{63}Cu .

O trabalho mostrado no Anexo 7 (7.1.2.21) e o 7.1.1.5 (não anexado aqui por ser muito extenso, veja o anexo do Memorial), descrevem um estudo sistemático das seções de choque fotonucleares medidas com fótons monocromáticos nos laboratórios de Saclay e Livermore, responsáveis pela maior parte dos dados de seções de choque de reações (γ,xn) . Nossas medidas, feitas por atividade residual, que identificam os canais de decaimento sem ambigüidades, permitiram-nos esclarecer uma discrepância existente entre os dados dos dois laboratórios, que persistia na literatura havia vários anos.

Mostramos que ambos os laboratórios mediam o mesmo número total de nêutrons, divergindo na separação de multiplicidades. Mostramos ainda que Livermore fazia a separação de multiplicidades correta, enquanto que Saclay apresentava melhores resultados em valor absoluto.

O trabalho mostrado no Anexo 8 (7.1.2.17) relata medidas da seção de choque de emissão de nêutrons ($\sigma_{e,Tn} = \sigma_{e,n} + 2\sigma_{e,2n}$) nos núclídeos ^{181}Ta , ^{208}Pb e ^{209}Bi . A medida foi realizada detectando-se os nêutrons emitidos, o que impediu a medida da taxa de produção de fotodesintegração, uma vez que não era possível a colocação de um radiador antes do alvo, devido à produção de nêutrons no próprio radiador. Os resultados foram analisados pelo método dos fótons virtuais de forma a determinar a intensidade das componentes $E2$ (isoescalar e isovetorial) e $M1$, presentes na seção de choque fotonuclear.

O Anexo 9 (7.1.2.15) apresenta resultados para as seções de choque das reações (γ,p) e $(\gamma,2p)$ nos isótopos 29 e 30 do Si, obtidos através da análise de medidas das seções de choque de eletrodesintegração e eletro+fotodesintegração correspondentes. Como a medida foi feita por atividade residual, a identificação dos estados finais não apresenta ambigüidades. Procuramos, nessa medida, determinar se a seção de choque $(\gamma,2p)$ teria magnitude suficiente para explicar o comportamento observado nos resultados obtidos nos isótopos do Zr (já comentado). Os resultados obtidos indicaram que a multiplicidade não poderia explicar os resultados do Zr, uma vez que o canal $2p$ é muito menor que o canal p ($\sim 2\%$ no ^{29}Si e $\sim 0,5\%$ no ^{30}Si). Mostramos, no entanto, que a seção de choque (γ,p) no ^{29}Si é 47 % maior que a (γ,n) . No ^{30}Si elas são da mesma ordem. Mostramos também que o canal $1p$ é predominantemente $E1$ em ambos os isótopos, mas que o canal $2p$ tem uma contribuição $E2$ da mesma ordem da $E1$.

O Anexo 10 (7.1.2.14), é um aprimoramento de meu trabalho de mestrado. Medimos as seções de choque (e,n) e $(e+\gamma,n)$ no ^{238}U . A faixa de energia da medida foi estendida a 60 MeV, complementando os dados obtidos anteriormente até 28 MeV. Foram utilizados na análise espectros de fótons virtuais $E1$ e $E2$ calculados em DWBA para núcleo finito. Os resultados obtidos mostram que a componente $E1$ do canal $1n$

exaure 40 % da respectiva regra de soma ponderada na energia (RSE1), enquanto que a componente $E2$ exaure entre 70 e 80 % da RSE2.

O trabalho mostrado no Anexo 11 (7.1.2.13) dá continuidade a uma das medidas de meu trabalho de doutorado. Foram medidas as seções de choque de eletrodesintegração e as taxas de produção de eletro+fotodesintegração das reações $(e,2n)$ e $(e,2p)$ no ^{63}Cu . As medidas foram estendidas até 60 MeV e os resultados anteriores incluídos na análise. Foram derivadas as seções de choque $(\gamma,2n)$ e $(\gamma,2p)$ e separadas suas componentes $E1$ e $E2$. Com essa medida mostramos que, também no ^{63}Cu , a seção de choque $(\gamma,2p)$ não apresenta intensidade suficiente (esgota apenas 1 % da RSE1) para explicar o comportamento observado na medida do Zr como sendo devido à multiplicidade de emissão das partículas.

Em resumo, mostramos que, através do uso do método dos fótons virtuais, podemos obter informações detalhadas sobre a absorção da radiação eletromagnética de multipolaridades $E1$ e $E2$ pelo núcleo atômico. O método, por envolver um processo de deconvolução, não é muito sensível a detalhes como forma e posição da seção de choque fotonuclear, mas é bastante acurado em relação à intensidade, obtendo resultados muito precisos para as seções de choque integradas. Estudamos, com esse método o decaimento das ressonâncias de $E1$ e $E2$ através de vários canais de emissão de partículas neutras e carregadas. Investigamos, em particular, a multiplicidade de emissão de partículas carregadas em vários núclídeos.

Exploramos a possível influência de excitações $E0$ e $E3$ no processo de análise e mostramos que ambas têm contribuição desprezível para a eletrodesintegração. No caso de $E0$, o fato das energias envolvidas estarem restritas à região em que a componente longitudinal dos elementos de matriz de transição são muito pequenos, limita a excitação desse modo. No caso de $E3$, a seção de choque de fotoabsorção é tão pequena que a ênfase do espectro $E3$ em relação aos de $E1$ e $E2$ não é suficiente para tornar as medidas de eletroprodução sensíveis a esse modo de excitação.

Estudamos os limites de validade do método, mostrando que a aproximação de núcleo pontual falha a partir de ~50 MeV para núcleos médios e mesmo a partir ~15 MeV

para pesados, no caso de espectros $E2$. Para espectros $E1$ esses limites não são tão críticos, havendo limitações para núcleos pesados a partir de ~ 50 MeV. Mostramos que esta falha está associada à importância relativa da componente longitudinal dos elementos de matriz de transição no espalhamento de elétrons. No caso de transições $E1$, a componente longitudinal é menor que a transversal até energias mais elevadas, enquanto que para $E2$ elas se tornam da mesma ordem já em torno de algumas dezenas de MeV.

3. Monocromação de fótons de bremsstrahlung

3.1. Introdução

A produção de feixes de fótons monocromáticos com energia suficiente para o estudo de propriedades nucleares tem sido, há décadas, um desafio para os físicos que trabalham nessa área. Diversos métodos foram desenvolvidos, como a aniquilação de pósitrons em voo, captura de nêutrons e marcação de fótons (*photon tagging*), cada qual com seus prós e contras.

Neste capítulo apresento uma discussão sobre um dos métodos de produção de fótons monocromáticos, classificado como marcação de fótons, que é a monocromação de fótons de bremsstrahlung. Existem essencialmente três tipos de marcação de fótons:

- aquela que envolve marcar fótons produzidos pela aniquilação em voo de pósitrons de alta energia, em que a energia e o tempo de um dos fótons (geralmente o fóton de energia mais alta, emitido na direção frontal no centro de massa) podem ser determinados através da medida da energia e do ângulo de emissão do outro fóton, em coincidência com o produto da reação;
- outra possibilidade é a que envolve o retroespalhamento Compton de fótons de um feixe de laser por elétrons de alta energia. A medida da energia do elétron que produziu o espalhamento em coincidência com o produto da reação permite a determinação da energia do fóton produzido;
- o terceiro processo trata da monocromação de fótons de bremsstrahlung, através da medida da energia do elétron espalhado, em coincidência com o sinal do detector do produto da reação. Esse método é o mais versátil e mais bem sucedido entre os descritos.

No item 3.2 apresento uma descrição detalhada do princípio de funcionamento do método de monocromação de fótons de bremsstrahlung, discutindo, nos itens subsequentes, suas características mais relevantes aos procedimentos experimentais de interesse.

3.2. Monocromação de fótons de bremsstrahlung

O sistema usado para a marcação de fótons de bremsstrahlung é conceitualmente simples e está esquematizado na Fig. 6. A figura mostra um feixe de elétrons monoenergético, de energia cinética E_1 , que incide sobre um radiador fino (tipicamente $\sim 10^{-3}$ comprimentos de radiação), produzindo um feixe de fótons de bremsstrahlung na direção frontal. Um espectrômetro magnético transmite os elétrons espalhados de energia E_2 para um dos detectores de seu plano focal. Um sinal desse detector denuncia a presença de um fóton de energia $k = E_1 - E_2$; a coincidência temporal entre esse sinal e o sinal do detector do produto de reação indica que a reação foi produzida pelo fóton associado. O espectrômetro magnético tem também uma segunda função, que é a de separar o feixe de elétrons primário (que passa praticamente ileso pelo radiador) do feixe de fótons.

Figura 6 — Representação esquemática do sistema de marcação de fótons. E_1 é a energia do elétron incidente, E_2 a energia do elétron espalhado e k a energia do fóton produzido.

O espectro de energia dos fótons, k , é contínuo no intervalo $]0, E_1[$, devido ao processo de bremsstrahlung, como ilustrado na Fig. 7. Somente os elétrons residuais com momentos contidos no intervalo de aceitação do espectrômetro podem ser registrados nos detectores do plano focal. Dessa forma, para uma dada escolha do campo magnético do espectrômetro, apenas uma certa faixa do espectro de bremsstrahlung é mapeada, sendo

que sua largura depende da relação entre k e E_1 , conforme ilustrado na Fig. 7 pelas regiões a e b . A faixa de operação tradicional de espectrômetros que operam como marcadores de fótons compreende a região $\sim 0,1E_1 \leq k \leq \sim 0,9E_1$. As limitações se devem, no extremo de altas energias do fóton, ao fato de os elétrons residuais terem energia muito baixa, confundindo-se com elétrons secundários produzidos no radiador e nas paredes do próprio espectrômetro. No outro extremo, em que $k \sim 0$, temos que a energia do elétron residual é próxima da do feixe primário, $E_2 \sim E_1$, inviabilizando sua separação e, portanto, sua detecção.

Figura 7 — Espectro de bremsstrahlung em função das energias k (do fóton, na abscissa) e E_2 (do elétron residual, na margem superior). As regiões a e b representam os intervalos marcados, definidos pela aceitação em momento do espectrômetro e pelo momento central do intervalo. E_1 é a energia do elétron incidente.

3.3. Características do sistema de marcação de fótons de bremsstrahlung

As principais características de um sistema de marcação de fótons, relevantes às experiências a serem realizadas, são a intensidade, faixa de energia de operação e resolução em energia do feixe de fótons e a eficiência de marcação. Outro parâmetro que pode ser relevante é o grau de polarização linear do feixe, caso essa potencialidade esteja sendo explorada. Discuto a seguir cada uma dessas características e seu impacto nas

possibilidades experimentais.

3.3.1. Intensidade

Em função de que no processo de marcação de fótons é necessária a detecção de uma partícula, correlacionada em tempo, para cada fóton utilizável do feixe, há um limite superior para o fluxo total de fótons, devido às taxas de contagens finitas que os detectores dos ramos de elétrons e de fótons (ou produtos de reação) podem aceitar. Com frequência, no entanto, o fator mais restritivo ao fluxo total de fótons costuma ser a taxa de coincidências acidentais entre o ramo de elétrons e o dos produtos de reação. Alguns fatores que contribuem para o aumento das coincidências acidentais são: a eficiência dos detectores, que é sempre menor que 100 %; a aceitação angular; e a aceitação em momento do espectrômetro. A aceitação angular está relacionada com a colimação tanto do feixe de fótons quanto, eventualmente, do de elétrons, se sua abertura for maior que a da janela do espectrômetro. Toda colimação está associada a perda de eventos e consequentemente aumento na taxa de acidentais. A exclusão dos elétrons com momentos fora da faixa de aceitação do espectrômetro resulta em fótons não correlacionados no sistema.

Podemos fazer uma estimativa da taxa de coincidências acidentais, calculando, inicialmente, a taxa de contagens obtida no detector de produtos da reação, devido aos fótons marcados:

$$N_x^{tag} = \frac{I}{e} \varepsilon_D N_r N_{alv} \int_{k_1}^{k_2} \sigma_{\gamma,x}(k) N_{br}(E_1, k, Z_r) k^{-1} dk \quad (16)$$

onde I é a corrente do feixe de elétrons; e é a carga elementar; ε_D é um fator de eficiência que leva em consideração todos os fatores geométricos e de eficiência de detecção da partícula x ; $\sigma_{\gamma,x}(k)$ é a seção de choque para produção da partícula x , por um fóton de energia k ; N_{br} é o espectro de bremsstrahlung para um radiador de número atômico Z_r ; e N_r e N_{alv} correspondem, respectivamente, ao número de átomos por cm^2 no radiador e no alvo. O intervalo de integração, $[k_1, k_2]$, corresponde à região do espectro de bremsstrahlung que está sendo marcada (aceitação em momento do espectrômetro). De forma análoga, a taxa de contagens devido aos fótons não marcados pode ser calculada

por:

$$N_x^{untag} = \frac{I}{e} \varepsilon_D N_r N_{div} \left[\int_0^{k_1} \sigma_{\gamma,x}(k) N_{br}(E_1, k, Z_r) k^{-1} dk + \int_{k_2}^{E_1} \sigma_{\gamma,x}(k) N_{br}(E_1, k, Z_r) k^{-1} dk \right] \quad (17)$$

sendo que a integração é feita sobre os intervalos não marcados $[0, k_1[$ e $]k_2, E_1]$.

A taxa de contagens total, N_e , nos detectores de elétrons, corresponde à taxa com que elétrons são espalhados dentro do ângulo sólido da janela do espectrômetro e dentro de seu intervalo de aceitação em momento.

A taxa de contagem correspondente às coincidências reais, no detector i do plano focal, $N_{T,i}$, será igual à taxa no detector de produtos da reação devida aos fótons marcados, já calculada: $N_{T,i} = N_x^{tag}$. Por outro lado, no detector de produtos da reação, a taxa de contagens é devida não só aos fótons marcados, mas também aos não marcados: $N_x = N_x^{tag} + N_x^{untag} = N_x^{tag}(1+f)$, onde o fator f leva em consideração os fótons não marcados dentro e fora do intervalo de marcação, eventos devido a outros canais de reação não distinguíveis dos de interesse e outras possíveis fontes de background. Dessa forma, a taxa de coincidências acidentais entre os dois detectores será aproximadamente $N_{RC} = N_{e,i} N_x \tau D^{-1}$, onde τ é a resolução temporal do sistema e D é o fator de utilização macroscópico do acelerador (*duty cycle*). A razão entre as taxas de coincidências reais e acidentais pode então ser escrita como:

$$R = \frac{N_{T,i}}{N_{RC}} = \frac{N_x^{tag}}{N_x^{tag}(1+f)} \frac{D}{N_{e,i} \tau} = \frac{D}{N_{e,i} \tau (1+f)} \quad (18)$$

Assim, fica claro que se o fator f aumentar, seja pela limitação do ângulo sólido de aceitação do espectrômetro, seja pela redução do intervalo de marcação, a taxa de coincidências acidentais também aumentará. Usando a definição de eficiência de marcação (veja discussão abaixo), podemos escrever a equação acima em função do número de fótons marcados no intervalo i : $N_{\gamma,i} = \varepsilon_i N_{e,i}$. Assim, $R = \frac{\varepsilon_i D}{N_{\gamma,i} \tau (1+f)}$. É

interessante notar que, por definição, $f \geq 0$ e D e ε_i são ambos menores ou iguais a 1. Dessa forma, podemos escrever a equação anterior como uma desigualdade: $R \leq (N_{\gamma,i} \tau)^{-1}$, ficando claro que, mesmo em um experimento ideal, o fluxo fica limitado a uma taxa de

$\sim 1/\tau$ para uma razão reais/acidentais igual a 1. Por causa dessas restrições é fundamental a utilização de detectores de produto de reação que apresentem o maior ângulo sólido possível, para maximizar as taxas de aquisição.

Podemos definir duas eficiências que são importantes em um sistema de marcação de fótons. A primeira delas, que é usualmente chamada de eficiência de marcação, é definida como a razão entre o número de fótons marcados em um certo intervalo de marcação e o número total de contagens no detector do plano focal correspondente. O segundo tipo, a probabilidade de marcação, é definida como a razão entre o número de fótons marcados em um certo intervalo e o número total de fótons nesse intervalo.

Vários fatores contribuem para reduzir tanto a eficiência quanto a probabilidade de marcação. Como já comentado, a eficiência de marcação é reduzida sempre que o feixe de fótons é colimado. Isso pode ser necessário por uma série de razões, desde a geometria do sistema de detecção ou do alvo, até eventualmente a produção de fótons linearmente polarizados (veja discussão no item 3.3.3). Já a colimação no ramo dos elétrons (usualmente provocada pela aceitação angular do espectrômetro) tende a reduzir a probabilidade de marcação, pois aumenta o número de fótons não marcados no alvo. Esse problema é mais crítico quando trabalhamos em baixas energias, uma vez que a divergência angular dos elétrons espalhados tende a aumentar com a diminuição da energia. Da cinemática do processo de bremsstrahlung pode-se definir (desprezando o recuo do núcleo) um ângulo característico para a divergência angular dos elétrons espalhados: $\theta_{ec} = \frac{m_0 k}{E_1 E_2} = \frac{k}{E_2} \theta_{\gamma c}$, onde $\theta_{\gamma c}$ é o ângulo característico de divergência do feixe de bremsstrahlung. Nota-se que baixas energias dos elétrons incidente ou espalhado tendem a aumentar θ_{ec} .

3.3.2. Intervalo de energia e resolução

O intervalo de energia em que os fótons são marcados, $k_{max} - k_{min}$, depende da aceitação em momento do espectrômetro e da escolha do momento central do intervalo. A resolução em energia dos fótons, $\Delta k/k$, está relacionada com as incertezas nas energias do elétron incidente e espalhado. A partir da conservação da energia envolvida no processo,

$k = E_1 - E_2$, podemos escrever:

$$\frac{\Delta k}{k} = \frac{[(\Delta E_1)^2 + (\Delta E_2)^2]^{1/2}}{E_1 - E_2}, \quad (19)$$

onde ΔE_1 e ΔE_2 são, respectivamente, as resoluções em energia dos elétrons incidente e espalhado.

A resolução em energia do feixe incidente depende da dispersão em energia do acelerador e da espessura do radiador (que tem que ser fino para minimizar o straggling). O termo ΔE_2 apresenta contribuições também da perda de energia no radiador e da resolução do espectrômetro e dos detectores do plano focal. Para condições definidas do alvo, do espectrômetro e do feixe incidente, a resolução em energia pode ser melhorada reduzindo-se o valor do momento central analisado (o que diminui a largura do intervalo de energia que está sendo analisado simultaneamente, conforme mostrado na Fig. 7. Outro aspecto favorável de se trabalhar com a escolha do momento central baixo reside no fato de que a energia máxima dos fótons marcados se aproxima da energia limite (E_1). Isso é particularmente importante em situações em que a seção de choque está crescendo com a energia, para minimizar o fator f discutido no item anterior, relacionado com reações produzidas por fótons não marcados. Deve-se, portanto, procurar um compromisso com a probabilidade de marcação, uma vez que a redução de E_2 também introduz dificuldades, como o aumento do ângulo de divergência e do fundo devido a elétrons secundários.

3.3.3. Polarização

Apesar de um feixe de fótons de bremsstrahlung ser, como um todo, não polarizado, feixes de fótons linearmente polarizados podem ser obtidos através da seleção de fótons e elétrons em ângulos sólidos fora do eixo de simetria do cone de bremsstrahlung. Esse processo, conhecido como método de colimação dupla fora do eixo [15], baseia-se no fato de que há uma assimetria na distribuição dos vetores de polarização dos fótons, que é correlacionada com os ângulos dos fótons emitidos e elétrons espalhados (veja discussão no capítulo 4).

Se por um lado o trabalho em baixa energia é restritivo em termos da aceitação angular dos elétrons espalhados, por outro lado esse efeito é benéfico para a seleção de fótons polarizados. Com o aumento da energia, o ângulo característico do cone de bremsstrahlung se reduz, tornando bastante complicado o processo de colimação assimétrica necessário para a obtenção de fótons polarizados. Em baixas energias os ângulos são maiores e a colimação é mais eficiente. No Anexo 12 (7.1.2.11) discutimos em detalhe a viabilidade da produção de fótons polarizados com o tagger do LAL, tanto com o primeiro quanto com o segundo estágios do acelerador (5 e 30 MeV, respectivamente).

3.3.4. O tagger do LAL

O espectrômetro de elétrons que operará como marcador de fótons no LAL foi recebido como doação do MAX-Lab de Lund, Suécia, onde operou como tagger [16] por vários anos. A Fig. 8 mostra um desenho esquemático do espectrômetro do LAL. A doação incluiu, além do espectrômetro e seu suporte mecânico, as câmaras de vácuo e suportes para os radiadores (controlados por motores de passo), mas infelizmente não incluiu a parte eletrônica do sistema de detecção do plano focal (apenas os cintiladores e guias de luz).

Figura 8 — Tagger do LAL (doado pelo MAX-Lab, de Lund, Suécia). a) Peças polares; b) bobinas; c) câmara de vácuo que acomoda o sistema de detecção do plano focal; e d) câmara de vácuo que acomoda o suporte do radiador (as câmaras são interligadas sem janelas).

Assim, foi necessária uma adaptação para a instalação das novas fotomultiplicadoras (doadas pelo KIPT, de Kharkov, Ucrânia) que são um pouco maiores que as originais. Nesse processo aproveitamos para polir e refazer as junções entre os cintiladores e os guias de luz, que apresentavam desgaste e opacidade. O trabalho foi feito com sucesso, o plano focal foi remontado com suas novas fotomultiplicadoras, novos suportes mecânicos e novas blindagens magnéticas (além de novo sistema de refrigeração) e calibrado usando fontes radioativas emissoras de elétrons. A Fig. 9 mostra a curva de calibração obtida. Esse foi o trabalho de mestrado do Sr. Reinaldo Guarino.

Figura 9 — Curva de calibração da corrente I em função do canal do plano focal para elétrons de 0,625 MeV (elétrons de conversão de uma fonte de ^{137}Cs).

4. A seção de choque de bremsstrahlung

4.1. Introdução

Como vimos no capítulo 2, a medida de taxas de produção de eletrodesintegração e de eletro+fotodesintegração, permite que se obtenham informações sobre as seções de choque fotonucleares correspondentes. O método, no entanto, depende fortemente da qualidade dos espectros de fótons virtuais e de bremsstrahlung utilizados no processo de análise. No item 4.2 apresento o primeiro trabalho que realizei nessa área, ainda durante meu período de pós-doc no NIST, que descreve um teste experimental da seção de choque de bremsstrahlung, realizado para identificar qual das várias expressões existentes seria a mais adequada para a análise dos experimentos.

O conhecimento da seção de choque diferencial de bremsstrahlung é fundamental para que se possa projetar e avaliar o desempenho de equipamentos como os descritos no capítulo 3, bem como das experiências que os utilizam. Isso se deve ao fato de que é preciso conhecer a intensidade e a distribuição angular dos fótons emitidos no processo de bremsstrahlung, assim como a dos elétrons espalhados, porque trabalha-se com janelas de aceitação (colimadores e/ou alvos ou detectores) de dimensões finitas, que selecionam uma fração do número total de partículas envolvidas no processo. A integração da seção de choque diferencial sobre as janelas de aceitação de fótons e elétrons permite que se estime as taxas de contagens em cada geometria. Além disso, como os eventos são selecionados também por um processo de coincidência temporal, a seleção geométrica tende a aumentar a taxa de coincidências acidentais e o conhecimento desse fator é fundamental para a avaliação do desempenho de um experimento. Na seção 4.3 apresento a seção de choque quádrupla diferencial de bremsstrahlung, $\frac{d^4\sigma}{d\Omega_k d\Omega_{p_1} dk}$, calculada sem as usuais aproximações de altas energias e pequenos ângulos e descrevo o programa de simulação por Monte Carlo, que nos permitiu fazer um estudo detalhado do

comportamento dessa seção de choque. Apresento também os principais resultados obtidos.

4.2. Teste Experimental da Seção de Choque de Bremsstrahlung

A análise conjunta de medidas de seção de choque de eletrodesintegração e de taxas de produção de eletro+fotodesintegração, para obter informações sobre as correspondentes seções de choque fotonucleares, depende fortemente dos espectros de fótons virtuais e de bremsstrahlung utilizados. Espectros de fótons virtuais $E1$ e $E2$, calculados em DWBA (aproximação de Born com ondas distorcidas) e corrigidos para o tamanho finito do núcleo apresentaram excelente acordo com experimentos. No caso dos espectros de bremsstrahlung, apesar da extensa literatura a respeito, testes específicos para o tipo de uso em questão (convolução com seções de choque fotonucleares) não existiam sobre toda a faixa de energia necessária.

Neste trabalho foi utilizada a ativação do ^{63}Cu pelas reações (e,n) ou (γ,n) como detector da intensidade de fótons de bremsstrahlung incidentes na amostra (seção de choque integrada sobre todos os ângulos de emissão de fótons e na energia dos fótons, ponderada na seção de choque de fotodesintegração do cobre). Utilizamos radiadores de Cu ($Z=29$), Mo ($Z=42$), Ta ($Z=73$) e Th ($Z=90$), de forma a mapear efeitos relacionados com a distorção coulombiana nos elétrons incidente e espalhado sobre os espectros de bremsstrahlung. Os resultados obtidos para a seção de choque (γ,n) , integrada até 30 MeV, permitiram uma avaliação adequada dos espectros analisados. Mostramos que o espectro interpolado de Seltzer ajusta bem os dados, que a expressão de Schiff é inaceitável, a menos de uma pequena região próxima ao *tip* do espectro; e que a expressão de Bethe e Heitler em aproximação ultra-relativística, com a correção coulombiana de Davies, Bethe e Maximon é adequada para energias do elétron acima de 30 MeV. O trabalho consta do Anexo 13 (7.1.2.20).

4.3. A Seção de choque diferencial de bremsstrahlung

Tratamos aqui de uma situação bastante diferente da tratada no item anterior, uma vez que lá estávamos lidando com a seção de choque integrada e, agora, necessitamos de sua forma diferencial. Como visto no capítulo 3, no tagger nem todos os elétrons espalhados no radiador são detectados no plano focal e, também, nem todos os fótons produzidos são detectados ou geram produtos que o sejam. Nesse sentido é importante que possamos calcular, através da integração da seção de choque de bremsstrahlung sobre as janelas do espectrômetro e do colimador de fótons, a eficiência do processo de monocromação.

Existem várias formas de se calcular a seção de choque de bremsstrahlung, quase todas usando uma ou mais aproximações. Vamos aqui tratar da expressão de Bethe e Heitler [17], que foi desenvolvida considerando, para os elétrons incidente e espalhado, funções de onda de partícula livre perturbadas em primeira ordem em Z (aproximação de Born em onda plana – PWBA). A escolha desse caminho não foi acidental, mas devido ao fato de que as expressões calculadas em PWBA apresentam formas analíticas razoavelmente simples, existem para energias relativísticas e não-relativísticas e com ou sem correções para blindagem eletrônica da carga nuclear (*screening*). Em geral, as expressões calculadas em PWBA tornam-se menos confiáveis quando *a)* o número atômico do radiador cresce, *b)* a energia do elétron incidente diminui, e *c)* a energia do fóton se aproxima do limite superior (ou seja, a energia do elétron espalhado se aproxima de zero). A primeira limitação é eliminada através do uso de radiadores de número atômico baixo, como o Be ($Z=3$), por exemplo. A limitação colocada no item *c)* é séria, mas o próprio tagger se encarrega de evitá-la, uma vez que a energia mínima dos elétrons que são analisados pelo espectrômetro corresponde a entre 10 e 20 % da energia cinética inicial. Desta forma o limite de altas energias de fótons fica eliminado da faixa de operação. Em relação ao item *b)*, como devemos trabalhar com elétrons acima de 5 MeV (mesmo no primeiro estágio), tomando-se o cuidado de evitar expressões que utilizam aproximações de altas energias ou pequenos ângulos, a aproximação é boa.

Um dos problemas que os usuários dessas expressões encontram é que elas são desenvolvidas no chamado *sistema de coordenadas do fóton*, um sistema que coloca o eixo z na direção do fóton emitido, conforme ilustrado na Fig. 10. Nesse sistema as expressões para a seção de choque diferencial são mais compactas. No entanto, quando se trabalha com um espectrômetro, como no caso do *tagger*, esse sistema de coordenadas é pouco útil, uma vez que ele não é fixo em relação ao laboratório. Quando se faz a mudança para o sistema do espectrômetro (no qual o eixo z aponta na direção do feixe de elétrons incidente e o plano xy é tangente às faces polares do eletroímã, veja Fig. 10) as expressões para a seção de choque tornam-se muito mais complicadas e difíceis de se trabalhar. É importante lembrar que essas expressões devem passar por um processo de integração quádrupla (sobre a janela de aceitação dos elétrons e do colimador de fótons) para que se possa avaliar a taxa de contagens. Para evitar esse processo, desenvolvemos um programa de simulação pelo método de Monte Carlo, que calcula a seção de choque no sistema do fóton, trazendo as condições de contorno (janelas de aceitação) do sistema do espectrômetro. Para maiores detalhes, veja o Anexo 12 (7.1.2.11).

Figura 10 — a) Sistema de referência do fóton, onde o eixo z é orientado na direção de emissão do fóton e o momento do elétron incidente pertence ao plano xz . b) Sistema de referência do espectrômetro, onde o eixo z coincide com a direção do elétron incidente. O plano xz coincide com o plano médio do espectrômetro e o plano xy é paralelo às bordas de entrada de suas faces polares.

O uso deste programa de simulação facilitou bastante o estudo das propriedades diferenciais da seção de choque de bremsstrahlung, permitindo uma avaliação acurada das potencialidades do *tagger* (discutidas no capítulo 3), bem como de outras características

importantes, que comento a seguir.

Uma componente linearmente polarizada do espectro de bremsstrahlung pode ser selecionada através do processo de colimação dupla fora do eixo de simetria, conforme discutido no Anexo 12 (7.1.2.11) e nas referências citadas. Na região de baixas energias do elétron incidente (5 a 30 MeV) podem ser conseguidas na prática as condições de pequena colimação, nas quais o comportamento da intensidade e do grau de polarização do feixe de fótons é determinado pelas propriedades diferenciais da seção de choque de bremsstrahlung. Estudamos, utilizando nosso programa de simulação, o comportamento da seção de choque de bremsstrahlung, quando integrada sobre a área de aceitação de pequenos colimadores, em função da energia do fóton, E_γ , e da posição dos colimadores dos elétrons, θ_e, ϕ_e e dos fótons $\theta_\gamma, \phi_\gamma$.

O estudo mostra que, nas posições dos colimadores correspondentes ao mínimo do momento transferido ao núcleo espalhador, q , a seção de choque apresenta um máximo pronunciado e estreito (em torno de uma energia definida E_γ) com elevado grau de polarização linear. O máximo pronunciado e estreito na seção de choque leva a um caráter quase-monocromático do espectro de fótons selecionados. Essa propriedade se mantém até a aceitação angular do colimador alcançar aproximadamente metade do ângulo eletrodinâmico crítico, θ_c . A posição do máximo de seção de choque, $E_{\gamma,max}$, depende da razão entre os ângulos θ_γ e θ_e . O grau de polarização linear decresce monotonicamente com o aumento da energia dos fótons, $E_{\gamma,max}$.

A Fig. 11 mostra o efeito estudado. Os losangos representam o quadrado do momento transferido, q^2 , os triângulos o grau de polarização linear, P , e os quadrados a seção de choque, σ , em função da energia do fóton. Esses resultados foram obtidos para elétrons de 30 MeV e colimadores pontuais (4,5 mrad).

Figura 11 - Seção de choque, momento transferido ao quadrado e grau de polarização em função da energia do fóton.

A potencialidade do método é ilustrada pelos resultados mostrados a seguir. Na Fig. 12 mostramos a variação da seção de choque com a energia do fóton, para três diferentes posições do colimador de fótons, mantendo o colimador de elétrons fixo em $\theta_e = 0,034$ rad. Nesse caso, $\Delta\phi_e = 0,1$ rad e $\Delta\phi_\gamma = \Delta\theta_\gamma = \Delta\theta_e = 4,5$ mrad. O ângulo θ_γ vale 17 mrad (pico em energia mais alta), 34 mrad e 136 mrad (pico em energia mais baixa).

Figura 12 - Mudança na energia do pico da seção de choque com a mudança na posição dos colimadores.

Para ilustrar o efeito do aumento das dimensões dos colimadores sobre os resultados, mostramos na Fig. 13 três espectros obtidos para colimadores com as posições fixas e tamanhos variáveis. Nota-se um alargamento do pico da seção de choque com o aumento da aceitação angular dos colimadores. Esse aumento, entretanto, não é suficiente para prejudicar o efeito de monocromação do feixe.

Figura 13 - Efeito de monocromatização em função da aceitação angular dos colimadores.

Condições dos colimadores para o pico mais estreito: $\theta_e = 0,068 \pm 0,0045$ rad; $\varphi_e = 0 \pm 0,05$ rad; $\theta_\gamma = 0,034 \pm 0,0045$ rad; $\varphi_\gamma = 0,15 \pm 0,05$ rad. Para o pico intermediário: $\Delta\theta_e$ e $\Delta\theta_\gamma$ duas vezes maiores. Para o pico mais largo: $\Delta\theta_e$ e $\Delta\theta_\gamma$ três vezes maiores. A linha contínua representa os resultados obtidos sem colimação no ramo dos elétrons.

A partir desse resultados, desenvolvemos um processo de produção de fótons quase monocromáticos e polarizados, utilizando mosaicos de detectores de Si, objeto da patente *Método de Seleção de Fótons Quase Monocromáticos Polarizados e Respectivo Dispositivo*, Anexo 14 (7.4.1).

5. Perspectivas

5.1. Introdução

Durante os últimos dez anos estive trabalhando no desenvolvimento de equipamentos voltados à modernização do LAL. Nesse contexto incluem-se os esforços na construção do acelerador microtron e de alguns de seus equipamentos periféricos, com ênfase no monocromador de fótons. O LAL tem tradição em pesquisa em física nuclear básica nas áreas de espectroscopia nuclear e reações fotonucleares, bem como em física aplicada com técnicas nucleares. Os novos instrumentos servirão para potencializar a atuação do LAL nessas áreas, aumentando sua eficiência e abrangência.

O novo acelerador terá corrente média 100 vezes maior que a do antigo MKII, enquanto que a corrente instantânea será 100 vezes menor, pois o fator de utilização (*duty cycle*) passará de 10^{-4} para 1. Essa característica traz várias vantagens, a maior delas sendo a possibilidade da realização de medidas em coincidência, o que viabiliza o uso do monocromador de fótons. Outro aspecto favorável diz respeito ao comportamento do feixe, que no caso do microtron, por ser um acelerador recirculado, apresenta parâmetros muito bons tanto em termos de óptica (dimensões, divergência), quanto em resolução em energia (0,1 % no microtron, contra 2,5 a 7,5 % no antigo MKII). Com isso deveremos ter menos background na área experimental. Há, no entanto, uma característica que vai piorar: a energia final. O antigo MKII chegava a 70 MeV (com restrições) e o microtron chegará a 31 MeV. Essa redução deverá restringir em alguns aspectos nossa capacidade experimental, mas o ganho com as outras vantagens (principalmente ser *cw*) certamente compensa a perda, como descrevo a seguir.

A nossa opção principal para o uso do microtron é como produtor de fótons, uma vez que a energia máxima de 31 MeV praticamente elimina a possibilidade de se realizar espalhamento de elétrons. Por outro lado, com o tagger teremos fótons monocromáticos na faixa de 3 a 27 MeV, permitindo o estudo de espectroscopia e de reações fotonucleares com possibilidades experimentais muito mais sofisticadas do que as anteriormente disponíveis. O uso de fótons de bremsstrahlung, sem monocromação, também apresenta

interesse, e o fato de dispormos de uma corrente 100 vezes mais intensa abre novas fronteiras, principalmente na área de estudo de formação de estados isoméricos (veja discussão abaixo). Por fim, na área de física aplicada, há a possibilidade de produção de fótons monocromáticos e polarizados, com energia sintonizável na faixa de dezenas de keV, através da radiação de polarização em um cristal. As possíveis aplicações desse tipo de radiação cobrem um espectro amplo, desde matéria condensada até física médica.

A Fig. 14 mostra, em vista isométrica, como deve ser a instalação dos novos equipamentos no prédio (reformado) do acelerador linear.

Figura 14 — Instalações do LAL. No andar superior encontra-se o acelerador microtron (injetor, booster e principal) e na sala experimental (andar inferior) o ramo do tagger e o de bremsstrahlung ou PRX.

Apresento a seguir alguns exemplos ilustrativos das possibilidades experimentais que devem se abrir para o LAL com o uso dos novos equipamentos.

5.2. Espalhamento de fótons em baixa energia

5.2.1. Estados isoméricos

A fotoexcitação de isômeros nucleares por reações de espalhamento inelástico de fótons é conhecida há 60 anos [18,19]. O estudo desse fenômeno, no entanto, concentrou-se em energias relativamente baixas (~ 2 MeV) ou relativamente altas, em torno do limiar de fotoprodução de partículas. Nesse caso, os resultados foram dominados pela fotoabsorção da RGE1 e a ênfase recaiu sobre as propriedades gerais da fotoexcitação ou o teste de modelos estatísticos para o decaimento gama de estados de alta energia de excitação. No caso de baixas energias, os esforços se concentraram na caracterização de estados discretos intermediários que alimentam o estado isomérico com alta probabilidade. Nessas condições, as seções de choque integradas de fotoexcitação de estados isoméricos variavam tipicamente de 0,01 a 1 keVmb. A Fig. 15 ilustra o processo de excitação e desexcitação.

Figura 15 — Representação esquemática do mecanismo de fotoexcitação ressonante para a excitação de um estado isomérico de energia E_{iso} . A banda hachurada representa um estado intermediário com energia E_j e largura radiativa total Γ . A razão de ramificação para o estado fundamental é representada por b_0 , enquanto b_{iso} representa a soma de todas as ramificações que levam ao estado isomérico.

Apenas nos últimos 10 anos foram realizados estudos sistemáticos na faixa de 2 a 7 MeV. Os resultados obtidos, que se iniciaram com uma medida da reação $^{180m}\text{Ta}(\gamma, \gamma')^{180}\text{Ta}$ [20], mostraram uma surpreendente seção de choque integrada de 100 keVmb. Medidas subsequentes em outros núclídeos apresentaram resultados semelhantes [21]. O conhecimento acurado dos processos de excitação e desexcitação de isômeros é importante por várias razões. No aspecto prático, espera-se que esse seja o caminho para a construção de um laser de gamas. Por outro lado, o acoplamento entre o estado isomérico e o fundamental apresenta grande importância em astrofísica nuclear. E, na física nuclear, a compreensão desse acoplamento, que aparentemente se dá através de estados intermediários que alimentam o isomérico, necessita de modelos e cálculos em estrutura nuclear.

O experimento é usualmente realizado em duas etapas. Inicialmente, irradia-se a amostra com fótons de bremsstrahlung, variando a energia final em passos estreitos, de forma a determinar a taxa de produção do isômero em função da energia. Em seguida, em uma experiência de fluorescência de ressonância nuclear (NRF), mede-se o espectro dos fótons espalhados no alvo. A Fig. 16 mostra os resultados obtidos no estudo da reação $^{115}\text{In}(\gamma,\gamma')^{115\text{m}}\text{In}$ [22].

Figura 16 — Taxa de produção do estado isomérico em função da energia máxima dos fótons de bremsstrahlung (parte superior). Dois pontos de inflexão, em 2,8 e 3,3 MeV, são assinalados, com indicação da seção de choque integrada até aquele ponto. A parte inferior mostra um espectro de NRF obtido com fótons de bremsstrahlung de energia máxima 4,6 MeV. Os colchetes assinalam os estados intermediários com intensidade significativa para população do estado isomérico. Da Ref. [22].

A curva da taxa de produção (parte superior da figura) mostra duas inflexões, em torno de 2,8 e 3,3 MeV. A figura mostra também que a seção de choque integrada para formação do estado isomérico, que é de 1 keVmb até 2 MeV, cresce para 7,6 keVmb em torno de 3,3 MeV. O espectro dos fótons espalhados (parte inferior da figura) mostra que a quase totalidade das transições (γ,γ') encontram-se concentradas em duas faixas em torno de 2,8 e 3,3 MeV. Dessa forma pode-se identificar os níveis intermediários responsáveis pela alimentação do estado isomérico.

Esse tipo de experiência tem perfeitas condições de ser realizado no LAL quando o microtron entrar em funcionamento.

5.3. Reações fotonucleares

O estudo da física fotonuclear tem se orientado no sentido de estudar detalhes finos dos processos fotonucleares, enfocando o importante papel que algumas pequenas amplitudes (contribuições de E2, interações de estado final, entre outras) apresentam no processo geral. Apesar de apresentarem, geralmente, contribuições muito pequenas para a seção de choque fotonuclear total, essas amplitudes são muito sensíveis aos detalhes da estrutura nuclear e dos mecanismos de reação. A descrição adequada desses processos é um excelente teste para as teorias e modelos desenvolvidos em anos recentes [23-25], que descrevem de forma adequada as propriedades estáticas e dinâmicas mais simples do núcleo atômico.

Apresento a seguir alguns exemplos de experimentos realizáveis nas novas instalações do LAL e que apresentam interesse na área de reações fotonucleares. Apesar de ser uma área muito estudada, em alguns casos uma simples medida de seção de choque (γ, n) ou (γ, p) configura um grande avanço no conhecimento disponível.

5.3.1. Fotodesintegração de núcleos leves

A fotodesintegração de núcleos leves é um campo fértil para o estudo de modelos de aglomerados e também a interação nucleon-nucleon. Um exemplo clássico nesse sentido é a fotodesintegração do dêuteron, provavelmente o processo mais estudado em física nuclear, que trouxe, e traz ainda, informações preciosas sobre o potencial nucleon-nucleon.

Como não poderia deixar de ser, um dos exemplos de como as novas instalações do LAL podem ser úteis no estudo da fotodesintegração de núcleos leves, refere-se à fotodesintegração do dêuteron. No Anexo 15 (7.1.1.2) apresentamos resultados de previsões teóricas e simulações para algumas observáveis relacionadas à reação $\bar{\gamma}d \rightarrow np$ em energias próximas ao limiar. As simulações foram feitas usando parâmetros esperados para as novas facilidades do LAL. Na região de energia próxima ao limiar os dados experimentais disponíveis são poucos e pouco acurados, não conseguindo distinguir entre os vários modelos existentes. Mostramos que, dentro das características esperadas para o

microtron e o tagger do LAL, poderemos obter resultados relevantes.

Outro exemplo de sistema leve que apresenta bastante interesse teórico e experimental é o Lítio. As reações que pretendemos estudar são as seguintes: ${}^6\text{Li}(\gamma, d_o){}^4\text{He}$ e ${}^7\text{Li}(\gamma, t_o){}^4\text{He}$, cujas características principais são mostradas na tabela abaixo.

Isótopo	Abundância	Reação	E_{limiar} (MeV)
${}^6\text{Li}$	7,5 %	(γ, d_o)	1,48
${}^7\text{Li}$	92,5 %	(γ, t_o)	2,47

Estas reações são favorecidas devido ao alto nível de clusterização dos isótopos do Li. Vários experimentos e considerações teóricas mostram que o ${}^6\text{Li}$ possui uma configuração do tipo α - d , enquanto o ${}^7\text{Li}$ apresenta uma configuração do tipo α - t [26-31]. A reação ${}^6\text{Li}(\gamma, d_o){}^4\text{He}$ é proibida pela lei de conservação de isospin para transições isovetoriais $E1$. Assim, essa reação ocorre predominantemente via transições isoescalares $E2$. Essa reação foi medida a energias em torno de 59 MeV e os resultados experimentais de distribuição angular confirmam a presença de uma forte componente $E2$ [32]. A baixas energias (em torno de 3 MeV), o canal d_o foi medido através da reação inversa ${}^2\text{H}(\alpha, \gamma){}^6\text{Li}$ [33]. Esse experimento mostrou que, nestas energias, apesar do domínio das transições $E2$, as transições $E1$ são importantes e não podem ser desprezadas [34]. Até o momento não existem medidas da reação direta ${}^6\text{Li}(\gamma, d_o){}^4\text{He}$ a baixas energias.

A reação ${}^7\text{Li}(\gamma, t_o){}^4\text{He}$ foi medida entre 6 e 50 MeV, com distribuições angulares tendo sido obtidas para energias de excitação entre 6 e 15 MeV [35,36]. Cálculos recentes utilizando o modelo de cluster, assumindo a configuração α - t para o ${}^7\text{Li}$ [31], indicam que as distribuições angulares mudam drasticamente entre 4 e 9 MeV, conforme mostra a Fig. 17. As medidas existentes, no entanto, não abrangem a região onde deve haver predominância de $E2$ e, além disso, não apresentam estatística suficiente para esclarecer a situação. Assim, a medida das distribuições angulares nessas energias será um excelente teste para o modelo de cluster.

Figura 17 — Distribuições angulares dos trítios para energias dos fótons de 4 MeV (a), 5 MeV (b) e 9 MeV (c). As linhas tracejadas representam resultados supondo apenas transições E1, as ponto-traço apenas transições E2 ($\times 10$) e as linhas contínuas os resultados supondo interferência E1-E2. Da Ref. [31].

Como parte do desenvolvimento experimental necessário para a nova fase do LAL, projetamos e construímos um sistema de detecção para partículas carregadas que utiliza detectores de Si divididos em faixas, produzidos na Ucrânia e caracterizados no LAL, veja o Anexo 16 (7.1.2.12). O sistema, projetado para detectar partículas carregadas emitidas em reações fotonucleares, apresenta grande ângulo sólido e boa resolução angular e em energia. Uma descrição mais detalhada é apresentada no Anexo 17 (7.1.2.10). O sistema, em uma montagem parcial, foi utilizado nas dependências do

IPEN/CNEN-SP, para a medida das distribuições angulares dos trítios emitidos na reação ${}^7\text{Li}(\gamma, t)$. Os resultados são apresentados no Anexo 18 (7.1.2.3).

5.3.2. Ressonâncias Gigantes

Fótons monocromáticos de alta resolução e intensidade, com energia sintonizável, constituem ferramenta extremamente poderosa no estudo das propriedades das ressonâncias gigantes fotonucleares. Algumas das características importantes que apresentam problemas em aberto são os canais de decaimento, o acoplamento entre a RGE1 e vibrações da superfície do núcleo, e o estudo da intensidade MI na região de energia próxima ao limiar de fotoprodução de partículas. Nesse último caso, a disponibilidade de feixes de fótons monocromáticos e linearmente polarizados é fundamental para viabilizar os experimentos.

O conhecimento detalhado da seção de choque de fotoabsorção e de seus canais de decaimento é necessário para inúmeras aplicações, desde cálculos de blindagem em instalações que usam aceleradores (cada vez mais numerosas) até cálculos de dose em tratamentos por radioterapia. Nesse contexto, um exemplo é particularmente esclarecedor. Trata-se do Ferro, elemento abundante na natureza e de extrema importância industrial e comercial, largamente utilizado na construção de aceleradores, seus periféricos e na estrutura dos prédios que os abrigam. As seções de choque do Fe foram medidas usando feixes de bremsstrahlung, sendo que os poucos resultados existentes, de seções de choque (γ, n) e (γ, p) obtidas por deconvolução, apresentam discrepâncias significativas. Há muito interesse, por parte da AIEA (Agência Internacional de Energia Atômica), no desenvolvimento de bancos de dados completos e confiáveis de seções de choque totais e parciais de elementos importantes para usos industriais, médicos e de blindagem. O LAL poderá suprir boa parte desses dados necessários quando o microtron e o tagger estiverem em operação.

5.4. Aplicações

Um laboratório como o LAL, com as instalações e os equipamentos que dispõe, apresenta grande potencial para aplicações em várias áreas. O LAL tem tradição na

prestação de serviços em problemas que requeiram o uso de técnicas nucleares (análise de materiais, contaminantes, irradiação de componentes eletrônicos, etc.). Há, no entanto, uma área que considero particularmente interessante e promissora. Trata-se da produção de fótons de R-X monocromáticos e polarizados, com energia sintonizável, através da radiação de polarização, PRX.

Quando um feixe de elétrons ultra relativísticos incide sobre um meio material, o bremsstrahlung produzido pode ser escrito como a soma de dois termos, um deles o bremsstrahlung de polarização, definido pela polarização dinâmica do átomo e o outro o bremsstrahlung estático da partícula nos campos coulombianos do núcleo e dos elétrons. Em um meio amorfo as distribuições angulares são bem separadas, pois o bremsstrahlung de polarização é praticamente isotrópico e não depende da energia do feixe incidente. Quando o meio material tem estrutura cristalina, o bremsstrahlung de polarização apresenta características muito interessantes, pois a cinemática do processo permite apenas alguns valores para o momento transferido. Com isso a radiação é emitida essencialmente na direção dos ângulos de Bragg, passando a ter um caráter coerente e uma seção de choque mais elevada. Esse tipo de radiação é também conhecido na literatura como radiação paramétrica, ou radiação paramétrica do tipo B. A Fig. 18 mostra um espectro de PRX produzido por elétrons de 8,3 MeV incidindo em um cristal de diamante [37]. Dois aspectos são particularmente relevantes nessa figura: a boa resolução em energia da radiação emitida e a separação entre o sinal e o fundo, característica muito importante para aplicações médicas e industriais.

Figura 18 — Espectro de PRX obtido quando elétrons de 8,3 MeV incidem sobre um cristal de diamante de 55 μm de espessura. A medida foi feita utilizando um detector de Si(Li) colocado a 44° em relação à direção do feixe de elétrons. Da Ref. [37].

Além dessas duas características excepcionais, a energia da radiação emitida varia como função do ângulo do cristal, como pode ser visto na Fig. 19, o que permite que a energia seja sintonizável através do controle do goniômetro.

Figura 19 — Dependência da energia da radiação PRX emitida em função do ângulo de inclinação do cristal. Da Ref. [37].

Um aspecto que deve ser enfatizado é o de que a PRX apresenta interesse também como objeto de estudo, pois apesar de teoricamente conhecida há anos, apenas recentemente estudos experimentais sistemáticos começaram a ser feitos [37,38].

6. Conclusões

A excitação do núcleo atômico através da radiação eletromagnética é uma ferramenta importante para o estudo de propriedades nucleares. Experimentos inclusivos, como a eletro e a fotodesintegração, permitiram a obtenção de resultados relevantes, mas em geral de caráter médio ou integrado. A obtenção de informações sobre propriedades relacionadas à dinâmica dos processos nucleares exige a realização de experiências exclusivas, muitas vezes em condições de cinemática completa.

O LAL operou durante 25 anos o acelerador linear MKII doado pela Universidade de Stanford na década de 60. A operação dessa máquina foi um sucesso em vários aspectos. Do ponto de vista científico propiciou a formação de muitos estudantes e pesquisadores e a publicação de trabalhos de relevância nas áreas de espectroscopia nuclear e reações fotonucleares. Do ponto de vista tecnológico, a operação do acelerador foi muito produtiva na formação de técnicos e em gerar capacitação na área de manutenção e desenvolvimento de aceleradores. A capacitação acumulada no aprimoramento dos vários sistemas do acelerador (microondas, vácuo, óptica) ao longo desse período, permitiu a realização do projeto de um novo acelerador, com características mais adequadas às necessidades da pesquisa básica, como operação em onda contínua (*cw*), corrente intensa (50 μ A) e óptica de feixe muito bem comportada (característica dos aceleradores recirculados).

Para que o potencial que o microtron oferece possa ser utilizado em sua plenitude, é necessário um conjunto de equipamentos periféricos que permitam a realização de medidas de caráter exclusivo, em linha. Sem isso o uso do acelerador não se alteraria na prática, permitindo essencialmente a realização de medidas por ativação (fora de linha).

O monocromador de fótons de bremsstrahlung do LAL permitirá a geração de fótons monocromáticos com resolução ~ 1 %, excelente ferramenta para experimentos de fotodesintegração e espalhamento de fótons. O LAL já dispõe de equipamento e capacitação para medidas de fótons de alta energia em alta resolução, o que viabiliza toda uma linha de pesquisa nas áreas de espalhamento de fótons e NRF. A possibilidade de

produzir fótons monocromáticos e linearmente polarizados abre também perspectivas interessantes de estudo dos modos de excitação MI nos núcleos.

Além disso desenvolvemos sistemas de detecção de partículas carregadas, utilizando detectores de Si, que tornam exeqüíveis experimentos de fotodesintegração por emissão de partículas carregadas em linha e com observação direta do alvo. Nessa categoria incluem-se estudos experimentais de canais de decaimento de ressonâncias gigantes, quebra de núcleos leves que apresentam estrutura de aglomerados e fotofissão.

Assim, podemos concluir que a operação do microtron do LAL, associado aos equipamentos periféricos desenvolvidos, permitirá excelentes perspectivas experimentais na área de fotoexcitação do núcleo. As características do acelerador são também propícias para seu uso em aplicações, como no caso da geração de raios-X monocromáticos e polarizados através da radiação de polarização em um cristal.

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Caderno de Anexos

(e,p) and (e, α) reactions in ^{90}Zr and ^{92}Zr

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The yields of protons and α particles from 2 mg/cm² targets of ^{90}Zr and ^{92}Zr have been measured in the incident electron energy range 20–100 MeV; the $^{90}\text{Zr}(e,p)$ and $^{90}\text{Zr}(e,\alpha)$ data were extended to 130 MeV. Photodisintegration plus electrodisintegration yields were also measured for electron energies above 50 MeV. The photodisintegration cross sections, derived from these data, rise continuously from 25 MeV onward for all four reactions. One satisfactory explanation of this phenomenon is that we are observing multiparticle emission following virtual photon absorption.

I. INTRODUCTION

Some time ago Tamae *et al.*¹ reported on a measurement of the $^{90}\text{Zr}(e,\alpha)$ cross section and the photonuclear cross section derived from it. The latter contained a peak resulting from the E1 giant resonance as well as an increasing cross section above 30 MeV. At the highest excitation energy measured by Tamae *et al.*,¹ 60 MeV, the cross section exceeded the giant dipole cross section. This curious feature led us to embark on a measurement of the (e,p) and (e, α) cross sections of ^{90}Zr and ^{92}Zr in the hope of reaching some understanding of this interesting phenomenon.

II. EXPERIMENT

The (e,p) and (e, α) cross sections for ^{90}Zr and ^{92}Zr have been measured in the incident electron energy range 20–100 MeV. The ^{90}Zr measurements were extended to 130 MeV. The experimental apparatus and the National Bureau of Standards (NBS) linear electron accelerator have been described previously.^{2–4} Target thicknesses, isotopic purities, and the separation energies for the decay channels of interest are presented in Table I. For electron energies between 50 and 100 MeV photodisintegration plus electrodisintegration yields were measured by interposing a 136 mg/cm² Ta radiator 7.6 cm upstream of the Zr targets. The effects of multiple scattering and electron energy loss in the 136 mg/cm² radiator used were checked with a 217 mg/cm² radiator at 60 and 100 MeV in a previous experiment.⁵ At these energies when the electrodisintegration yields were subtracted from the photodisintegration plus electrodisintegration yields for both radiators, the (counts/Coulomb) (mg/cm²) of radiator were the same within counting statistics.

Five silicon surface barrier charged particle detectors were mounted in the spectrometer focal plane. Three of these detectors were 10×40 mm rectangular detectors and two were 22 mm circular detectors. All detectors could be totally depleted. The absolute efficiency of the *i*th counter is defined to be

$$(\text{Efficiency})_i = \frac{\int C(p)_i d\Omega dp}{4\pi S_0}, \quad (1)$$

TABLE I. Separation energies and target properties.

Reaction	Separation energy	
	^{90}Zr (MeV)	^{92}Zr (MeV)
(γ ,p)	8.36	9.40
(γ ,n)	11.98	8.64
(γ ,2p)	15.43	17.11
(γ ,np)	19.83	17.34
(γ ,2n)	21.29	15.84
(γ ,n2p)	26.55	24.91
(γ ,2np)	29.21	24.20
(γ ,n3p)	35.96	35.79
(γ ,3np)	37.82	35.67
(γ ,n4p)	48.16	44.99
(γ ,4np)	50.50	45.05
(γ ,n5p)	55.40	57.01
(γ ,5np)	62.72	56.87
(γ , α)	6.68	2.97
(γ ,2 α)	13.02	10.87
(γ ,ap)	16.32	13.58
(γ ,an)	18.13	14.09
(γ ,3 α)	19.01	17.97
(γ ,2ap)	22.93	21.58
(γ ,2an)	24.01	21.39
(γ ,4 α)	25.04	24.94
(γ ,anp)	26.80	23.50
(γ ,3ap)	29.41	29.33
(γ ,3an)	29.51	27.87
(γ ,2apn)	33.09	31.16
(γ ,4an)	35.24	34.39
(γ ,4ap)	36.01	36.91
Enrichment (%)	97.62	95.13
Thickness (mg/cm ²)	2.02	2.00

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where $C(p)_i$ stands for the number of counts recorded by a focal plane counter when a calibrated alpha source of strength S_0 is located in the target position of the spectrometer. Relative values of the efficiencies obtained from alpha source measurements agreed within counting statistics with a series of measurements made with each counter measuring protons and alphas of the same energy from a smooth portion of a (e,p) or (e, α) energy spectrum. The centroid of the distribution $C(p)_i$ provided the spectrometer momentum calibration for the i th counter. ^{241}Am and ^{238}U calibrated alpha sources provided by the NBS radioactivity section were used for these measurements. The dependence of the counter efficiency on beam spot size was determined by direct comparison of the counting rate with a 2 mm beam spot and with a 22 mm beam spot. The counting rates under these two extreme conditions differed by only 4.5%. Since the beam size did not exceed 10 mm even at the lowest electron energy, we assumed our counting efficiency was independent of electron energy.

For each target complete electrodisintegration energy spectra, $d^2\sigma/d\Omega dT$, were taken at 90° for incident electron energies, E_0 , of 20, 22, 24, 30, 40, 50, 60, 80, and 100 MeV. Complete photodisintegration plus electrodisintegration energy spectra were taken at 50, 60, 80, and 100 MeV. At other incident electron energies, E_0 , between 20 and 100 MeV, the proton and α -particle cross sections were sampled in 2 MeV steps at a single magnetic field setting of the spectrometer roughly corresponding to the peak of the α -particle cross section. These cross sections were also measured in 5 MeV steps for ^{90}Zr between 100 and 130 MeV. These "bite" cross section measurements were converted into complete cross sections by using ratios, $R_\alpha(E_0)$ and $R_p(E_0)$, obtained from the complete cross section measurements. The ratios $R_\alpha(E_0)$ and $R_p(E_0)$ were defined to be

$$R_i(E_0) = \frac{\int_{T_{\text{bite}}}^{T_{\text{max}}} (d^2\sigma/d\Omega dT)_i dT}{\int_{T_{\text{bite}}} (d^2\sigma/d\Omega dT)_i dT} \quad (2)$$

The contribution to the (e, α) and (e,p) yields from the energy region, which lies above the maximum energy which the spectrometer magnet can bend (≈ 25 MeV), was estimated by using the α -particle data of Flowers *et al.*⁶ for ^{92}Mo and ^{94}Mo , and the proton data of Schumacher *et al.*⁷ for Cu, and Tonapetyan *et al.*⁸ for ^{93}Nb . In the energy region above 20 MeV both the proton and α -particle energy distributions are approximately given by

$$\frac{d^2\sigma}{d\Omega dT_i} = \left[\frac{d^2\sigma}{d\Omega dT_i} \right]_0 e^{-\lambda T_i}, \quad (3)$$

with

$$\lambda_\alpha = (0.16 \pm 0.02)/\text{MeV}$$

and

$$\lambda_p = (0.021 \pm 0.003)/\text{MeV}.$$

These values of λ are nearly constant for a wide range of mass numbers A ($\lambda_p = 0.021/\text{MeV}$ for ^{93}Nb at a bremsstrahlung end-point energy of 340 MeV and $0.023/\text{MeV}$ for Cu at end-point energies of 150 and 340 MeV). The values $R_p(E_0)$ and $R_\alpha(E_0)$ are given in Tables II and III. Where a second set of values for $R_p(E_0)$ is given, it includes an estimate of the proton yield above 25 MeV. These larger values were used to evaluate the (e,p) cross sections. The high-energy tail for the α particles made a negligible contribution. The larger value of $R_\alpha(E_0)$ at each E_0 contains the α yield below 5.6 MeV displayed in Fig. 2. The cross sections were evaluated excluding the yield below 5.6 MeV since we believe it is produced by a contaminant. Since $R_\alpha(E_0)$ and $R_p(E_0)$ were slowly varying functions of E_0 except within a few MeV of the threshold energies, these ratios between measured values were obtained by a linear interpolation and only the error estimate of the bite area and the ratios $R_\alpha(E_0)$ and $R_p(E_0)$ were included in the error estimate of the total cross section.

The linac beam currents were measured with an NBS beam dump during electrodisintegration measurements

TABLE II. The ratio, R_p , defined in Eq. (2) as a function of the incident electron energy, E_0 . Where a second set of ratios is given, it includes an estimate of the proton yield above 25 MeV (see the text for details).

E_0	$^{90}\text{Zr}(e,p)$	$^{90}\text{Zr}(e + \gamma, p)$	$^{92}\text{Zr}(e,p)$	$^{92}\text{Zr}(e + \gamma, p)$
20	60.8 $\pm$ 1.9		79.9 $\pm$ 4.3	
22	20.24 $\pm$ 0.20		13.8 $\pm$ 0.2	
24	18.9 $\pm$ 0.2		10.6 $\pm$ 0.2	
30	17.4 $\pm$ 0.2		10.9 $\pm$ 0.1	
	17.6 $\pm$ 0.2		11.0 $\pm$ 0.1	
40	17.0 $\pm$ 0.2		11.3 $\pm$ 0.1	
	17.9 $\pm$ 0.2		12.3 $\pm$ 0.1	
50	15.9 $\pm$ 0.1	16.0 $\pm$ 0.1	10.7 $\pm$ 0.1	10.9 $\pm$ 0.1
	17.0 $\pm$ 0.2	17.1 $\pm$ 0.2	11.9 $\pm$ 0.2	12.1 $\pm$ 0.2
60	16.7 $\pm$ 0.1	16.4 $\pm$ 0.1	11.7 $\pm$ 0.1	11.8 $\pm$ 0.1
	17.7 $\pm$ 0.1	19.9 $\pm$ 0.1	12.5 $\pm$ 0.2	12.5 $\pm$ 0.2
80	16.1 $\pm$ 0.1	16.5 $\pm$ 0.1	11.9 $\pm$ 0.1	11.9 $\pm$ 0.1
	17.3 $\pm$ 0.2	17.9 $\pm$ 0.1	13.3 $\pm$ 0.3	13.5 $\pm$ 0.1
100	15.8 $\pm$ 0.1	15.9 $\pm$ 0.1	11.7 $\pm$ 0.1	11.6 $\pm$ 0.1
	18.1 $\pm$ 0.1	18.4 $\pm$ 0.2	15.1 $\pm$ 0.2	14.5 $\pm$ 0.2

TABLE III. The ratio, R_α , defined in Eq. (2) as a function of the incident electron energy, E_0 . The second set of ratios given at each E_0 do not include the alpha yield below 5.6 MeV (see the text for details).

E_0	$^{90}\text{Zr}(e,\alpha)$	$^{90}\text{Zr}(e+\gamma,\alpha)$	$^{92}\text{Zr}(e,\alpha)$	$^{92}\text{Zr}(e+\gamma,\alpha)$
20	4.65 $\pm$ 0.31		2.94 $\pm$ 0.11	
	4.19 $\pm$ 0.20		2.87 $\pm$ 0.11	
22	4.24 $\pm$ 0.20		2.91 $\pm$ 0.10	
	3.79 $\pm$ 0.17		2.81 $\pm$ 0.09	
24	3.73 $\pm$ 0.13		2.92 $\pm$ 0.09	
	3.43 $\pm$ 0.12		2.81 $\pm$ 0.09	
30	3.79 $\pm$ 0.15		3.21 $\pm$ 0.07	
	3.50 $\pm$ 0.14		3.11 $\pm$ 0.07	
40	4.12 $\pm$ 0.12		3.30 $\pm$ 0.08	
	3.76 $\pm$ 0.11		3.15 $\pm$ 0.08	
50	4.23 $\pm$ 0.10	4.04 $\pm$ 0.06	3.63 $\pm$ 0.08	3.61 $\pm$ 0.05
	3.86 $\pm$ 0.10	3.73 $\pm$ 0.06	3.61 $\pm$ 0.08	3.46 $\pm$ 0.05
60	4.50 $\pm$ 0.05	4.42 $\pm$ 0.05	4.01 $\pm$ 0.07	4.01 $\pm$ 0.05
	4.06 $\pm$ 0.04	4.02 $\pm$ 0.04	3.80 $\pm$ 0.06	3.76 $\pm$ 0.04
80	4.84 $\pm$ 0.07	4.81 $\pm$ 0.06	4.43 $\pm$ 0.07	4.43 $\pm$ 0.05
	4.35 $\pm$ 0.05	4.43 $\pm$ 0.05	4.06 $\pm$ 0.06	4.16 $\pm$ 0.05
100	4.81 $\pm$ 0.06	4.87 $\pm$ 0.04	4.62 $\pm$ 0.06	4.62 $\pm$ 0.05
	4.48 $\pm$ 0.07	4.54 $\pm$ 0.04	4.39 $\pm$ 0.04	4.37 $\pm$ 0.04

and with a ferrite beam current monitor calibrated relative to the beam dump during radiator-in measurements. Alpha particles and protons were separated by varying the bias of the silicon transmission detectors so that protons were not stopped in the depletion layer while alphas were always stopped except at the highest fields where the range of the alphas exceeded the thickness of the detectors.

III. RESULTS

The proton energy spectra, $(d^2\sigma/d\Omega dT_p)$, from ^{90}Zr and ^{92}Zr bombarded with 20 MeV electrons are shown in Fig. 1. All of the structure seen in the ^{90}Zr energy spectrum has been reported previously.⁹⁻¹⁴ The strong state which occurs at a center-of-mass proton energy of 7.94 MeV, corresponding to an excitation energy of 16.28 MeV, is a 1^- analog state which has been widely studied. The peak at 6.40 MeV results from the decay of the 16.28 MeV state to the 1.51 MeV, $\frac{3}{2}^-$ second excited state in ^{89}Y . The peak at 6.00 MeV is the ground state decay of a 1^- analog state with an excitation energy of 14.43 MeV. The energy difference between the small peak at 4.9 MeV and the 6.40 MeV peak suggests that the 4.9 MeV peak is due to the decay of the 14.43 MeV 1^- state to the 1.51 MeV, $\frac{3}{2}^-$ second excited state in ^{89}Y .

The prominent peak near 4 MeV has been reported extensively.^{9-12,14,15} This peak is thought to be a pseudoresonance produced by the opening of the neutron channel at an effective threshold of 12.568 MeV.¹² Since the ground state of ^{89}Zr is $\frac{3}{2}^+$, ground-state neutron transitions are strongly inhibited by the angular momentum barrier, and the neutron channel remains effectively closed until the first excited $\frac{1}{2}^-$ state at 0.588 MeV is energetically feasible. Since the total photon-absorption cross section should be smooth, the behavior of the (γ,p) cross section implies a very sharp rise in the (γ,n) cross section near threshold.

Although the $^{92}\text{Zr}(\gamma,n)$ reaction has been well studied,¹⁶ the $^{92}\text{Zr}(\gamma,p)$ reaction has not been reported in the literature. The strong peak seen at 7.38 MeV in Fig. 1(b) probably results from ground-state transitions from a 1^- ana-

FIG. 1. The proton spectra produced at 90° when 20 MeV electrons are incident on targets of ^{90}Zr and ^{92}Zr .

log state at an excitation energy of 16.8 MeV. The smaller peaks seen at 7.03 and 5.90 MeV are probably the decays of the 16.8 MeV state to the $\frac{3}{2}^-$ excited states at 0.556 and 1.474 MeV in ^{91}Y . The pseudoresonance found in the ^{90}Zr spectrum is, of course, not seen here, because the neutron separation energy is less than the proton separation energy.

The α spectra for ^{90}Zr and ^{92}Zr , ($d^2\sigma/d\Omega dT_\alpha$), observed at 90° using 60 MeV incident electrons are shown in Figs. 2(a) and (b). They have a peak characteristic of evaporation spectra for α -particle energies near the Coulomb barrier height. The high-energy tails of the α spectra are usually associated with preequilibrium processes.⁶

Both of these spectra display a group of α particles near 3 MeV. It is hard to understand how α particles of such low energy could survive the Zr Coulomb barrier. We believe that they are produced in the $^{12}\text{C}(e,3\alpha)$ reaction because (1) there is a dark spot on the targets the size and shape of the beam and (2) after a 30-minute exposure to 30 MeV electrons there was a component in the decay curve of the annihilation radiation having a half-life consistent with that of ^{11}C . For this reason the low-energy yield in the α spectra has been excluded in the evaluation of the cross sections.

FIG. 2. The α -particle spectra produced at 90° when 60 MeV electrons are incident on targets of ^{90}Zr and ^{92}Zr .

IV. DATA ANALYSIS

The data were treated in much the same way as described in Ref. 4. The measured spectra, $d^2\sigma/d\Omega dT$, were integrated over particle energy to obtain $d\sigma/d\Omega$ (90°), and the latter was multiplied by 4π to yield the total cross section, $\sigma_{e,x}(E_0)$. The photodisintegration plus electrodisintegration yields, $Y_{e,x}(E_0)$, were obtained by multiplying by the same factors. These two measured quantities, $\sigma_{e,x}(E_0)$ and $Y_{e,x}(E_0)$, are plotted as a function of total incident electron energy, E_0 , in Figs. 3–6 for the four reactions studied here.

The relationship between the photodisintegration cross section, $\sigma_{\gamma,x}(E_0)$, and the measured electrodisintegration cross section is

$$\sigma_{e,x}(E_0) = \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(\omega) N^{\lambda L}(E_0, \omega, Z) \frac{d\omega}{\omega}, \quad (4)$$

where $N^{\lambda L}(E_0, \omega, Z)$ is the virtual photon intensity spectrum of multipolarity λL . The yield with the radiator in is

$$Y_{e,x}(E_0) = \sigma_{e,x}(E_0 - 2\Delta E_0) + n_r \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(\omega) K(E_0 - \Delta E_0, \omega) \frac{d\omega}{\omega}. \quad (5)$$

Here n_r is the number of nuclei/cm² in the tantalum radiator and ΔE_0 is the electron energy loss in half the radiator thickness. For the bremsstrahlung cross section, $K(E_0, \omega)$, we have used the Davies-Bethe-Maximon cross section as given in Ref. 17.

The cross sections, $\sigma_{e,x}(E_0)$, and yields, $Y_{e,x}(E_0)$, of Figs. 3–6 have been simultaneously fitted using the E1

FIG. 3. The measured $\sigma_{e,p}(E_0)$ for ^{90}Zr as a function of total incident electron energy E_0 (lower points). The upper points represent the yield $Y_{e,p}(E_0)$ obtained when a 0.136 g/cm² tantalum foil was placed in the electron beam ahead of the target. The smooth curves are the best fits to the data and were obtained by combining the histograms representing the E1 and E2 (γ, p) cross sections (right-hand scale) in Eqs. (1) and (2) with the E1 and E2 DWBA virtual photon spectra and by making use of the Davies, Bethe, and Maximon (DBM) bremsstrahlung cross section. The size effect correction described in the text has been applied to the virtual photon spectra.

FIG. 4. The $\sigma_{e,p}(E_0)$ for ^{92}Zr . See caption of Fig. 3. FIG. 6. The $\sigma_{e,p}(E_0)$ for ^{92}Zr . See caption of Fig. 3.

and $E2$ virtual photon spectra calculated from the distorted-wave Born approximation (DWBA) code of Soto Vargas *et al.*¹⁸ This code assumes a point nucleus and hence neglects the momentum dependence of the nuclear form factors and the interference between the nuclear form factors and the Coulomb distortion effects. We have introduced a simple expression to correct for the momentum dependence of the form factors (i.e., reduced matrix elements) which reproduces the measured rms nuclear charge radius R . This correction consists in multiplying $N^{\lambda L}(E_0, \omega, Z)$ inside the integrals of Eq. (4) and (5) by

$$F^L(qR) = \left[\left| \frac{\omega}{q} \right| \frac{j_L(qR)}{j_L(\omega R)} \right]_{q=q_{\text{rms}}}^2 \quad (6)$$

The rms values of q were obtained in plane-wave Born approximation (PWBA) as described in Ref. 2. Because of the strong forward peaking of the $E1$ ($e, e'x$) cross section from which the (e, x) cross section and virtual photon spectra are derived, the model dependence of the nuclear form factor correction is very small for $E1$ transitions as long as the form factors reproduce the measured ground-state charge and transition radii.

The fits to the electrodisintegration and photodisintegration plus electrodisintegration cross sections are shown in Figs. 3–6, the photodisintegration cross sections

being represented by very coarse histograms. The $E2$ strength is represented by a single bin at the energy of the isovector $E2$ resonance. Durgapal and Onley¹⁹ have written a second order Born approximation (SOBA) computer code which is valid for Zr. This code takes into account a certain class of nuclear form factors as well as the interference between the form factors and the Coulomb distortion. The results of using the SOBA virtual photon spectra to analyze our data are the second entries in Table IV. The χ^2 test for goodness of fit are worse using the SOBA spectra without $E2$ isovector strength, but with $E2$ isovector strength both SOBA and DWBA with a multiplicative form-factor correction give roughly equally good fits.

The most remarkable feature of the derived $^{90}\text{Zr}(\gamma, p)$ and (γ, α) cross sections is that they are still rising at the highest excitation energies we sampled, 130 MeV. The same remarks apply to the $^{92}\text{Zr}(\gamma, p)$ and (γ, α) cross sections up to 100 MeV. Thus we confirm the principle result of Tamae *et al.*¹ for the $^{90}\text{Zr}(\gamma, \alpha)$ cross section and, in addition, show that the “apparent cross section rise” is also seen in the $^{90}\text{Zr}(\gamma, p)$, $^{92}\text{Zr}(\gamma, p)$, and $^{92}\text{Zr}(\gamma, \alpha)$ decay channels. We shall argue below that we believe that the apparent cross section rise is not real but due to our inability to measure more than one particle in the final state. Our $^{90}\text{Zr}(e, \alpha)$ excitation function rises about 20% faster than that of Tamae *et al.*¹ in the excitation energy interval of 20 to 60 MeV and the average value of ours in this energy range is about 27% smaller than theirs.

V. DISCUSSION

This rather strange behavior raises the question whether (1) it is an artifact of the virtual photon spectrum used in the analysis⁶ or (2) it results from multiparticle emission following virtual photon absorption. We have already shown⁵ that the electric dipole virtual photon spectrum, including Coulomb distortions and nuclear size corrections, quite adequately describes the electroexcitation of a discrete level. We cannot be quite so confident concerning electrodisintegration in the quasideuteron region.

Chang *et al.*²⁰ and Tiator and Wright^{21,22} have pointed out that the virtual photon spectra for light systems are

FIG. 5. The $\sigma_{e,a}(E_0)$ for ^{90}Zr . See caption of Fig. 3.

TABLE IV. Integrated cross sections. There are two entries for each reaction. The first is evaluated using the DWBA virtual photon spectrum with the form factor correction and the second using the SOBA spectrum.

90Ba spectrum.

Reaction	Reduced χ^2	E1 fit $\int_{30}^{100} \sigma^{E1} d\omega$ (MeV mb)	$\int_{30}^{100} \sigma^{E1} d\omega$ (MeV mb)	$\int_{30}^{130} \sigma^{E1} d\omega$ (MeV mb)
$^{90}\text{Zr}(\gamma, \alpha)$	1.63 2.94	3.5 ± 0.3 3.7 ± 0.3	53.6 ± 2.0 55.6 ± 2.0	95.5 ± 5.3 93.5 ± 5.4
$^{90}\text{Zr}(\gamma, p)$	0.74 1.26	281 ± 25 305 ± 26	813 ± 77 785 ± 79	1730 ± 170 1570 ± 173
$^{92}\text{Zr}(\gamma, \alpha)$	1.10 2.24	3.5 ± 0.4 2.3 ± 0.3	47.2 ± 2.2 50.2 ± 2.3	
$^{92}\text{Zr}(\gamma, p)$	1.35 2.33	149 ± 8 156 ± 8	664 ± 45 661 ± 64	

Reaction	Reduced χ^2	E1 + isovector $\int_{30}^{100} \sigma^{E1} d\omega$ (MeV mb)	E2 fit $\int_{30}^{100} \sigma^{E1} d\omega$ (MeV mb)	$\int_{30}^{130} \sigma^{E1} d\omega$ (MeV mb)	$\int \sigma^{E2} d\omega$ (MeV mb)
$^{90}\text{Zr}(\gamma, \alpha)$	0.81 2.11	2.4 ± 0.3 2.9 ± 0.2	48.9 ± 2.0 52.1 ± 2.0	84.3 ± 5.2 91.3 ± 5.4	0.53 ± 0.12 6.9 ± 0.4
$^{90}\text{Zr}(\gamma, p)$	0.65 0.91	271 ± 20 280 ± 15	744 ± 147 695 ± 156	1389 ± 168 1470 ± 171	11.9 ± 4.6 76 ± 15
$^{92}\text{Zr}(\gamma, \alpha)$	0.85 1.27	3.3 ± 0.4 3.7 ± 0.3	42.6 ± 2.3 47.8 ± 2.3		0.69 ± 0.14 1.90 ± 0.24
$^{92}\text{Zr}(\gamma, p)$	0.60 1.27	141 ± 9 138 ± 15	545 ± 46 569 ± 88		18.7 ± 2.8 50.5 ± 5.5

substantially modified by recoil effects. To test these ideas we have assumed that for excitation energies greater than ~ 40 MeV the dominant absorption takes place via the quasideuteron mechanism in which the $(A-2)$ system is simply a spectator and that the observed protons and α particles are generated in the final state interactions of the neutron-proton pair produced by the initial photon absorption. We have generated a virtual photon spectrum that includes exact three-body kinematics, two nucleons, and a scattered electron. Since our range of excitation en-

ergies extends beyond the region where the neglect of the Coulomb, transverse-Coulomb interference, and transverse-transverse terms of the coincidence $(e, e'x)$ cross section is justified, we include all of the terms in a numerical calculation of the virtual photon spectrum for $A=2$. Including these terms forced us to abandon the forward peaking approximation of Tiator and Wright²¹ and to work either in the long-wavelength limit (LWL) or to include model-dependent form factors in our calculation.²³ We assumed the virtual photon spectrum to be

$$N_{\text{QD}}^{E1}(E_0, \omega, A, A) = \{ [1 - (e^{-L/\omega})] N_{\text{FWBA}}^{E1}(E_0, \omega) R_{\text{FF}}(E_0, \omega, Z_{\text{Tr}}, A_{\text{Zr}}) + (e^{-L/\omega}) N_{\text{QD}}^{E1}(E_0, \omega, A_D) R_{\text{FF}}(E_0, \omega, Z_D, A_D) \} R_C(E_0, \omega, Z), \quad (7)$$

where $L \approx 80$ MeV is the Levinger quenching factor,²⁴ and the functions $R_{\text{FF}}(E_0, \omega, Z, A)$ and $R_C(E_0, \omega, Z)$ are corrections for form factor and Coulomb distortion effects and are described more fully in Ref. 23. In the analysis of our experiment we integrate over all energies and angles of the emitted protons or α particles at each E_0 . Because of the recoil of the correlated nucleon pair, decay protons of the same kinetic energy, T_p , can be produced by the absorption of a range of virtual photon ener-

gies and an even wider range of virtual photon momentum transfer. The $N_{\text{QD}}^{E1}(E_0, \omega)$ we require is given by

$$N_{\text{QD}}^{E1}(E_0, \omega) = \int d\Omega_p dT_p d\Omega_{e,e'} \times \frac{d^3\sigma_{e,p}[E_0, T_p(\theta_{e,e'}, \phi_{e,e'})]}{d\Omega_p dT_p d\Omega_{e,e'}}, \quad (8)$$

where the integration over the $(e, e'p)$ coincidence cross section²³ is carried out subject to the constraint that only

FIG. 7. The quasideuteron virtual photon spectrum $N_{QD}^{E1}(100, \omega)$ defined in Eq. (8) and the infinite mass, no recoil, spectrum $N_0^{E1}(100, \omega)$.

those values of $\theta_{e,e'}$, $\phi_{e,e'}$, and θ_p which produce the same ω are included in the integration. Experimental values^{25,26} for the deuteron photodisintegration angular distribution coefficients, $a_0(\omega)$ and $a_2(\omega)$, were used in the evaluation of Eq. (8). After the integration of Eq. (8) over $d\Omega_{ee'}$, the result was divided by

$$a_0(\omega_\gamma)P_0(\cos\theta_{ep}) + a_2(\omega_\gamma)P_2(\cos\theta_{ep})$$

and other electron kinematic variables to obtain the quasideuteron virtual photon spectrum. Without the $(\theta_{ee'}, \phi_{ee'})$ dependence introduced by recoil effects, the transverse-Coulomb and transverse-transverse terms of the coincidence cross section would not contribute to $N_{QD}^{E1}(E_0, \omega)$. $N_{QD}^{E1}(E_0, \omega)$ is shown in Fig. 7 along with the plane wave virtual photon spectrum. This alteration in the virtual photon spectrum is not adequate to eliminate the high energy rise in $\sigma(\gamma, p)$ and $\sigma(\gamma, \alpha)$ if only $E1$ virtual photon absorption is important. The high energy rise in the (e,p) cross section can be eliminated by multiplying the Levinger quenching factor, $e^{-L/\omega}$, in Eq. (7) by an arbitrary function of ω ; however, the same modification of the virtual photon spectrum of Eq. (7) will not eliminate the high energy rise in the (e,α) cross section but only reduce its magnitude slightly. Moreover, such modifications to the virtual photon spectra will not allow satisfactory simultaneous fits to be obtained to both our electroexcitation data and our bremsstrahlung data, measured between 50 and 100 MeV.

Thus we are left with the alternative that the cross section rise is produced by multiparticle emission following virtual photon absorption by the quasideuteron. The quantities we have called (e,p) and (e,α) cross sections²⁻⁴ are really photoproton and photoalpha yield cross sections and should be written as

$$\sigma_p = \sum_i i\sigma(\gamma, ip)$$

and

$$\sigma_\alpha = \sum_i i\sigma(\gamma, i\alpha).$$

(9)

FIG. 8. A comparison of the measured ⁹⁰Zr proton yield cross section, $\sum_i i\sigma(\gamma, ip)$, a histogram fit similar to that of Fig. 3, with the neutron yield cross section, $\sum_i i\sigma(\gamma, in)$, reconstructed from the data of Ref. 27.

In Fig. 8 we compare our ⁹⁰Zr photoproton yield cross section, the histogram of Fig. 3, with the photoneutron yield cross section,

$$\sigma_n = \sum_i i\sigma(\gamma, in), \quad (10)$$

reconstructed from the results of Ref. 27. The two curves come together for energies above 90 MeV suggesting that the multiplicities in the two channels are the same. In the range of 40–90 MeV some protons have low enough energies so that they cannot escape the Coulomb barrier and there neutron emission is more important. As Table I shows, the separation energies for multiple proton and alpha emission are not so high as to preclude the possibility of this explanation, especially if nucleon emission occurs predominantly in the preequilibrium stage.

The (γ, p) cross section for ⁹⁰Zr integrated to 30 MeV of 271 ± 20 MeV mb is consistent with the value of 235 ± 20 MeV mb obtained by Brajnik *et al.*,¹⁴ especially since their cross section does not contain all transitions and ours includes some $(\gamma, 2p)$ reactions. The ⁹²Zr (γ, p) cross section integrated to 30 MeV is depressed with respect to that of ⁹⁰Zr because more dipole strength goes into neutron emission as it does for any target with excess neutrons.

In our previous experiments^{2,4} the problem of multiple particle emission was certainly not so conspicuous, and except for the nucleus, ⁶⁴Zn, we obtained (γ, p) and (γ, α) cross sections which were flat at energies above the giant resonance. Here the yields seem to be dominated by multiparticle emission. It would be interesting to see if a precompound model calculation, analogous to the one reported by Blann *et al.*,²⁸ would reproduce these results.

On the experimental side it is possible and useful to study induced radioactivities, because characteristic gamma rays and half-lives serve to identify specific channels. For example, after a 130 MeV irradiation the gamma-ray spectrum from the ⁹⁰Zr target was examined, and the relative intensities of the gamma rays from ⁸⁸Y (1836 keV) and ⁸⁸Zr (393 keV) were determined. This comparison led to the conclusion that the ⁹⁰Zr $(\gamma, 2n)$ and (γ, pn) cross sections are comparable, lending support to our ideas

about the multiplicities. In the future it would be interesting to study the $^{90}\text{Zr}(e,3p)$ reaction through the induced radioactivity.

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Electrodisintegration of ^{63}Cu and ^{65}Cu

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The (e, α) cross section for ^{63}Cu has been measured in the electron energy range 14–34 MeV. The results have been analyzed using the distorted-wave Born approximation $E1$ and $E2$ virtual photon spectra and the $E1$ and $E2$ components of the corresponding (γ, α) cross section were obtained. To assess the accuracy of the virtual photon analysis, the $(e, 2n)$ cross section for ^{63}Cu was also measured and the obtained $(\gamma, 2n)$ cross section is compared with direct measurement of this cross section performed with annihilation gamma rays.

$$\left[\text{NUCLEAR REACTIONS } ^{65}\text{Cu}(e, \alpha) \text{ and } ^{63}\text{Cu}(e, 2n). \text{ Measured } \sigma_{e, \alpha}(E_0) \text{ and } \sigma_{e, 2n}(E_0). \text{ Deduced } \sigma_{\gamma, \alpha}^{E1}(E), \sigma_{\gamma, \alpha}^{E2}(E), \text{ and } \sigma_{\gamma, 2n}^{E1}(E). \right]$$

I. INTRODUCTION

The (e, α) cross section has been measured for several nuclei^{1,2} and the results have been analyzed using DWBA virtual photon spectra^{3,4} to obtain the $E1$ and $E2$ components of the corresponding (γ, α) cross section.

In this paper we use the same technique described in Refs. 1 and 2 to study the (e, α) cross section in ^{65}Cu . This technique takes advantage of the fact that the $E2$ virtual photon spectrum is enhanced relative to the $E1$ spectrum, as shown in Fig. 1, while the bremsstrahlung spectrum contains all multipoles in equal amounts. The $E1$ and $E2$ components of the (γ, α) cross section are obtained from combined measurements of the electrodisintegration cross section and the corresponding yield of electrodisintegration plus photodisintegration in-

duced by bremsstrahlung.

The (γ, α) cross section is known for many nuclei and in almost all studied cases it has a resonant shape. However, for two nuclei, ^{90}Zr (Ref. 5) and ^{65}Cu (Ref. 6), the (γ, α) cross section, obtained from measurements of the (e, α) cross section, showed a nonresonant behavior. For these nuclei the (γ, α) cross section increases continuously with the photon energy up to the maximum energy studied, which was ~ 60 MeV. An attempt was made to explain the strange behavior of the (γ, α) cross section in ^{90}Zr , using a preequilibrium exciton model combined with the quasideuteron model.⁵

This has motivated us to study the (e, α) cross section in ^{65}Cu , since it can easily be measured by residual activity. In Ref. 6 the (e, α) cross section was measured by directly counting the alpha particles emitted at 90° to the electron beam.

In order to assess the accuracy of the technique employed in the analysis we have measured the $(e, 2n)$ cross section for ^{63}Cu . There are several reasons for this choice. The $(\gamma, 2n)$ cross section is well known for this nucleus and it is above the isoscalar $E2$ and below the isovector $E2$ resonances. Thus we can be sure it is an $E1$ process. Even though there are many experimental tests of the $E1$ virtual photon spectra, and all show excellent agreement between calculation and experiment, it has been pointed out by Ströher⁷ that no conclusive tests of the $E1$ virtual photon spectra exist, since the investigated reaction data were analyzed with the assumption of a pure $E1$ excitation, in cases where $E2$ excitation could give a significant contribution. Furthermore, the $(e, 2n)$ cross section can

FIG. 1. $E1$ and $E2$ virtual photon spectra for 30 MeV electrons inelastically scattered from a copper nucleus.

TABLE I. Target properties and separation energies.

Target	^{63}Cu	^{65}Cu
Enrichment (%)	99.89	99.69
Thickness (mg/cm^2)	10.06 ± 0.10	9.86 ± 0.10
$S(2n)$, MeV	19.7	17.8
$S(\alpha)$, MeV	5.8	6.76

also be measured by radioactivity and yields the same gamma ray line used to measure the (e, α) cross section, as discussed in the next section. Consequently, the results also test the reliability of our measurements.

II. THE EXPERIMENT

The experiment was performed using the electron linear accelerator of Universidade de São Paulo. Table I gives target thicknesses and enrichments. The (e, α) and $(e, 2n)$ cross sections were measured by counting the induced residual activity. For the (e, α) cross section in ^{65}Cu we measured the 67.4 keV gamma ray which follows the decay of $^{61}\text{Co} \rightarrow ^{61}\text{Ni} + \beta^-$, with a half-life of (1.650 ± 0.005) h.⁸ For the $(e, 2n)$ cross section we measured the same 67.4 keV gamma ray, since ^{61}Cu decays to ^{61}Ni with a half-life of (3.41 ± 0.01) h.⁹ Figure 2 shows a typical gamma ray spectrum. It is important to notice that there are no nearby lines which could contribute to uncertainties in deriving the cross sections. As a check of our data, we measured both half-lives after irradiating the targets with 30 MeV electrons, and the decay constants obtained are compared with values from the literature in Table II.

FIG. 2. Typical gamma ray spectrum observed in the decay of ^{61}Co , obtained from the reaction $^{65}\text{Cu}(e, \alpha)^{61}\text{Co}$.

TABLE II. Decay constants.

Nucleus	λ (min^{-1})	Ref.
^{61}Co	$(6.95 \pm 0.29) \times 10^{-3}$ $(7.001 \pm 0.021) \times 10^{-3}$	This work 8
^{61}Cu	$(3.56 \pm 0.18) \times 10^{-3}$ $(3.39 \pm 0.01) \times 10^{-3}$	This work 9

In order to obtain the electrodisintegration plus photodisintegration yields, a $0.717 \text{ g}/\text{cm}^2$ copper radiator was placed in the electron beam, ahead of the target, without any spacing between the radiator and the target.

III. RESULTS AND ANALYSIS

The electrodisintegration cross section $\sigma_{e,x}(E_0)$ may be obtained from the photonuclear cross section $\sigma_{\gamma,x}^{\lambda L}(E)$ through an integral over the virtual photon spectrum $N^{\lambda L}(E_0, E, Z)$:

$$\sigma_{e,x}(E_0) = \int_0^{E_0 - m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) \times N^{\lambda L}(E_0, E, Z) \frac{dE}{E}. \quad (1)$$

In Eq. (1), E_0 stands for the total electron energy and E stands for the excitation energy of multipolarity λL . In the same spirit the yield with the radiator in is

FIG. 3. $\sigma_{e,2n}(E_0)$ for ^{63}Cu (circles) and the yield of electrodisintegration plus photodisintegration (squares). The smooth curves are the best fit to the data and were obtained by combining the histogram shown in Fig. 5 with the $E1$ virtual photon spectrum and the Davies-Bethe-Maximom (DBM) bremsstrahlung cross section in Eqs. (1) and (2).

FIG. 4. $\sigma_{e,\alpha}(E_0)$ for ^{63}Cu (circles) and the yield of electrodisintegration plus photodisintegration (squares). The smooth curves are the best fit to the data and were obtained by combining the $E1$ and $E2$ histograms of Fig. 7 with the virtual photon spectra and the DBM bremsstrahlung cross section in Eqs. (1) and (2). The triangles show the (e,α) cross section from Ref. 6.

$$Y_{e,x}(E_0) = \sigma_{e,x}(E_0 - 2\Delta E_0) + N_r \int_0^{E_0 - m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) \times K(E_0 - \Delta E_0, E, Z_r) \frac{dE}{E}, \quad (2)$$

where N_r is the number of nuclei/cm² in the copper radiator, $K(E_0, E, Z_r)$ is the bremsstrahlung cross section in copper, and ΔE_0 is the electron energy loss in half the radiator thickness. We have used the Davies-Bethe-Maximon¹⁰ bremsstrahlung cross section. The size correction discussed in Ref. 2 was

FIG. 5. $^{63}\text{Cu}(\gamma,2n)$ cross section. The histogram is the result derived from this work and the points show the measurement of Fultz *et al.*¹

applied to the virtual photon spectra. This correction, for 35 MeV electrons, amounts to 2% for the $E1$ and 8% for the $E2$ spectra.

The measured $(e,2n)$ and (e,α) cross sections (circles) and the corresponding yields of electrodisintegration plus photodisintegration (squares) are shown in Figs. 3 and 4. In Fig. 4 the triangles show the (e,α) cross section from Ref. 6. In order to compare it with our results, we have multiplied their 90° cross section by 4π , since previous measurements have shown that the angular distribution is nearly isotropic.^{11,12}

The cross sections $\sigma_{e,x}(E_0)$ and the yields $Y_{e,x}(E_0)$ have been simultaneously fitted, using Eqs. (1) and (2), with the photonuclear cross sections represented by histograms.

As discussed previously, for the $(\gamma,2n)$ cross section we used only $E1$ multipolarity. The histogram in Fig. 5 shows the $(\gamma,2n)$ cross section obtained from our measurements. The experimental points show the $(\gamma,2n)$ cross section measured by Fultz *et al.*¹³ There is good agreement between the $(\gamma,2n)$ cross section derived from our electrodisintegration measurement and the cross section measured with annihilation gamma rays. Our integrated $(\gamma,2n)$ cross section, up to 27.8 MeV (the maximum energy measured by Ref. 13) is 79.3 ± 2.2 MeVmb, while Fultz obtains 76 MeVmb. The difference of 4.5% in the integrated strength is well within the uncertainties of the absolute values of both experiments.

Figure 6 shows the ratio of measured to calculated $(e,2n)$ cross sections (circles) and measured to calculated electrodisintegration plus photodisintegration yields (triangles). This is the ratio between the measured points and the calculated full

FIG. 6. Ratio of measured to calculated $(e,2n)$ cross section (circles) and measured to calculated yield of electrodisintegration plus photodisintegration (triangles).

FIG. 7. $^{65}\text{Cu}(\gamma, \alpha)$ cross sections. The $E1$ and $E2$ histograms are the results from this work. The points show the (γ, α) cross section from Ref. 6.

curves of Fig. 3. This figure shows the compatibility between electrodisintegration and photodisintegration, that is, between the $E1$ virtual photon spectrum and the bremsstrahlung spectrum. These results give confidence to our measurement and the analysis employed to derive the photonuclear cross section.

Figure 7 shows the histograms obtained for the $E1$ and $E2$ components of the (γ, α) cross section. The $E1$ component exhausts 0.97 ± 0.14 percent of the dipole sum and the $E2$ component 3.2 ± 1.2 percent of the energy weighted sum rule for isoscalar $E2$. The points show the (γ, α) cross section from Ref. 6, multiplied by 4π . It was obtained by the unfolding of the (e, α) cross section assuming a pure $E1$ process, but this would introduce only a small error in the derived (γ, α) cross section. The strong

FIG. 8. Ratio of the (e, α) cross section from Ref. 6 and this work.

TABLE III. (γ, α) strength for nuclei in the $A = 60$ region. $E1$ sum: $60 \text{ NZ}/A \text{ MeV mb}$; $E2$ sum: $0.22Z^2 A^{-1/3} \mu\text{b/MeV}$.

Nucleus	$\int_0^{30} \sigma_{\gamma, \alpha}(E) dE$ (MeV mb)	Fraction of $E1$ sum	Fraction of $E2$ sum	Ref.
^{56}Fe	21 ± 3	2.1 ± 0.3	7 ± 1	2
^{58}Ni	43 ± 4	3.9 ± 0.4	21 ± 3	1
^{59}Co	17 ± 2	1.7 ± 0.2	5 ± 1	2
^{60}Ni	41 ± 4	3.5 ± 0.4	21 ± 5	1
^{62}Ni	17 ± 2	1.5 ± 0.2	8 ± 2	1
^{64}Zn	78 ± 16	6.9 ± 1.5	25 ± 3	2
^{65}Cu	10 ± 1	1.0 ± 0.1	3 ± 1	This work

disagreement between our (γ, α) cross section and that from Ref. 6 does not come from the analysis, but from disagreement in the experimental data. In Fig. 8 we show the ratio between the (e, α) cross section of Ref. 6 and the present measurement. Their cross section is only 0.13 of our value at 15 MeV and reaches 0.79 of our value at 34 MeV.

In Table III we compare the $E1$ and $E2$ strengths in the (γ, α) cross section of ^{65}Cu with other nuclei in this mass region. The (γ, α) cross section in ^{65}Cu is smaller than for other nuclei, but like the other nuclei studied, the $E2$ component exhausts a larger fraction of the $E2$ sum, relative to the $E1$. The strength of the (γ, α) cross section shows large variations in these nuclei, being 78 MeV mb for ^{64}Zn and only 10 MeV mb for ^{65}Cu . However, these large differences result mostly from differences in binding energies, Coulomb barrier heights, and competition with other channels, as already pointed out by Dodge *et al.*² The (γ, α) cross section in ^{65}Cu is similar in shape and magnitude to ^{62}Ni or ^{59}Co . The dominant statistical nature of the (γ, α) cross section in the nuclei listed in Table III is evident from the spectrum of the emitted alpha particles, which is of evaporation type, peaking at the energy of the Coulomb barrier.^{1,2} In the present work we did not observe the emitted alpha particles, but their spectrum is known from previous work.^{11,12} It has to be pointed out that in this respect the alpha particle spectrum from ^{90}Zr (Ref. 5) is also of the evaporation type, peaking at the energy of the Coulomb barrier height. Since, now, the only nucleus with a nonresonant (γ, α) cross section is ^{90}Zr , the results obtained for ^{65}Cu suggest that the ^{90}Zr (e, α) cross section should be measured again before developing calculations for reaction mechanisms involving cascade processes that could account for a nonresonant (γ, α) cross section.⁵

IV. CONCLUSIONS

The measurement of the $(e,2n)$ and $(e+\gamma,2n)$ cross sections for ^{63}Cu tests the $E1$ virtual photon spectrum in a situation where $E2$ contributions can be ruled out. The agreement between photodisintegration and electrodisintegration is very good. The $(\gamma,2n)$ cross section obtained from this measurement agrees well with the shape and absolute magnitude of available $(\gamma,2n)$ data.

The (e,α) cross section in ^{65}Cu is in disagreement with previous measurements in magnitude and shape. The (γ,α) cross section derived from our measurements has the expected resonant shape and fits well in the systematics of nuclei in this mass re-

gion. The $E1$ and $E2$ components of the (γ,α) cross section are both small but the $E2$ component is relatively more important.

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Virtual Photon Spectra for Finite Nuclei

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Abstract The experimental results of an isochromat of the virtual photon spectrum, obtained by measuring the number of ground-state protons emitted by the 16.28 MeV isobaric analogue state in ^{90}Zr as a function of electron incident energy in the range 17-105 MeV, are compared with the values predicted by a calculation of the E1 DWBA virtual photon spectra for finite nuclei. It is found that the calculations are in excellent agreement with the experimental results. The DWBA virtual photon spectra for finite nuclei for E2 and M1 multipoles are also assessed.

1. INTRODUCTION

The interest in electrodisintegration experiments to study the multipolarity of nuclear transitions has motivated an experiment to test virtual photon calculations with great accuracy. An isochromat of the E1 virtual photon spectrum was measured by counting the number of ground state protons, emitted by the 16.28 MeV, 1^- , isobaric analogue state in ^{90}Zr , as a function of incident electron energy in the range 17-105 MeV¹. When this experiment was performed, available virtual photon spectra were evaluated taking into account the distortion of incoming and outgoing electron waves in the Coulomb field of a point nucleus. It had already been shown that the plane wave approximation could only be used for very light nuclei and that the distorted wave Born approximation (DWBA) calculations² yielded spectra of correct magnitude as a function of Z . However, there were doubts about the range of validity of the point nucleus approximation. Since in electrodisintegration experiments the outgoing electrons are not detected, the measurements integrate over all possible values of momentum transfer. Because the electron scattering cross section is forward peaked, the electrodisintegration cross section is dominated by low momentum transfer and it was often argued that size corrections should be negligible. The isochromat experiment showed that for an E1 transition in a medium weight nucleus such as ^{90}Zr , size effects become evident for electrons of incident energy greater than 30 MeV. At that time two types of size corrections were available. One based on a plane wave calculation^{1,3} and also an evaluation

of virtual photon spectra for a finite nucleus, using a model for nuclear charge and current distributions, but performed in second order Born approximation (SOBA)⁴. It was found^{1,5} that both size corrections yielded acceptable results, but SOBA made a prediction closer to the experimental values. However SOBA is inadequate to describe heavy nuclei because these require a DWBA calculation. More recently Zamani-Noor and Onley⁶ have developed DWBA virtual photon calculations for finite nuclei. The comparison of their results with experimental data for the (e,n) cross section in ¹⁸¹Ta is discussed in ref.5. However the isochromat data of ref.1 is a more precise test for the virtual photon calculation. The purpose of this paper is to compare the experimental results of ref.1 with the E1 DWBA calculations for finite nuclei⁶. The results of these calculations for other multipoles will also be discussed.

2. THE VIRTUAL PHOTON METHOD

The virtual photon method relates electro- and photodisintegration through virtual photon spectra. In photodisintegration experiments, the photon is simply a means of injecting a certain amount of energy and angular momentum into the nucleus, which then has a variety of channels through which it may decay. Let us assume we are dealing with an Uranium nucleus and that it chooses the fission channel. We may draw diagrammatically the process of photofission as in fig. 1(a). The same nucleus can interact with a passing electron, which will create a time-varying field at the target site and this field may be regarded as a pulse of electromagnetic radiation. This is an old idea in classical electrodynamics due to Weiszacker⁷ and Williams⁸, but in quantum mechanics the intermediate radiation is a virtual photon and the process is depicted as in fig. 1(b). The nuclear part of the diagram is the same as in fig. 1(a). In both cases the interaction is electromagnetic, but there are important differences that force the nucleus to reveal more about itself, when interacting with an electron, than in a real photon interaction. Notice that the electron lines are curved because the electrons are moving in a Coulomb field which for a heavy nucleus can distort the electron wave function considerably.

Fig.1 - The photofission process (a) and the electrofission process (b) are shown diagrammatically.

The experimentalist has the choice of using real or virtual photons. Both originate with electrons, but real photons are created separately in a converter foil (see fig. 2) and travel a long path to the target. In this case the target experiences radiation with a broad spectrum which is known as the bremsstrahlung spectrum. On the other hand, a virtual photon is created and absorbed in interaction with one and the same nucleus, which is thus both converter and target. In a sense the target is, in this case, just closer to the source. If we make a classical analogy and think of the electron as an antenna, which is emitting electromagnetic radiation, there is the familiar difference between near field and far field. Far field radiation is purely transverse and has effectively a plane wave front; near field radiation is not planar and has also longitudinally polarized components. Near field or virtual radiation also has a different spectrum from bremsstrahlung and the shape and magnitude of this spectrum has to be known in order to analyze electrodisintegration measurements. Onley and his collaborators^{4,6,9,10} have generated the virtual photon spectra in the distorted wave approximation, first using a numerical calculation that treats the nucleus as a point^{9,10} and more recently for finite nuclei^{4,6}.

ELECTRODISINTEGRATION

$$\sigma_{e,x}(E_0) = \int_0^{E_0-m} \sum_{\lambda L} \sigma_{Y,x}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z) \frac{dE}{E}$$

ELECTRO + PHOTODISINTEGRATION

$$\gamma_{e,x}(E_0) = \sigma_{e,x}(E_0) + N_r \int_0^{E_0-m} \sigma_{Y,x}(E) K(E_0, E, Z) \frac{dE}{E}$$

Fig.2 - Schematic view of an electrodisintegration experiment (a) and photodisintegration experiment (b).

If the nucleus is subjected to a beam of radiation with spectrum $N(E)$ the resulting cross section is

$$\sigma = \int N(E) \sigma_\gamma(E) \frac{dE}{E} \quad (1)$$

where $\sigma_\gamma(E)$ is the photoexcitation cross section as a function of photon energy E ($\hbar = c = 1$) and $N(E)/E$ is the number of incident photons per unit energy interval. In the case of bremsstrahlung photons the spectrum depends on the electron incident energy E_0 , the atomic number Z of the converter or radiator and the number N_r of atomic nuclei/cm² in the radiator. In this case eq.(1) becomes

$$\sigma(E_0) = N_r \int_0^{E_0-m} \sigma_\gamma(E) K(E_0, E, Z) \frac{dE}{E} \quad (2)$$

where m is the electron rest energy, $K(E_0, E, Z)$ is the bremsstrahlung cross section and thus $N_r K(E_0, E, Z)/E$ is the number of photons per unit energy interval.

In the virtual photon case, assuming that the scattered electron is not detected (inclusive experiments like electrodisintegration) the cross section is rather similar to eq.(2).

$$\sigma(E_0) = \int_0^{E_0-m} \sum_{\lambda L} N^{\lambda L}(E_0, E, Z) \sigma_{\gamma}^{\lambda L}(E) \frac{dE}{E} \quad (3)$$

Here λL stands for the spin and parity or multipole class of the electromagnetic transition.

The basic difference is that the plane wave real photon spectrum has the same strength in all multipoles and in eq.(2):

$$\sigma_{\gamma}(E) = \sum_{\lambda L} \sigma_{\gamma}^{\lambda L}(E) \quad (4)$$

The virtual photon spectra by contrast, increase rapidly with L as illustrated in fig.3, which shows $N^{\lambda L}$ for $EL = E1, E2$, and $E3$ for $E_0 = 30$ MeV in tantalum. This enhancement has already been exploited in various experiments to examine multipole transitions of orders higher than $E1$, which a real photon spectrum does not reveal (e.g. references (11-22)). Combining electro- and photodisintegration measurements the experimentalist can change the multipole composition seen by the target.

b) Calculation of virtual photon spectra

In order to discuss the accuracy of the calculations we review them briefly.

For any system which interacts with the electromagnetic field the interaction may be written in the form

$$H_{int} = \int (\vec{J} \times \vec{A} - \rho \phi) d^3r \quad (5)$$

where $\vec{A}$, and ϕ are the potentials created by the electron. The nucleus is represented by the nuclear current density $\vec{J}$ and transition charge density ρ . These are constrained by continuity

Fig.3 - E1, E2 and E3 DWBA virtual photon spectra for a finite nucleus.

$$\vec{\nabla} \times \vec{J} + (d\rho/dt) = 0 \quad (6)$$

Since electrons are scattered by the Coulomb field of the nucleus the distorted wave (DWBA) formalism is used. Thus the static part of the Coulomb field is included in the unperturbed hamiltonian and hence H_{int} includes only the radiative part of the interaction.

Electron wave functions obey the Dirac equation and the DWBA formalism breaks the wave functions into partial waves. These are labelled with the Dirac angular momentum quantum number K , which may be a positive or negative integer and specifies both the total (j) and the orbital (ℓ) quantum numbers

$$j = |K| - 1/2 \quad \text{and} \quad \ell = |K + 1/2| - 1/2 \quad (7)$$

The basic probability amplitudes for the electron to change from angular momentum state K_0 , to state K_f , while emitting a photon in state λL are called

$R^{\lambda L}(K_0, K_f)$. In terms of these, the virtual photon spectra of eq.(3) can be obtained^(4,6,9,10)

$$N^{\lambda L}(E_0, E, Z) = \frac{\alpha K_f (E_0 + m)(E_f + m)E^4}{\pi K_0 (2L + 1)} \times \sum_{K_0 K_f} (2j_0 + 1)(2j_f + 2) |(j_0, j_f L; -1/2, 1/2) R^{\lambda L}(K_0, K_f)|^2 \quad (8)$$

Here E_f is the electron final energy and α is the fine structure constant.

In the point nucleus approximation the electron does not penetrate the nucleus and in the plane wave approximation the distortion of electron waves is ignored. Under both approximations the calculation of $N^{\lambda L}$ is relatively simple and one obtains analytical expressions for them⁴. However, the approximations will work for suitably low Z nuclei so that the distortion in electron waves is negligible. It can be used when the Coulomb energy is small compared with the incident electron energy:

$$\frac{\alpha Z}{R_N} \ll E_0 \quad (9)$$

In eq.(9) R_N is the nuclear radius. The assumption that the nuclear and electron waves do not interpenetrate significantly amounts to a condition

$$\bar{q} R_N \ll 1$$

where $\bar{q}$ is the average momentum transfer. How stringent this is depends on the process we are looking at, being least for magnetic and electric dipoles, for which the electron angular distribution exhibits a strong forward peak and for light nuclei (small R_N).

For medium and heavy nuclei conditions eqs. (9) and (10) are inevitably violated, and to evaluate $R^{\lambda L}$ in DWBA for a finite nucleus we have to assume a model to describe nuclear density functions for values of $r < R_N$, because $R^{\lambda L}$ involves an integral over all space (electron coordinates).

It has been shown⁴ that the results are not particularly sensitive to the details of the charge distribution. This should be expected since electron scattering at low

momentum transfer ($E_0 \lesssim 150 \text{ MeV}$) is unable to detect the details of the charge distribution and the only quantity that can be extracted from such experiments in the nuclear root-mean-square radius. In order to evaluate virtual photon spectra for finite nuclei, Onley and collaborators⁴ take the ground state charge distribution $\rho(r)$ to be the standard Fermi shape.

The transition probabilities $R^{\lambda L}$ contain the radial parts of the nuclear matrix elements, R_N^{EL} and R_N^{ML} , in which the nuclear current densities are expanded in multipoles (for details see ref.6). Thus

$$R_N^{EL} = \int_0^{R_N} \{ -j_{L+1}(E_0 r') |L/(L+1)|^1 \\ j_{L+1}(r') + j_{L-1}(E_0 r') J_{L-1}(r') \} r'^2 dr' \quad (11)$$

and

$$R_N^{ML} = \int_0^{R_N} j_L(E_0 r') J_L(r') r'^2 dr' \quad (12)$$

where j_L is the spherical Bessel function.

For the nuclear current densities Zamani-Noor and Onley use:

$$J_{L-1}(r) = r^{L-1} \rho_0(r) \quad (13)$$

$$J_L = d\rho_0/dr \quad (14)$$

and

$$J_{L+1} = 0 \quad (15)$$

which are derived under the assumption of irrotational incompressible flow in the nucleus.

The great difficulty in evaluating virtual photon spectra in DWBA is that the evaluation of $R^{\lambda L}$ requires to perform an integral over all space. Since the interaction involved is electromagnetic, and therefore long-ranged, there is no mechanism for cutting off this integral. To overcome this, Onley and his collaborators^{6,10} have developed an asymptotic series for the remote part of the integral and use numerical integration for the near part.

3. TESTS OF DWBA VIRTUAL PHOTON SPECTRA FOR FINITE NUCLEI

a) The E1 spectrum

In ref.1 the cross section $\sigma_{e,p_0}(E_0)$ for the proton decay of the 16.28 MeV, 1-, isobaric analogue state in ^{90}Zr was measured as a function of incident electron energy in the range 17-105 MeV. The same decay was also measured using bremsstrahlung photons.

The photonuclear absorption cross section integrated over the level width of the 16.28, 1-, isobaric analogue state in ^{90}Zr that results in protons populating the ground state of ^{89}Y is related to the photon width, Γ_γ , the ground state proton width, Γ_{p_0} , and the total width, Γ , of this level by

$$\int \sigma_{\gamma,p_0}(E) dE = (\pi\lambda)^2 [(2I+1)/(2I_0+1)] \Gamma_{p_0} \Gamma_\gamma / \Gamma \quad (16)$$

where I and I_0 are the spins of the excited and ground states.

Since

$$\sigma_{e,p_0}(E_0) = (1/E) |N^{E1}(E_0, E, Z)| \int \sigma_{\gamma,p_0}^{E1}(E) dE \quad (17)$$

using eqs.(16) and (17) with the experimental results for $\sigma_{e,p_0}(E_0)$ and calculating $N^{E1}(E_0, 16.28, 40)$, the invariant quantity $\Gamma_{p_0} \Gamma_\gamma / \Gamma$ can be obtained. By an analogous procedure, replacing in eq.(17) the number of virtual photons per unit energy interval by the appropriate number of real photons, the same invariant quantity $\Gamma_{p_0} \Gamma_\gamma / \Gamma$ can be obtained from the photoexcitation measurements. In ref.1 it was found that using DWBA E1 virtual photon spectra for a point nucleus, the value of this invariant quantity decreased as the electron energy increased, departing from that determined from the photoexcitation measurement, showing clearly that size corrections were necessary. When the virtual photon spectra were evaluated in the second order Born approximation, taking into account the finite size of the nucleus⁴, then the results obtained for $\Gamma_{p_0} \Gamma_\gamma / \Gamma$ were constant with the electron incident energy and compatible with that obtained from photodisintegration. The results obtained are

65.4 ± 0.6 eV from photodisintegration
 66.1 ± 0.3 eV from electrodisintegration
 using SOBA

Fig.4 - Ratio between the number of 16.28 MeV E1 virtual photons obtained with DWBA for a finite nucleus and that obtained with SOBA.

In fig.4 we show the ratio between the number of 16.28 MeV E1 photons obtained using DWBA for a finite nucleus⁶ and that obtained using SOBA. For $E_0 = 17.5$ MeV the number of virtual photons predicted by DWBA for a finite nucleus is 4 percent bigger than that predicted by SOBA. This can be understood because distortion is more important for lower values of the electron incident energy and indicates that the second order Born approximation does not take completely into account the distortion effects in ^{90}Zr . At 105 MeV SOBA underestimates size corrections by 2 percent.

Fig. 5 shows $\Gamma_{p_0}\Gamma_\gamma/\Gamma$ evaluated using SOBA and DWBA for a finite nucleus. In both cases $\Gamma_{p_0}\Gamma_\gamma/\Gamma$ is constant and independent of the electron incident energy. The dashed lines represent in each case the weighted average value, which are:

65.2 ± 0.3 with DWBA for a finite nucleus and 66.1 ± 0.3 eV for SOBA. Compared to 65.4 ± 0.6 eV obtained from photodisintegration, the DWBA result is in better agreement.

Fig.5 - The invariant quantity evaluated using SOBA and DWBA for a finite nucleus. The dashed lines represent in each case the weighted average value.

We conclude that the E1 virtual photon spectra evaluated in DWBA for a finite nucleus is in excellent agreement with experiment.

b) The E2 spectrum

There is no known E2 isolated level in medium weight or heavy nuclei which decays by particle emission. Thus an isochromat of the E2 spectrum cannot be measured.

In ref. 5 the electrodisintegration of ^{181}Ta by one neutron emission was used to assess the reliability of the recent DWBA calculations for a finite nucleus. In ^{181}Ta the isoscalar E2 resonance decays dominantly by one neutron emission because charged particle decay is inhibited by the Coulomb barrier. Thus the (e, n) cross section contains basically E1 and E2 excitations. Figs. 6 and 7 show E1 and E2 DWBA virtual photon spectra for a point and finite nucleus ($Z = 73$).

Size corrections affect the E2 spectrum far more. It was shown⁵ that the use of virtual photon spectra for a point nucleus made electro- and photodisintegration incompatible in heavy nuclei, but those for a finite nucleus yielded good agreement with experiment. The DWBA calculations for a finite nucleus⁶ make electrodisintegration and photodisintegration compatible and the E2 strength necessary to simultaneously fit both cross sections is in good agreement with the known systematics.

Fig.6 - DWBA E1 virtual photon spectra for a point and a finite nucleus.

This test does not have the same accuracy as the measurement of the E1 isochromat, but it indicates that the E2 size correction has the appropriate magnitude.

The M1 Spectrum

Fig. 8 shows DWBA M1 spectra for a point and a finite nucleus ($Z = 73$). The magnitude of the size correction is much bigger for M1 than for E2 (see fig.

Fig.7 - DWBA E2 virtual photon spectra for a point and a finite nucleus.

7). This is surprising if one reasons along lines suggested by plane wave (PWBA) calculations. In PWBA size effects are small for the transverse components of the virtual photon spectra³ and this behaviour is reproduced by DWBA, at least for E1 and E2 spectra. The PWBA E1 spectrum is dominantly transverse and the magnitude of its size corrections in DWBA is relatively small as compared to that of the E2 spectrum, which is dominantly longitudinal^{3,5}. Since the PWBA M1 spectrum is purely transverse, size corrections should have little effect on the corresponding DWBA virtual photon spectrum.

In a recent paper on ¹⁹⁷Au, Campos *et al.*²⁴ derived the M1 and E2 strengths observed on the one neutron decay channel of the photonuclear cross sections. They used both finite and point DWBA M1 spectra in their analysis and got consistent results only with the point DWBA M1 spectrum. It must be emphasized that both E1 and E2 spectra used in the analysis were finite DWBA calculations. The M1 cross section obtained using finite DWBA M1 spectra was too big (larger

Fig.8 - DWBA M1 virtual photon spectra for a point and a finite nucleus.

than the absorption in the energy region studied), probably due to the excessive correction on the spectrum, which made it too small. The results obtained for M1 strength using point DWBA M1 spectra were compatible with previous experimental results and with theoretical estimates and suggest that the spectra used in the analysis have the correct magnitude.

4. CONCLUSIONS

The E1 virtual photon spectra evaluated in the distorted wave Born approximation for a finite nucleus are in excellent agreement with experimental results.

Fig.9 - Isochromats of the DWBA M1 virtual photon spectra for a point and a finite nucleus. The virtual photon energy is 7 MeV and the target nucleus is $Z=73$.

The E2 spectra are also in good agreement with experimental results, but the experimental test in this case does not have the same accuracy as that for the E1.

For magnetic multipoles the results of the calculations for a finite nucleus seem to be unreliable. They predict a very large correction due to the finite size of the nucleus, while these corrections are expected to be very small as a consequence of the transverse nature of these multipoles. The failure of the calculations for a finite nucleus may result from inadequate model description of the nuclear currents. Eventhough this strongly affects the magnetic multipoles, it has little effect on the electric multipoles. For the latter it is found^{4,6} that the results are rather independent of the detailed shape of the charge distribution.

They depend primarily on the values of the nuclear root-mean-square radius and the appropriate transition radius.

For magnetic multipoles the most reliable calculation available present at is the DWBA for a point nucleus. For the M1 multipole it yields $B(M1)$ strengths, derived from electrodisintegration measurements, which are consistent with results obtained with other probes, like polarized photon scattering.

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Resumo

Os resultados experimentais referentes à medida de uma isocromata do espectro de fótons virtuais, a qual foi efetuada detectando-se o número de prótons que decaem para o estado fundamental, emitidos pelo estado análogo de energia 16.28 MeV no ^{90}Zr , em função da energia do elétron incidente na faixa de 17-105 MeV, são comparados com os valores previstos pelo cálculo do espectro de fótons virtuais E1, efetuado na aproximação DWBA e considerando o tamanho finito do núcleo. Verifica-se que os cálculos estão em excelente acordo com os resultados experimentais. Os cálculos de espectros de fótons virtuais, em DWBA e para núcleos finitos, para multipolos E2 e M1 também são discutidos.

Nuclear size effects in virtual photon spectra

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Abstract. The (e, n) cross section has been measured for ^{181}Ta from threshold to 30 MeV. The results are compared with the cross sections predicted by virtual photon theory using different corrections for nuclear size effects.

NUCLEAR REACTIONS $^{181}\text{Ta}(e, n)$, $E = 9\text{--}30$ MeV; measured $\sigma(E)$; virtual photon analysis; deduced E2 strength.

1. Introduction

During the last decade the study of the giant multipole resonances together with the emergence of virtual photon calculations using the distorted-wave Born approximation (DWBA) of Onley, Wright and collaborators [1, 2] have jointly stimulated a renaissance of interest in electrodisintegration experiments. Early calculations of the M1, E2 and E3 virtual photon spectra in the plane-wave Born approximation (PWBA) showed that these spectra were greatly enhanced relative to the E1 spectrum. The enhancement was understood to result from the momentum dependence of higher multipole cross sections and the fact that in the electron scattering event the momentum transfer to the nucleus can greatly exceed the energy transfer. Experimental results [3–5] showed that these PWBA calculations were grossly underestimated for all but the lightest nuclei.

Gargaro and Onley [1] were the first to calculate the virtual photon spectra taking into account the distortion of the incoming and outgoing electron waves in the Coulomb field of a point nucleus. These calculations showed that the enhancement of higher multipole virtual photon spectra, already seen in PWBA calculations, was greatly increased as the atomic number of the nucleus increased.

The enhancement of the E2 relative to the E1 DWBA virtual photon spectra has recently been exploited in several electrodisintegration experiments to study E2 absorption by nuclei, e.g. references [6–14]. In these experiments the relationship between the electrodisintegration cross section $\sigma_{e,x}(E_0)$ and the photonuclear cross section $\sigma_{\gamma,x}(E)$ is used:

$$\sigma_{e,x}(E_0) = \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z) dE/E. \quad (1)$$

Here $N^{\lambda L}(E_0, E, Z)/E$ stands for the number of virtual photons of multipolarity λL exchanged with a nucleus of atomic number Z when an electron of energy E_0 produces an excitation of energy E .

Since the DWBA virtual photon spectra used in equation (1) are evaluated for a point nucleus, it is assumed that the transition probabilities for photodisintegration and electrodisintegration, $B(EL, E)$ and $B(EL, q)$, are equal. Here q is the momentum transferred to the nucleus in electrodisintegration. This assumption is exact as $q \rightarrow E$ but surely fails for $q \gg E$.

Several experimental checks of the DWBA virtual photon spectra have been made [3–5] but the most complete and accurate of these tests was an experiment performed at the National Bureau of Standards (NBS) [15]. An isochromat of the E1 virtual photon spectrum was measured by counting the number of ground-state protons emitted by the 16.28 MeV isobaric analogue state in ^{90}Zr as a function of incident electron energy in the range 17–105 MeV. The experimental results reproduce well the DWBA spectra for a point Zr nucleus for electron energies up to 30 MeV. A radiator was used for electron energies of 60–100 MeV to measure the photodisintegration cross section for the 16.28 MeV level. These results showed that the Davies–Bethe–Maximon bremsstrahlung cross section [16] yields the same result for $\Gamma_\gamma \Gamma_{p_0} / \Gamma$ as the electrodisintegration below 30 MeV, where size corrections for the finite extent of the Zr nucleus are minimal. As E_0 increases to 105 MeV the need for such corrections became manifest. It was shown that the size corrections of Dodge *et al* [6] and of Durgapal and Onley [17] made the electrodisintegration results compatible with the photodisintegration results, for the full range of electron incident energies.

The size correction of Dodge *et al* [6] consists in multiplying $N^{\lambda L}(E_0, E, Z)$ inside the integral of equation (1) by

$$F^{\lambda L}(qR) = \left[\left(\frac{E}{q} \right)^L \frac{j_L(qR)}{j_L(ER)} \right]^2_{q=q_{\text{RMS}}} \quad (2)$$

where R is the radius of the nuclear ground state and q is the momentum transfer. The RMS values of q are obtained in PWBA as described in reference [6]. This size correction is of the same magnitude as the one proposed by Shottner [18], also based on a PWBA calculation. For the electric dipole transition at 16.28 MeV in the ^{90}Zr nucleus $F^{\lambda L}(qR)$ is 0.89 for 100 MeV electrons and is negligible for electrons of less than 30 MeV.

Durgapal and Onley [17] take into account the size effects by evaluating the virtual photon spectra in second-order Born approximation (SOBA) using a model for the nuclear charge and current density distributions. Their model reproduces values of the charge and transition radii obtained from electron scattering experiments. The resulting spectra are slightly smaller than those obtained with the use of equation (2).

The results obtained in the NBS experiment for $\Gamma_\gamma \Gamma_{p_0} / \Gamma$ are:

65.4 ± 0.6 eV	from photodisintegration
66.1 ± 0.3 eV	from electrodisintegration using SOBA
63.8 ± 0.3 eV	from electrodisintegration using equation (2).

Even though both size corrections yield acceptable results, the result obtained with SOBA is closer to the photodisintegration result. The SOBA virtual photon calculation takes into account not only the fact that $B(EL, E)/B(EL, q)$ is not equal to one, but also deviations from a point nuclear Coulomb field. Thus, in principle, results obtained with SOBA should be the most reliable.

However, SOBA cannot be used for heavy nuclei because these require a full DWBA calculation. Zamani-Noor and Onley [19] have recently developed a DWBA calculation for the virtual photon spectra with size effects, using charge and current density distributions given by the Tassie hydrodynamical model [20]. The results of this calculation, which are

Figure 1. DWBA E1 virtual photon spectra for $E_0 = 12, 22$ and 30 MeV for a point and a finite nucleus.

Figure 2. DWBA E2 virtual photon spectra for $E_0 = 12, 22$ and 30 MeV for a point and a finite nucleus.

in agreement with the results from SOBA for Zr, show that for heavy nuclei size effects are important even at low electron incident energies (< 30 MeV).

Figures 1 and 2 show E1 and E2 virtual photon spectra for $E_0 = 12, 22$ and 30 MeV and $Z = 73$. In these figures the full curves represent the DWBA spectra for a point nucleus. The broken curves represent the DWBA spectra with size effects [19]. For E1 the size effects decrease the spectra by $\sim 10\%$ and for E2 by $\sim 50\%$.

The size correction based on PWBA calculations and given by equation (2) decreases the E1 spectra by 2% and the E2 spectra by 7% for $E_0 = 30$ MeV and virtual photons of 10 MeV. The correction becomes smaller for lower electron energies and also for higher virtual photon energies. These size corrections are much smaller than those shown in figures 1 and 2.

In order to test if the magnitude of the size corrections in the recent calculation of Onley and collaborators [19], which we will designate by DWBA-SZ, are in agreement with experimental results we have measured the electrodisintegration of ^{181}Ta .

2. The experiment

The $^{181}\text{Ta}(e, n)$ cross section was measured by residual activity, following the 93.3 keV γ ray line that results from the decay of ^{180}Ta to ^{180}Hf , which has a half-life of 8.11 ± 0.02 h. We have used seven targets obtained from two foils which were ~ 4 and ~ 10 mg cm^{-2} thick. The target thicknesses, determined by weighing, are listed in table 1. The targets were activated in the electron beam of the São Paulo Electron Linear Accelerator. The LINAC current was measured using a Faraday cup.

The measured cross section was corrected for the photodisintegration produced by bremsstrahlung in the targets. This correction was between 1 and 4% depending on the electron incident energy and the target thickness. For electron kinetic energies of 24.5, 27.5 and 30.5 MeV we have also measured the photodisintegration induced by bremsstrahlung, using a 0.329 g cm^{-2} copper radiator, placed in the electron beam just ahead of the target. The LINAC current was measured using a secondary emission monitor calibrated relative to the Faraday cup. The measured yields were corrected for the energy loss of electrons in the radiator and the electrodisintegration produced in the targets was subtracted.

The results obtained for the electrodisintegration cross section are shown in figure 3. Each of the data points is the weighted mean of several measurements with different targets. The uncertainty in the absolute scale is 10%. The results obtained for the photodisintegration yield are given in table 2.

Table 1. Target thickness.

Target	Thickness (mg cm^{-2})
1	3.82 ± 0.12
2	3.81 ± 0.11
3	3.84 ± 0.12
4	10.19 ± 0.30
5	10.23 ± 0.31
6	10.24 ± 0.31
7	10.34 ± 0.31

Figure 3. Measured electrodisintegration cross section against the electron kinetic energy. The uncertainty in the absolute scale is 10%. Each data point is the weighted mean of several measurements. The full curve is a fit of the data, with E1 and E2 components in the (γ, n) cross section, using DWBA virtual photon spectra for a finite nucleus.

3. Discussion

In order to compare our results with the predictions of virtual photon theory we have to discuss which multipoles should enter in evaluating equation (1). Besides the dominant E1 in the energy range covered by this experiment there are several electric multipoles: the isoscalar E2 at $63A^{-1/3}$ MeV, the isoscalar E3 at $110A^{-1/3}$ MeV and the isovector E2 at

Table 2. Photodisintegration yield obtained with a 0.329 g cm^{-2} radiator.

Number of measurements	Electron energy (MeV)	$\sigma_{\text{Br},n}$ (mb)
6	24.5	2.74 ± 0.06
4	27.5	2.61 ± 0.06
6	30.5	3.01 ± 0.06

Table 3. Integrated photonuclear cross sections.

λL	T	E_p (MeV)	$\int \sigma_{\gamma}^{EL}(E) dE$ (MeV mb)
E1	1	$80 A^{-1/3}$	2613
E2	0	$63 A^{-1/3}$	27
E2	1	$130 A^{-1/3}$	162
E3	0	$110 A^{-1/3}$	1

$130A^{-1/3}$ MeV. Table 3 gives the integrated photonuclear cross sections corresponding to one energy-weighted sum for each of these resonances. We have made the approximation

$$\int \sigma_{\gamma}^{EL}(E) dE = E_p^{2L-2} \int \frac{\sigma_{\gamma}^{EL}(E) dE}{E^{2L-2}} \quad (3)$$

where E_p is the peak position of the resonance.

Figure 4. E1, E2 and E3 DWBA virtual photon spectra for a finite nucleus.

Table 4. Isoscalar E2 resonance.

Reaction	E_p (MeV)	Γ (MeV)	% EWSR	Reference
(p, p')	11.2 ± 0.2	3.9	42 ± 8	[22]
(e, e')	9.54 ± 0.20	2.07 ± 0.35	8 ± 4	[23]
	11.47 ± 0.22	3.13 ± 0.55	22 ± 8	
(e, e')	9.5 ± 0.2	1.8 ± 0.6	29 ± 5	[24]
	11.3 ± 0.2	2.2 ± 0.7	63 ± 8	
	12.6 ± 0.2	1.3 ± 0.8	11 ± 5	

Figure 4 shows DWBA—SZ E1, E2 and E3 virtual photon spectra for 30 MeV electrons. Because the photonuclear cross section corresponding to one E3 sum is so much smaller than the E1 it makes negligible contribution to our measured cross section, even though the E3 spectrum is enhanced relatively to the E1. Since electrons of the energies used here cannot excite the monopole giant resonance [21] we have to consider only E1 and E2 electric multipoles in equation (1).

The isoscalar E2 giant resonance in ^{181}Ta has been studied in (p, p') and (e, e') experiments. Table 4 summarises the available results and suggests that these results are not in agreement as to the shape and strength of this resonance. Figure 5(a) shows the

Figure 5. (a) Experimental results for the isoscalar E2 resonance in ^{181}Ta . (b) The results shown in (a) are normalised to the same peak height.

photonuclear E2 cross section corresponding to the parameters given in table 4. In figure 5(b) the results of references [22] and [23] are normalised to the peak height of reference [24] and it shows clearly that the three experiments are in good agreement as to the shape and position of the isoscalar E2 resonance. They disagree only in the strength, but this is a consequence of the different models used to extract it.

The isovector E2 was measured for ^{181}Ta by Miura *et al* [24] and Hicks *et al* [23]. Figure 6 shows the predicted electrodisintegration yields for E1, isoscalar E2 and isovector E2. The E1 component was evaluated using the (γ, n) data from Livermore [25] in equation (1). For the isoscalar E2 we used a Lorentz line shape that reproduces the shape shown in figure 5(b) and exhausts one sum. For the isovector E2 we used a Lorentz line shape with $E_p = 23.2$ MeV and $\Gamma = 7$ MeV, in accordance with reference [23] with a strength corresponding to one sum to represent the E2 absorption, σ_{abs}^{E2} . However, since the

Figure 6. The contributions of various multipoles to the electrodisintegration cross section (see text for details). The insert shows the ratio of calculated yields for $(E1 + M1)$ and $(E1 + E2)$ to pure E1.

isovector E2 is in a region where the dominant decay channel is $(\gamma, 2n)$ we assumed that

$$\sigma_{\gamma, n}^{E2} = \sigma_{\text{abs}}^{E2} \frac{\Gamma_n}{\Gamma_{\text{tot}}} = \sigma_{\text{abs}}^{E2} \frac{\sigma_{\gamma, n}}{\sigma_{\gamma, n} + \sigma_{\gamma, 2n}}. \quad (4)$$

Since we are only interested in showing the relative magnitudes of the various multipoles the yields shown in figure 6 were evaluated using DWBA virtual photon spectra for a point nucleus to speed up computing.

We have now to discuss the contributions from magnetic multipoles. Since there is no systematics for M1 and M2 absorption in heavy nuclei we assumed an M1 strength of $20 \mu_0^2$ uniformly distributed between 7.6 and 8.6 MeV and for M2 we assumed the strength observed for ^{208}Pb [26]. The predicted yields are also shown in figure 6. It is clear that M2 can be neglected relative to E1. The M1 can make a detectable contribution if a detailed measurement is performed in the threshold region. In the insert of figure 6 we illustrate the expected increase in the measured electrodisintegration cross section due to an M1 strength of $20 \mu_0^2$ or to one E2 isoscalar sum, plotting the ratio of the predicted yields to a pure E1. The ratios shown were evaluated using DWBA-SZ.

From the above discussion we are led to take into account contributions from E1, E2 and M1 in the analysis of our experiment.

The measured (γ, n) cross section [25] contains all multipoles but does not distinguish them. We can write

$$\sigma_{\gamma, n}(E) = \sigma_{\gamma, n}^{E1}(E) + \sigma_{\gamma, n}^{E2}(E) + \sigma_{\gamma, n}^{M1}(E) \quad (5)$$

and

$$\sigma_{\gamma, n}^{E1}(E) = \sigma_{\gamma, n}(E) - \sigma_{\gamma, n}^{E2}(E) - \sigma_{\gamma, n}^{M1}(E). \quad (6)$$

Inserting (6) in (1) we write

$$\begin{aligned} \sigma_{e, n}(E_0) = & k_1 \int_0^{E_0 - m} \sigma_{\gamma, n}(E) N^{E1}(E_0, E, Z) dE/E \\ & + k_2 \int_0^{E_0 - m} \sigma_{\gamma, n}^{E2}(E) [N^{E2}(E_0, E, Z) - N^{E1}(E_0, E, Z)] dE/E \\ & + k_3 \int_0^{E_0 - m} \sigma_{\gamma, n}^{M1}(E) [N^{M1}(E_0, E, Z) - N^{E1}(E_0, E, Z)] dE/E \end{aligned} \quad (7)$$

where for $\sigma_{\gamma, n}(E)$ we use the measured (γ, n) cross section [25] and k_i are constants to be determined from the fit of equation (7) to our data.

The constant k_1 allows for a difference between our absolute scale and that of the (γ, n) cross section. This constant can also be determined, independently, by comparing our measured bremsstrahlung yield given in table 3 with the calculated values

$$\sigma_{\text{Br}, n} = N_r \int_0^{E_0 - m} \sigma_{\gamma, n}(E) K(E_0, E, Z) dE/E \quad (8)$$

where $\sigma_{\gamma, n}$ is the photonuclear cross section [25], K is the Davies-Bethe-Maximon bremsstrahlung cross section, which shows excellent agreement with photodisintegration experimental results [16] and N_r is the number of nuclei/cm² in our radiator. Using equation (8), k_1 is the ratio between measured and calculated yields. From our measurements we obtain $k_1 = 1.29 \pm 0.04$.

For $\sigma_{\gamma,n}^{E2}$ isoscalar and isovector we have used the shapes previously discussed and used to evaluate the yields of figure 6. Thus, by fitting equation (7) to our data, k_2 will give the fraction of the E2 isovector and isoscalar sums.

For $\sigma_{\gamma,n}^{M1}$ we assumed a strength of $1\mu_0^2$ uniformly distributed between 7.6 and 8.6 MeV excitation energy, since the M1, if it is located above the (γ, n) threshold, will be distributed into many levels and our measurement is not sensitive to the detailed structure. In equation (7) k_3 will give the M1 strength in units of μ_0^2 .

We fitted our electrodisintegration data with equation (7) using the DWBA virtual photon spectra corrected for size effects with equation (2) and obtained a good fit with $k_1 = 1.16 \pm 0.02$, an E2 intensity of $(70 \pm 10)\%$ of the E2 sum and an M1 intensity of $(17 \pm 5)\mu_0^2$. The quality of the fit was good and the results would seem quite reasonable. However, this result is unacceptable, since the value of k_1 obtained from the fit is

Figure 7. Ratio of measured to calculated electrodisintegration cross section. (a) Using DWBA virtual photon spectra corrected for size effects with equation (2) and assuming that the (γ, n) cross section is pure E1. (b) Using DWBA virtual photon spectra for a finite nucleus and assuming that the (γ, n) cross section is pure E1. (c) Using DWBA virtual photon spectra for a finite nucleus and assuming that the (γ, n) cross section contains contributions from both E1 and E2.

incompatible with the value determined from our photodisintegration measurements. This is an important point and could explain why previous results of electrofission measurements, analysed with DWBA virtual photon spectra for a point nucleus yielded conflicting results. (See for example references [12–14].) A measurement of the photodisintegration yield, under the same experimental conditions as the electrodisintegration, is a necessary constraint.

It is impossible to fit our electrodisintegration data using $k_1 = 1.29 \pm 0.04$ and the DWBA virtual photon spectra for a point nucleus or corrected for size effects by equation (2). Figure 7(a) shows the ratio of measured to calculated electrodisintegration for E1 only ($k_1 = 1.29$, $k_2 = k_3 = 0$) (using the size corrections of equation (2)). Since the ratio becomes smaller than one, it would require negative E2 and/or M1 strength to fit the data. The conclusion is that with this size correction photodisintegration and electrodisintegration are incompatible. Figure 7(b) shows the same ratio evaluated with DWBA-SZ. The ratio is greater than one, indicating that other multipoles contribute.

We fitted our data using DWBA-SZ for the virtual photon spectra in equation (7). Since it turned out that k_3 was compatible with zero, we fitted the data only in terms of E1 and E2. The full curve in figure 3 is the result of this fit. The results obtained are: $k_1 = 1.31 \pm 0.02$ (in good agreement with the value determined by our photodisintegration measurements) and an E2 strength of $140 \pm 37\%$ of the isoscalar E2 sum and $31 \pm 8\%$ of the isovector E2 sum in the (γ, n) channel. Figure 7(c) shows the ratio of measured to calculated yield. The agreement is excellent.

The DWBA calculation for a finite nucleus from Zamani-Noor and Onley [19] makes electrodisintegration and photodisintegration yields compatible since it yields the same value for k_1 as obtained, independently, from our photodisintegration measurements. The isoscalar E2 intensity is consistent with the systematics for heavy nuclei [27, 28]. The uncertainty in the E2 strength could be decreased by extending the measurement to higher electron incident energies.

The fact that the present measurement is consistent with the absence of M1 does not exclude the possibility of M1 strength near threshold. Actually in the ratio shown in figure 7(c) the first point is high as expected when there is M1. In order to be more conclusive about the M1 component we would need to measure several points from threshold to 9 MeV with good statistics.

4. Conclusions

Our results indicate that the recent calculation of DWBA virtual photon spectra for a finite nucleus is in agreement with experimental results. With these calculations photodisintegration and electrodisintegration results are compatible. For heavy nuclei the difference between the spectra for a finite and point nucleus is large for the E2 spectra even for electron energies around 10 MeV. This invalidates the multipole decomposition obtained in several previous experiments of electrodisintegration and electrofission in heavy nuclei.

Since the analysis of electrodisintegration experiments in terms of virtual photon spectra involves, usually, a photonuclear cross section measured in an independent experiment, reliable results can only be obtained if the possibility of differences between the absolute scales of the electrodisintegration and photodisintegration cross sections are taken into account.

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M1 strength in ^{197}Au

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Abstract. The electrodisintegration of ^{197}Au was measured by residual activity from threshold to 23 MeV. The results are analysed with DWBA virtual photon spectra for finite nuclei (E1 and E2 multipoles, point DWBA for M1) to obtain M1 and E2 strengths.

NUCLEAR REACTIONS: $^{197}\text{Au}(e, n)$; measured $\sigma_{e,n}(E)$; $E=9\text{--}23$ MeV; virtual photon analysis deduced E2 and M1 strengths.

1. Introduction

There is little information about M1 strengths in heavy nuclei. Light- and medium-weight nuclei have been extensively studied by electron scattering at backward angles [1–3]. For heavy nuclei, this technique presents difficulties, as discussed in [4], because nuclear size and Coulomb distortion combine to shift the peak of the form factor first maximum to low momentum transfers ($\sim 0.1 \text{ fm}^{-1}$). Since existing accelerators and spectrometers are only able to reach the second maximum, where the cross section is greatly reduced, extrapolation to the photon point becomes strongly model dependent and $B(\text{M1}) \uparrow$ cannot be obtained with reasonable accuracy, even when its intensity is measurable.

However, these difficulties do not occur in electrodisintegration experiments, because they are integrated over all angles of scattered electrons and thus dominated by low momentum transfer. The recent calculations of virtual photon spectra for finite nuclei [5, 6] made possible the interpretation of electrodisintegration cross sections, $\sigma_{e,x}$, for heavy nuclei in terms of the multipole composition of the corresponding photonuclear cross section, $\sigma_{\gamma,x}^{\lambda L}$. The relationship between these cross sections is

$$\sigma_{e,x}(E_0) = \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z) \frac{dE}{E}. \quad (1)$$

Here $N^{\lambda L}(E_0, E, Z)/E$ stands for the number of virtual photons of multipolarity λL exchanged with a nucleus of atomic number Z when an electron of energy E_0 produces an excitation of energy E . The validity of equation (1) to relate photo and electrodisintegration has been verified in several experiments [6, 7]. In particular, the recent calculations of virtual photon spectra for finite nuclei [5] have been compared with experimental results by Wolyneć *et al* [7].

For the electron energies used in this work, as discussed in [7], the only multipoles that make a detectable contribution to the electrodisintegration cross section are E1, M1 and E2. Even though the E1 contribution is dominant, differences in shape, magnitude and in the dependence with incident electron energy between virtual photon spectra of distinct multipolarities allow the separation of the multipole components in the photonuclear process.

The sensitivity of the (e, n) cross section to E2 and M1 contributions is illustrated here by comparing the calculated cross section for a pure E1 absorption with that of an (E1 + E2 + M1) absorption process in ^{197}Au . For this purpose equation (1) can be rewritten:

$$\begin{aligned}\sigma_{e,n}(E_0) = & \int_0^{E_0-m} \sigma_{\gamma,n}(E) N^{E1}(E_0, E, Z) \frac{dE}{E} \\ & + \int_0^{E_0-m} \sigma_{\gamma,n}^{M1}(E) [N^{M1}(E_0, E, Z) - N^{E1}(E_0, E, Z)] \frac{dE}{E} \\ & + \int_0^{E_0-m} \sigma_{\gamma,n}^{E2}(E) [N^{E2}(E_0, E, Z) - N^{E1}(E_0, E, Z)] \frac{dE}{E}\end{aligned}\quad (2)$$

where we have assumed that the photonuclear cross section $\sigma_{\gamma,n}(E)$ has only E1, E2 and M1 components:

$$\sigma_{\gamma,n}(E) = \sigma_{\gamma,n}^{E1}(E) + \sigma_{\gamma,n}^{E2}(E) + \sigma_{\gamma,n}^{M1}(E)$$

or

$$\sigma_{\gamma,n}^{E1}(E) = \sigma_{\gamma,n}(E) - \sigma_{\gamma,n}^{M1}(E) - \sigma_{\gamma,n}^{E2}(E). \quad (3)$$

For $\sigma_{\gamma,n}(E)$ we use experimental results [8]. It has to be pointed out that the experimental photonuclear cross section contains all multipoles, but the results cannot distinguish them because their contribution is very small compared with the dominant E1 component. $\sigma_{\gamma,n}^{E2}(E)$ was represented by a Lorentz shape with peak energy $E_p = 10.8$ MeV and width 4 MeV, exhausting one isoscalar E2 EWSR. The isovector E2 resonance located at ~ 23 MeV makes a negligible contribution to the (e, n) process in the energy range covered by our experiment. It is well known that in heavy nuclei the M1 strength is located in the vicinity of the neutron emission threshold and is distributed into many levels. Since an electrodisintegration cross section is sensitive to the position and integrated strength, but not to the detailed structure, we assumed an M1 strength uniformly distributed between 8.1 and 9.4 MeV, with an integrated strength of $30\mu_0^2$. This strength lies between the $50\mu_0^2$ obtained by Lone *et al* [1, 9] from an (e, e') experiment at 180° and the $20\mu_0^2$ estimated by Bohr and Mottelson [10]; these are the only available results for M1 in ^{197}Au .

Figure 1 shows the ratio of the calculated (e, n) cross section assuming E1, E2 and M1 absorption, to that obtained for a pure E1 absorption process. An M1 strength of $30\mu_0^2$ makes a sizeable contribution to the electrodisintegration cross section from threshold up to ~ 20 MeV, while the contribution of the isoscalar E2 component is rather small.

In order to be more sensitive to the isoscalar E2 component the experiment has to go to higher incident electron energies [7].

Figure 1. Predicted ratio of the electrodisintegration cross section with E1, E2 and M1 components to that for a pure E1 process. (a) Curve A, $(E1 + E2)/E1$ assumes one E2 sum; (b) curve B, $(E1 + M1)/E1$ assumes $B(M1) = 30\mu_0^2$; (c) curve C, $(E1 + M1 + E2)/E1$ assumes that E2 exhausts 100% of the EWSR and a $B(M1)$ of $30\mu_0^2$.

2. The experiment

The absolute (e, n) cross section of ^{197}Au was obtained by measuring the activity induced in the target. The isotope ^{196}Au decays to ^{196}Pt , with $T_{1/2} = (6.183 \pm 0.010)$ [11], emitting a 356 keV γ ray of absolute intensity $I_\gamma = (87.6 \pm 0.1)\%$ [11]. The γ rays were counted using a coaxial Ge(Li) detector system. The absolute efficiency was determined by counting the 356 keV γ rays ($I_\gamma = (62 \pm 1)\%$ [11]) emitted by a calibrated ^{133}Ba source mounted in the same geometry used for counting the targets. The targets, high purity ^{197}Au , 1.35 mg cm^{-2} thick (produced by A D Mackay, USA), were activated with the electron beam of the linear electron accelerator of the University of São Paulo. During the irradiations the incident charge was monitored with a Faraday cup. The (e, n) cross section is shown in figure 2. The error bars refer only to statistical uncertainty.

The yield of the reaction $^{197}\text{Au}(e + \gamma, n)^{196}\text{Au}$ was measured at 23 MeV using bremsstrahlung photons, which were produced by interposing in the beam a natural copper radiator (0.717 g cm^{-2} thick) just ahead of the target. During these irradiations (six in total) the linac current was monitored by a secondary emission monitor (SEM), calibrated relative to the Faraday cup, prior and after each irradiation. To obtain the bremsstrahlung yield the electrodisintegration contribution was subtracted, after properly correcting for energy loss in the radiator (Δr):

$$Y_{\text{br, n}}(E_0 - \Delta r/2) = Y_{\text{e+br, n}}(E_0) - \sigma_{\text{e, n}}(E_0 - \Delta r). \quad (4)$$

The results obtained for the electrodisintegration plus bremsstrahlung yield, $Y_{\text{e+br, n}}$, are shown in table 1.

3. Analysis and discussion

As already pointed out, M1 and isoscalar E2 are the only transitions, apart from E1, that make a detectable contribution to the (e, n) process in the energy region under

Figure 2. Measured $\sigma_{e,n}$ cross section of ^{197}Au . The curve is the result of fitting the data with E1, E2 and M1 multipoles (see text).

study. The measured (e, n) cross section was fitted by a linear regression using the least-squares method. Rewriting equation (2) we have:

$$\begin{aligned}
 \sigma_{e,n}^{\text{exp}}(E_0) = & k_1 \int_0^{E_0-m} \sigma_{\gamma,n}(E) N^{\text{E1}}(E_0, E, Z) \frac{dE}{E} \\
 & + k_2 \int_0^{E_0-m} \sigma_{\gamma,n}^{\text{E2}}(E) [N^{\text{E2}}(E_0, E, Z) - N^{\text{E1}}(E_0, E, Z)] \frac{dE}{E} \\
 & + k_3 \int_0^{E_0-m} \sigma_{\gamma,n}^{\text{M1}}(E) [N^{\text{M1}}(E_0, E, Z) - N^{\text{E1}}(E_0, E, Z)] \frac{dE}{E}
 \end{aligned} \quad (5)$$

Table 1. Measured (electrodisintegration + bremsstrahlung) yield for electrons of energy 23 MeV incident on a 0.717 g cm^{-2} natural copper radiator.

Measurement	$Y_{e,n} \text{ (mb)}$
1	9.16 ± 0.23
2	9.22 ± 0.48
3	8.99 ± 0.16
4	9.25 ± 0.43
5	9.27 ± 0.15
6	9.79 ± 0.12
Weighted mean	(9.40 ± 0.07)

with $\sigma_{\gamma,n}$, $\sigma_{\gamma,n}^{\text{E2}}$ and $\sigma_{\gamma,n}^{\text{M1}}$ as already defined. The constant k_1 was determined by comparing our experimental result for $Y_{\text{br},n}(E_0 - \Delta r/2)$, given by (4), with its calculated value:

$$k_1 = Y_{\text{br},n}^{\text{exp}}(E_0 - \Delta r/2) \left(N_r \int_0^{E_0 - \Delta r/2} \sigma_{\gamma,n}(E) [K(E_0 - \Delta r/2, E, Z)] \frac{dE}{E} \right)^{-1}. \quad (6)$$

Here N_r stands for the number of nuclei/cm² in our radiator and $K(E_0 - \Delta r/2, E, Z)$ is the Davies–Bethe–Maximon [12, 13] bremsstrahlung cross section.

The constant k_1 in equation (5) reflects the difference between our absolute scale and that from [8]. On the other hand, k_2 and k_3 , the linear regression coefficients that best fit our experimental (e, n) cross section, are directly related to the depletion of the isoscalar E2 EWSR and to $B(\text{M1}) \uparrow$, respectively. We obtain from equation (6), $k_1 = (0.99 \pm 0.03)$. Thus our absolute cross section is in very good agreement with the absolute photonuclear cross section for ^{197}Au [8] measured at Saclay.

The fit of the experimental $\sigma_{\gamma,n}(E)$ cross section yields an isoscalar E2 component exhausting $(70 \pm 64)\%$ of the corresponding EWSR and an M1 intensity (from 8.1 to 9.4 MeV) of $(44 \pm 6) \mu_0^2$. The result of the fit is shown by the full curve in figure 2. As could be expected from the results shown in figure 1, we obtain an undetermined value for the E2 component because its contribution to the (e, n) cross section up to ~ 25 MeV is small. In this energy region the effect of the E2 component is masked by the dominant E1 and M1 contributions. However, the experimental systematics for heavy nuclei [14, 15], and in particular results for ^{197}Au [15], show that the isoscalar E2 giant resonance exhausts $\sim 100\%$ of the corresponding EWSR in heavy nuclei. Thus we fitted our data again, setting $k_2 = 1$, corresponding to one E2 sum depletion, and obtained for the M1 strength: $B(\text{M1}) \uparrow = (41 \pm 3) \mu_0^2$.

Furthermore, we repeated the analysis using the Livermore Laboratory experimental values for $\sigma_{\gamma,n}(E)$ [16]. In this case, we get from (6), $k_1 = (1.13 \pm 0.03)$ and, assuming one E2 EWSR depletion, we get for the M1 strength: $B(\text{M1}) \uparrow = (45 \pm 3) \mu_0^2$. As shown in table 2, the results obtained for $B(\text{M1}) \uparrow$, assuming one isoscalar E2 sum and using either the Livermore or the Saclay data, are compatible with the $(44 \pm 6) \mu_0^2$ obtained using the Saclay (γ, n) cross section and leaving the E2 intensity as a free parameter.

Throughout this work we have used, for electric multipoles, DWBA spectra for finite nuclei [5]. It is well known that size corrections are needed for heavy nuclei and that the DWBA E1 and E2 spectra for finite nuclei are in agreement with experimental results [7].

It has already been shown that size effects are small for the transverse components of virtual photon spectra [17]. The E1 and E2 DWBA spectra for finite nuclei confirm this behaviour. The E1 spectrum is dominantly transverse and the magnitude of size

Table 2. M1 intensity.

$\sigma_{\gamma,n}(E)$	K_1	E2 EWSR (%)	$B(\text{M1}) \uparrow (\mu_0^2)$
Saclay [8]	(0.99 ± 0.03)	(70 ± 64)	(44 ± 6)
		100 fixed	(41 ± 3)
Livermore [16]	(1.13 ± 0.03)	100 fixed	(45 ± 3)

corrections is relatively small as compared with that for the E2 spectrum, which is dominantly longitudinal (see [7] and [17]). Since M1 is purely transverse, size corrections should have little effect on the corresponding virtual photon spectra. However, the DWBA M1 spectra for finite nuclei [5] differ considerably from the results obtained for a point nucleus (factors of two or more in certain parts of the spectrum for a heavy nucleus). Such unexpected behaviour led us to perform other tests. We found that isochromats of the M1 spectrum for a finite nucleus yield unacceptable results. As an example, for $Z=73$ the number of 7.0 MeV virtual photons decreases as the electron energy increases for E_0 above 18 MeV. It is known that an isochromat of real or virtual photons increases monotonically with the incident electron energy, because the number of radiated photons of a given energy never decreases as the electron energy increases.

This behaviour of the M1 spectra for finite nuclei suggests that the model used to describe nuclear current densities, which is derived from the assumption of irrotational incompressible flow in the nucleus, is inadequate. Even though this affects strongly the magnetic multipoles, it has little effect on the electric multipoles. For the latter it is found [5] that the results are rather independent of the detailed shape of the charge distribution. They depend primarily on the values of the nuclear root-mean-square radius and the appropriate transition radius.

For M1 we have used both point and finite DWBA spectra, but consistent results were obtained only with the first type. As expected, the M1 strength derived using finite DWBA spectra was much larger than that obtained with point DWBA. So large, in fact, that the M1 component of the photonuclear cross section obtained is bigger than the cross section measured with monochromatic photons (between 8.1 and 9.4 MeV), which should include all components. This result is physically unacceptable because it implies the existence of a negative E1 strength in this energy region.

Pondering over the problems that seem to affect the finite DWBA spectra we consider the point DWBA to be the most reliable calculation currently available.

4. Conclusions

Electrodisintegration experiments can be used to obtain information on M1 strength located above the (γ, n) threshold in heavy nuclei.

The obtained value, $B(M1) \uparrow = (44 \pm 6) \mu_n^2$, is compatible with previous experimental results [1, 9] and with theoretical estimates [9, 10], according to which $20 < B(M1) \uparrow < 50 \mu_n^2$ in ^{197}Au .

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Isovector Giant Quadrupole Resonance in ^{63}Cu

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Abstract The decay of the isovector $E2$ giant resonance in ^{63}Cu has been studied by measuring the $(e, 2n)$ cross section, in the incident electron energy range 22-45 MeV. The photodisintegration induced by bremsstrahlung was also measured. The electrodisintegration results have been analyzed using the distorted wave Born approximation $E1$ and $E2$ virtual photon spectra to obtain these multipole components in the corresponding $(\gamma, 2n)$ cross section. It is found that the isovector $E2$ giant resonance decays dominantly by two-neutron emission in ^{63}Cu . This decay channel exhausts 65 percent of the energy weighted $E2$ sum.

1. INTRODUCTION

The possible existence of a giant quadrupole resonance was discussed soon after the explanation of the nuclear photoeffect¹. The latter was explained as a dipole oscillation of the nucleus². Using the excitation energy $E_x(E1)$ of the giant dipole resonance (GDR) as normalization, the excitation energy of the corresponding quadrupole mode, $E_x(E2)$, was calculated based on hydrodynamical analogies and found to be 1.6 times $E_x(E1)$ ³. Microscopic models for the $E1$ were developed later⁴, but since the dipole resonance has no experimentally identifiable isoscalar part, there was no immediate need to consider separately isoscalar and isovector interactions. The attempt to unify the macroscopic and microscopic treatment of the giant dipole resonance⁵ led to the incorporation of isospin-dependent particle-hole interactions into the microscopic model⁴. The unified description leads to the prediction of a splitting of all giant modes into isovector and isoscalar parts. The isoscalar parts are always lower in energy than the unperturbed state because of the attractive nature of the isoscalar particle-hole force, while the isovector parts are higher in energy. For the case of quadrupole excitations the unperturbed energy of $2\hbar\omega_0 = 80A^{-1/3}$ MeV changes into $-60A^{-1/3}$ MeV for $\Delta T = 0$ and $-130A^{-1/3}$ MeV for $\Delta T = 1$ ⁶.

For the isoscalar giant quadrupole resonance there is already a good systematics in the literature^{7,8}. Most of the data about this collective mode come from hadron inelastic scattering, which selectively excites isoscalar modes. For the isovector quadrupole mode there is comparatively little experimental evidence. Since isovector excitations are suppressed in hadronic interactions, the evidence comes from electromagnetic experiments: capture reactions and electron scattering. Capture reaction measurements can give direct indication of the isoscalar and isovector nature of excitations by comparing magnitude and sign of the appropriate Legendre coefficients in, e.g., (γ, p_0) and (γ, n_0) angular distributions. However they observe only that part of the isovector quadrupole mode that decays to the ground state of the residual nucleus, which is usually only a few percent of the total strength. Inclusive electron scattering excites all modes and it is very difficult to separate them. As an example, fig. 1 shows the decomposition of the electron scattering cross section in ^{63}Cu into isovector $E1$ and $E2$ and isoscalar $E2$ resonances². From this decomposition the isovector quadrupole is located around 32 MeV excitation energy and exhausts 82% of the corresponding sum rule.

It has to be pointed out that inclusive electron scattering has to take recourse to general theoretical expectations and to comparison with hadron inelastic scattering, which suppresses isovector excitations, in order to establish evidence for the isonature of a certain excitation. Furthermore, since the isovector $E2$ resonance appears in the tail region of the $E1$, the strength is very dependent on the radiative tail subtraction, as can be seen in fig. 1. Apart from that, there is an increased model dependence of these results as compared to those for the isoscalar resonance. While for the latter it is generally accepted that it can be represented accurately enough by a surface oscillation (Tassie model, Goldhaber-Teller model), this may no longer be true for isovector oscillations, as has been shown for the dipole case¹⁰. Another point is that isovector sum rules are not as model independent as isoscalar sums, because of the effects of exchange contributions. The isovector energy weighted sum rule, which is obtained from the isoscalar one replacing Z^2 by NZ , has to be regarded as a lower limit for the isovector strength.

Fig.1 - Isovector E1 and E2 and isoscalar E2 giant resonances observed in electron scattering in ^{63}Cu .

Fig. 2 shows the location of the isovector E2 resonance in several nuclei. The data are from electron scattering and capture reactions and were obtained from the compilations of Bertrand⁷ and Pittman¹¹ plus the recent data of ref. 9. As fig. 2 shows, the isovector E2 is located at $120\text{--}130A^{-1/3}$ MeV for most nuclei. According to ref. 11 the isovector E2 mode exhausts $\sim 100\%$ of the energy weighted sum in most nuclei.

Since we have poor information on the isovector E2 mode as yet, a better knowledge of this mode is of great interest. The reason is that we can learn about features that are not present in the isoscalar mode. The main properties which influence the isovector strengths, which play little or no role for the isoscalar case, are velocity-dependent and exchange effects in nuclear interactions⁶. Once we have learned to measure the isovector resonances accurately and to evaluate results independent of model, we will have more stringent tests of nuclear models.

Fig.2 - Systematics for the excitation energy of the isovector E2 giant resonance versus the atomic mass A.

Apart from the lack of accurate data on the isovector quadrupole mode, there is no data about its decay properties. There is recent evidence that the decay of giant resonances is dominantly statistical¹². If this is also the behaviour of the isovector quadrupole mode, we would expect that in ^{63}Cu it would decay dominantly by emission of two neutrons, since this is the dominant decay mode of photoabsorption around 22 MeV excitation energy. The (γ, n) decay channel becomes vanishingly small above 27 MeV¹³. This has motivated us to search for the isovector E2 in the $(\gamma, 2n)$ decay channel.

In this paper we report measurements of the $(\gamma, 2n)$ electrodisintegration cross section in ^{63}Cu . In order to perform the multipole decomposition, the virtual photon method is used.

2. THE VIRTUAL PHOTON METHOD

During the last decade the study of the giant multipole resonances, together with the emergence of virtual photon calculations using the distorted-wave Born approximation (DWBA) of Onley and collaborators¹⁴⁻¹⁷, have jointly stimulated a renaissance of interest in electrodisintegration experiments.

The basic premise of the virtual photon technique is that the electrodisintegration cross section, $\sigma_{e,x}(E_0)$, for incident electrons of energy E_0 , can be expressed in terms of the photonuclear cross section, $\sigma_{\gamma,x}^{\lambda L}(E)$ associated with the absorption of real photons of energy E and multipolarity λL through the virtual photon spectrum, $N^{\lambda L}(E_0, E, Z)$

$$\sigma_{e,x}(E_0) = \int_0^{E_0 - m_e} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z) \frac{dE}{E} \quad (1)$$

In fig. 3 the E1 and E2 DWBA virtual photon spectra for electrons of incident energy 45 MeV are shown for Cu. The advantage of an electrodisintegration experiment is that it enhances E2 absorption, because the E2 spectrum is bigger than E1. Even though virtual photon spectra for finite nuclei are already available¹⁷, throughout this work we will use DWBA spectra for a point nucleus, because for the energy range considered here in ⁶³Cu size corrections are negligible.

Fig.3 - E1 and E2 virtual photon spectra for 45 MeV electrons inelastically scattered from a copper nucleus.

The enhancement of the $E2$ relative to the $E1$ virtual photon spectra has already been exploited in several experiments to study $E2$ absorption by nuclei, e.g. references 18-21. In these experiments the photodisintegration cross section induced by bremsstrahlung, $\sigma_{br,x}(E_0)$, is also measured by interposing a radiator ahead of the target:

$$\sigma_{br,x}(E_0) = N_r \int_0^{E_0-m} \sum_{\lambda L} \sigma_{Y,x}^{\lambda L}(E) K(E_0, E, Z) \frac{dE}{E} \quad (2)$$

where N_r is the number of nuclei/cm² in the radiator and $K(E_0, E, Z)$ is the bremsstrahlung cross section. By using the radiator the multipole composition of the photons seen by the target is changed. For ⁶³Cu and the energies involved here the virtual spectrum contains about twice as many $E2$ photons as $E1$ photons, but the bremsstrahlung spectrum has no preferred multipoles.

Several experimental checks of the DWBA virtual photon spectra have been made²¹⁻²⁵, but the most accurate of these tests was an experiment where an isochromat of the $E1$ virtual photon spectrum was measured by counting the number of ground-state protons emitted by the 16.28 MeV isobaric analogue state in ⁹⁰Zr as a function of incident electron energy²⁵. The experimental results agree with those predicted by the virtual photon calculation within the experimental error of one percent.

The bremsstrahlung cross section which enters in eq.(2) has also been subjected to experimental tests²⁶. It was shown that the Davies-Bethe-Maximon bremsstrahlung cross section, which will be used in this work, is in excellent agreement with photodisintegration experimental results above 28 MeV.

3. THE EXPERIMENT

The ⁶³Cu($e, 2n$) cross section was measured by residual activity. The residual nucleus ⁶³Cu decays by β^+ (62%) and electron capture (38%) to ⁶¹Ni²⁷. We have followed the decay by detecting the 67.4 keV γ -ray line that results from the decay to ⁶¹Ni, using a Ge-Li detector system.

The targets (five) used are 99.89 enriched ^{63}Cu with a thickness of $10.16 \pm 0.10 \text{ mg/cm}^2$ measured by weighing. The targets were activated in the electron beam of the São Paulo Linear Electron Accelerator. The LINAC current was measured by a Faraday Cup.

We have also measured the photodisintegration induced by bremsstrahlung using a 0.358 g/cm^2 copper radiator, placed in the electron beam just ahead of the target. The LINAC current was measured using a secondary emission monitor calibrated relative to a Faraday Cup. The measured yields were corrected for energy loss of electrons in the radiator and the electrodisintegration produced in the targets was subtracted.

The results obtained for the electrodisintegration cross section and the bremsstrahlung induced photodisintegration are shown in fig. 4 by the open and full circles respectively.

Fig.4 - $\sigma_{e,2n}(E_0)$ for ^{63}Cu (open circles) and the photodisintegration induced by bremsstrahlung (full circles). The statistical errors are of the size of the points or smaller. The smooth curves (E1+E2) and σ_{br} are the best fit to the data (see text).

4. DISCUSSION

In order to decompose our experimental electrodisintegration cross section in terms of E1 and E2 contributions we write eq.(1) as:

$$\begin{aligned} \sigma_{e,2n}(E_0) = & \int_0^{E_0-m_e} \sigma_{\gamma,2n}(E) N^{E1}(E_0, E, Z) \frac{dE}{E} \\ & + \int_0^{E_0-m_e} \sigma_{\gamma,2n}^{E2}(E) \left[N^{E2}(E_0, E, Z) - N^{E1}(E_0, E, Z) \right] \frac{dE}{E} \end{aligned} \quad (3)$$

Since the $(\gamma, 2n)$ threshold is at 19.7 MeV, in the energy range covered by our experiment the multipoles that are expected to contribute are the isovector E1 and isovector E2. Around 30 MeV there should also be the isoscalar E3 mode, but it gives negligible contribution to the electrodisintegration cross section (see ref.21). In eq.(3) we use for $\sigma_{\gamma,2}(E)$ available photonuclear data²⁸. Since the photonuclear data do not distinguish different multipoles we are assuming:

$$\begin{aligned} \sigma_{\gamma,2n} &= \sigma_{\gamma,2n}^{E1} + \sigma_{\gamma,2n}^{E2} \\ \text{or} \\ \sigma_{\gamma,2n}^{E1} &= \sigma_{\gamma,2n} - \sigma_{\gamma,2n}^{E2} \end{aligned} \quad (4)$$

For $\sigma_{\gamma,2n}^{E2}(E)$ we will use a Lorentzian shape with peak at $120A^{-1/3}$ ($E_p = 30$ MeV) and a width of 8 MeV. The systematics for the isovector E2 indicates that it is a broad resonance as illustrated in fig.1.

The curves labelled (E1+E2) and σ_{br} shown in fig.4 are the best fits to our data using eq.(3) for the electrodisintegration cross section and eq.(2) for the bremsstrahlung yield simultaneously. Fig.5 shows the photonuclear data from ref. 28 decomposed into E1 and E2, as obtained from the fit to our data. Just to illustrate the effect, curve labelled E1 in fig. 4 is the result of integrating the cross section of fig. 5 assuming it is all E1. The difference between curves (E1+E2) and E1 is 40% at 45 MeV. The calculated photodisintegration yield, of course, remains unchanged if we assume that the photonuclear cross section is all E1. The difference between curves (E1+E2) and E1 shows the sensitivity of electrodisintegration cross section to the multipole composition of the photonuclear process. The isovector E2 component obtained from the fit exhausts 65 percent of the energy weighted sum. This result

Fig. 5 - The data points show the photo-nuclear cross section $(\gamma,2n)$ from ref.28. The curves show the decomposition into E1 and E2 components.

is consistent with the electron scattering result of ref. 9, which finds that the isovector $E2$ exhausts 82 percent of the corresponding sum. Since we observed only the dominant decay channel, the remaining strength can be in other channels $(\gamma,pn;\gamma,2p;\dots)$.

5. CONCLUSIONS

We have measured the $(e,2n)$ cross section for ^{63}Cu . Combining this measurement with photodisintegration induced by bremsstrahlung we have shown that the $(e,2n)$ cross section is the dominant decay channel for the isovector $E2$ resonance in this nucleus, exhausting 65 percent of the energy weighted sum.

The electrodisintegration cross section can be used as an important tool to measure the decay of $E2$ strength in nuclei.

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Resumo

O decaimento da ressonância gigante isovetorial de quadrupolo elétrico no ^{63}Cu foi estudado através da medida da secção de choque $(\gamma, 2n)$, usando-se elétrons de energia 22-45 MeV. A fotodesintegração induzida por bremsstrahlung foi também medida. Os resultados da eletrodesintegração foram analisados usando-se os espectros de fótons virtuais $E1$ e $E2$ calculados na aproximação de Born de ondas distorcidas, a fim de se obter a composição multipolar de processo $(\gamma, 2n)$ correspondente. Obtemos que a ressonância gigante isovetorial $E2$ decai predominantemente pela emissão de dois nêutrons. Esse canal de decaimento esgota 65 por cento da correspondente regra da soma ponderada em energia.

Comment on photoneutron cross sections

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The differences between the Saclay and Livermore photoneutron cross sections are discussed. It is shown that the differences between Saclay and Livermore (γ, n) and $(\gamma, 2n)$ cross sections arise from the neutron multiplicity sorting.

[NUCLEAR REACTIONS Photoneutron cross sections analyzed.]

In the last 20 years, photoneutron cross sections have been measured for many nuclei using monoenergetic photons. Most of this work was carried out at two laboratories: Saclay and Livermore. The available results are compiled in the *Atlas of Photoneutron Cross Sections Obtained with Monoenergetic Photons*.¹ There are also in the literature a few review articles on the subject, but none of these publications has addressed the problem of the differences between the measurements performed at Saclay and Liver-

more. In this paper we compare the Saclay and Livermore measurements for the nuclei listed in Table I.

The typical differences between Saclay and Livermore data are illustrated in Fig. 1, where the (γ, n) measurements from Saclay and Livermore are shown. The results from Livermore are multiplied by 1.06 in order to show both cross sections in the same absolute scale. The cross sections are in good agreement up to the $(\gamma, 2n)$ threshold. Above this energy there is an important difference: The

TABLE I. Nuclei measured by Saclay (S) and Livermore (L).

Nucleus	$\int \sigma_{\gamma, n}(E_{\gamma}) dE_{\gamma}^a$ (MeV mb)	$\int \sigma_{\gamma, 2n}(E_{\gamma}) dE_{\gamma}^a$ (MeV mb)	Reference	r
⁸⁹ Y	1279 S	74 S	2	1.255 ± 0.005
	960 L	99 L	3	
¹¹⁵ I	1470 S	278 S	4	0.942 ± 0.004
	1354 L	508 L	5	
¹¹⁷ Sn	1334 S	220 S	4	1.012 ± 0.007
	1380 L	476 L	5	
¹¹⁸ Sn	1377 S	258 S	4	
	1302 L	531 L	5	
¹²⁰ Sn	1371 S	399 S	4	0.987 ± 0.004
	1389 L	673 L	5	
¹²⁴ Sn	1056 S	502 S	4	0.929 ± 0.006
	1285 L	670 L	5	
¹³³ Cs	1828 S	328 S	4	1.106 ± 0.007
	1475 L	503 L	6	
¹⁵⁹ Tb	1936 S	605 S	7	1.062 ± 0.001
	1413 L	887 L	8	
¹⁶⁵ Ho	2090 S	7665	7	1.136 ± 0.007
	1735 L	744 L	9	
¹⁸¹ Ta	2180 S	790 S	7	1.218 ± 0.018
	1300 L	881 L	10	
¹⁹⁷ Au	2588 S	479 S	11	1.004 ± 0.013
	2190 L	777 L	12	
²⁰⁸ Pb	2731 S	328 S	11	1.296 ± 0.011
	1776 L	860 L	13	

^aFrom Ref. 1.

FIG. 1. (γ, n) cross sections from Saclay (solid line) and Livermore (experimental points) for ^{159}Tb . The Livermore data are multiplied by 1.06 in order to show both measurements at the same absolute scale.

FIG. 2. $(\gamma, 2n)$ cross sections from Saclay (solid line) and Livermore (experimental points) for ^{159}Tb . The Livermore data are multiplied by 1.06 in order to show both measurements at the same absolute scale.

Livermore cross section vanishes a few MeV above the $(\gamma, 2n)$ threshold, in good agreement with the predictions of the statistical model, while the Saclay cross section has a tail. In Ref. 7, the observed tail of the Saclay cross section is interpreted as arising from fast neutrons that would have escaped detection in the Livermore measurements, leading to the conclusion that for ^{159}Tb the contribution of the "direct effect" in the photoneutron cross section is $n_d = 23 \pm 4\%$. In Table II the percentages of direct neutrons inferred by Saclay are given for several nuclei.

Figure 2 shows the $(\gamma, 2n)$ cross sections from Saclay and Livermore. The $(\gamma, 2n)$ cross sections differ in shape and magnitude, the Livermore one being much bigger. Even though up to the $(\gamma, 2n)$ threshold the (γ, n) cross section from Livermore, $\sigma_{\gamma, n}^L$, and that from Saclay, $\sigma_{\gamma, n}^S$, differ by

only 6% in the absolute scale, their integrated cross sections up to 28 MeV are 1413 and 1936 MeVmb, respectively. While the integrated (γ, n) cross section from Saclay is 37% bigger than the Livermore result, their integrated $(\gamma, 2n)$ cross section is 47% smaller.

In order to understand these differences we reconstructed the total neutron measurements from Saclay and Livermore:

$$\sigma_{\gamma, \text{tn}} = \sigma_{\gamma, n} + 2\sigma_{\gamma, 2n} + 3\sigma_{\gamma, 3n} \quad (1)$$

Since $\sigma_{\gamma, \text{tn}}$ is the cross section measured and the (γ, n) and $(\gamma, 2n)$ cross sections are obtained by neutron multiplicity sorting, it is important to compare $\sigma_{\gamma, \text{tn}}^S$ and $\sigma_{\gamma, \text{tn}}^L$. This comparison has not been discussed in the literature. Figure 3 shows $\sigma_{\gamma, \text{tn}}$ from Saclay and Livermore for ^{159}Tb . In Fig. 4 the ratio $r = \sigma_{\gamma, \text{tn}}^S / \sigma_{\gamma, \text{tn}}^L$ is shown. The ratio is reasonably constant and the least-squares fit of a constant yields the

TABLE II. Percentage of direct neutrons obtained by Saclay.

Nucleus	n_d (%)	References
^{94}Mo	25 ± 7	14
^{96}Mo	15 ± 7	14
^{98}Mo	10 ± 7	14
^{100}Mo	11 ± 7	14
^{139}La	28 ± 5	7
^{140}Ce	12 ± 3	15
^{142}Ce	10 ± 3	15
^{nat}Sm	10 ± 3	15
^{159}Tb	23 ± 4	7
^{165}Ho	23 ± 4	7
^{nat}Er	11 ± 3	15
^{175}Lu	15 ± 3	15
^{181}Ta	22 ± 2	7
^{197}Au	20 ± 4	11
^{208}Pb	15 ± 4	11
^{238}U	14 ± 2	16

FIG. 3. $\sigma_{\gamma, \text{tn}}$ from Saclay divided by $\sigma_{\gamma, \text{tn}}$ from Livermore. $\sigma_{\gamma, \text{tn}} = \sigma_{\gamma, n} + 2\sigma_{\gamma, 2n} + 3\sigma_{\gamma, 3n}$.

FIG. 4. $\sigma_{\gamma, \text{tn}}$ from Livermore multiplied by 1.06 (experimental points) and $\sigma_{\gamma, \text{tn}}$ from Saclay (solid line). $\sigma_{\gamma, \text{tn}} = \sigma_{\gamma, \text{n}} + 2\sigma_{\gamma, 2\text{n}} + 3\sigma_{\gamma, 3\text{n}}$.

value $r = 1.062 \pm 0.011$. In order to compute r we interpolated $\sigma_{\gamma, \text{tn}}^{\text{S}}$ and $\sigma_{\gamma, \text{tn}}^{\text{L}}$, since their data points are not at the same photon energies. One important conclusion can be derived from Fig. 3: Both laboratories are detecting the same number of neutrons for all photon energies. If there were fast neutrons escaping detection in the Livermore measurements above 20 MeV, r should increase above this energy. The value of the constant r is the difference in the absolute scale of both measurements. Figure 4 shows $\sigma_{\gamma, \text{tn}}^{\text{L}}$ multiplied by 1.06 and $\sigma_{\gamma, \text{tn}}^{\text{S}}$, just to illustrate the good agreement between them, when they are plotted at the same absolute scale.

Since both laboratories agree as to the total number of neutrons detected, the differences in their (γ, n) and $(\gamma, 2\text{n})$ cross sections arise from the separation of the total counts into (γ, n) and $(\gamma, 2\text{n})$ events (neutron multiplicity sorting).

If we assume that the excess (γ, n) cross section in the Saclay measurement is due to $(\gamma, 2\text{n})$ events interpreted as two (γ, n) events, that is, if we compute

$$\sigma_{\gamma, 2\text{n}}^{\text{S}} = \sigma_{\gamma, 2\text{n}}^{\text{L}} + \frac{1}{2}(\sigma_{\gamma, \text{n}}^{\text{S}} - 1.06\sigma_{\gamma, \text{n}}^{\text{L}}), \quad (2)$$

FIG. 5. $\sigma_{\gamma, 2\text{n}}$ from Livermore (experimental points) and the modified (see text) $\sigma_{\gamma, 2\text{n}}$ from Saclay (solid line).

we obtain for $\sigma_{\gamma, 2\text{n}}^{\text{S}}$ the solid line shown in Fig. 5. The modified $\sigma_{\gamma, 2\text{n}}$ cross section from Saclay agrees well with the $(\gamma, 2\text{n})$ cross section from Livermore multiplied by 1.06 (data points).

The same analysis carried out for ^{159}Tb was repeated for all nuclei listed in Table I. The results obtained for these nuclei are presented in detail in a report from our Institute, available upon request.¹⁷ Here we show only the results obtained by fitting a constant r to the ratio $\sigma_{\gamma, \text{tn}}^{\text{S}}/\sigma_{\gamma, \text{tn}}^{\text{L}}$. The values of r and the standard deviation of the mean are listed in Table I. For all these nuclei we obtained results similar to ^{159}Tb :

(a) Both laboratories are detecting the same number of neutrons versus the incident photon energy.

(b) $\sigma_{\gamma, \text{n}}^{\text{S}}$ is bigger than $r\sigma_{\gamma, \text{n}}^{\text{L}}$ above the $(\gamma, 2\text{n})$ threshold.

(c) $\sigma_{\gamma, 2\text{n}}^{\text{S}}$ is smaller than $r\sigma_{\gamma, 2\text{n}}^{\text{L}}$.

(d) If $\sigma_{\gamma, 2\text{n}}^{\text{S}}$ is modified using Eq. (2), then $\sigma_{\gamma, 2\text{n}}^{\text{S}}$ and $\sigma_{\gamma, \text{n}}^{\text{S}}$ are in agreement with $r\sigma_{\gamma, 2\text{n}}^{\text{L}}$, respectively.

The analysis performed here shows that the differences in shape and magnitude in the (γ, n) and $(\gamma, 2\text{n})$ cross sections measured by Saclay and Livermore arise from the neutron multiplicity sorting.

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Electrodisintegration of ^{208}Pb , ^{209}Bi and ^{181}Ta

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Abstract. The electrodisintegration cross section has been measured for ^{208}Pb , ^{209}Bi and ^{181}Ta , from threshold to 22 MeV, by detecting the emitted neutrons. The quantity measured is $\sigma_{\text{e},1\text{n}} = \sigma_{\text{e},\text{n}} + 2\sigma_{\text{e},2\text{n}}$. The results are analysed with DWBA virtual photon spectra for finite nuclei to obtain M1 and E2 strengths.

NUCLEAR REACTIONS ^{208}Pb , ^{209}Bi and ^{181}Ta ; measured $\sigma(\text{e}, \text{n}) + 2\sigma(\text{e}, 2\text{n})$; $E = 8\text{--}22$ MeV; virtual photon analysis; deduced E2 and M1 strengths.

1. Introduction

It has been shown recently [1] that the new calculations of virtual photon spectra for finite nuclei [2, 3] are in good agreement with experimental results, that is they correctly relate the electrodisintegration cross section, $\sigma_{\text{e},x}(E_0)$, and the photonuclear cross section, $\sigma_{\gamma,x}(E)$. This relationship is given by

$$\sigma_{\text{e},x}(E_0) = \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z) \frac{dE}{E}. \quad (1)$$

Here $N^{\lambda L}(E_0, E, Z)/E$ stands for the number of virtual photons of multipolarity λL exchanged with a nucleus of atomic number Z when an electron of energy E_0 produces an excitation of energy E . As discussed in [1], even though the sum in equation (1) goes up to infinity, for electron incident energies below about 30 MeV, the multipoles that make detectable contributions to the electrodisintegration cross section are E1, M1 and E2, the cross section being particularly sensitive to the M1 strength, which is located in the vicinity of the particle emission threshold (see figure 6 of [1]).

We have measured the electrodisintegration of ^{208}Pb , ^{209}Bi and ^{181}Ta by neutron emission, aiming to obtain the intensities of isoscalar and isovector E2 absorption, as well as the M1 intensity located above the (γ, n) threshold.

2. The experiment

The $\sigma_{\text{e},1\text{n}}$ cross sections ($\sigma_{\text{e},1\text{n}} = \sigma_{\text{e},\text{n}} + 2\sigma_{\text{e},2\text{n}}$) were measured by counting the neutrons. The neutron detector system consists of four $^{10}\text{BF}_3$ counters embedded in paraffin. The efficiency of the system was determined by measuring the (e, n) cross section in

Table 1. Target thickness, composition and separation energies.

Nucleus	Thickness (mg cm ⁻²)	Composition	Separation energies	
			γ, n (MeV)	$\gamma, 2n$ (MeV)
¹⁸¹ Ta	19.3 ± 0.1 ^a	Natural	7.6	14.2
²⁰⁸ Pb	3.8 ± 0.8 ^b	99% enriched	7.4	14.1
²⁰⁹ Bi	2.86 ± 0.01 ^b	Natural	7.5	14.4

^a Self-supporting.^b Deposited over Al backing.

⁶³Cu and comparing it with the absolute measurements of this cross section performed at the National Bureau of Standards [4]. The detector efficiency is constant for the range of electron energies used here. The efficiency was also measured with a ²⁵²Cf source, which was often used to check the system during the experiment. The targets were bombarded in the electron beam of the São Paulo Electron Linear Accelerator. Table 1 gives the composition, thickness and separation energies for the targets used. The accelerator current was monitored with an ionisation chamber that was calibrated relative to a Faraday cup prior to and after each irradiation. The stability was better than 1%.

3. Results and analysis

The measured cross sections are shown in figures 1–3. In order to analyse the data using equation (1) we have to use experimental values for the photonuclear cross section $\sigma_{\gamma, Tn} = \sigma_{\gamma, n} + 2\sigma_{\gamma, 2n}$. These results are available from both the Saclay and

Figure 1. $\sigma_{e, Tn}$ for ¹⁸¹Ta. The curve is the result of the fit to the data.

Figure 2. $\sigma_{e, Tn}$ for ^{208}Pb . The curve is the result of the fit to the data.

Livermore laboratories, apart from the nucleus ^{209}Bi which has only been measured at Livermore. Since it has been shown that Saclay and Livermore are in very good agreement as to the shape of all their measured $\sigma_{e, Tn}$ cross sections [5], it is irrelevant which data are used in equation (1). We chose to use the Saclay data, when available,

Figure 3. $\sigma_{e, Tn}$ for ^{209}Bi . The curve is the result of the fit to the data.

because their measurements start at the threshold and have more experimental points with smaller error bars.

As in [1] we will assume $\sigma_{\gamma, \text{Tn}} = \sigma_{\gamma, \text{Tn}}^{\text{E1}} + \sigma_{\gamma, \text{Tn}}^{\text{E2}} + \sigma_{\gamma, \text{Tn}}^{\text{M1}}$ and write equation (1) as

$$\begin{aligned} \sigma_{\gamma, \text{Tn}}(E_0) = & k_1 \int_0^{E_0-m} \sigma_{\gamma, \text{Tn}}(E) N^{\text{E1}}(E_0, E, Z) \frac{dE}{E} \\ & + k_2 \int_0^{E_0-m} \sigma_{\gamma, \text{Tn}}^{\text{E2}(T=0)}(E) [N^{\text{E2}}(E_0, E, Z) - N^{\text{E1}}(E_0, E, Z)] \frac{dE}{E} \\ & + k_3 \int_0^{E_0-m} \sigma_{\gamma, \text{Tn}}^{\text{E2}(T=1)}(E) [N^{\text{E2}}(E_0, E, Z) - N^{\text{E1}}(E_0, E, Z)] \frac{dE}{E} \\ & + k_4 \int_0^{E_0-m} \sigma_{\gamma, \text{Tn}}^{\text{M1}}(E) [N^{\text{M1}}(E_0, E, Z) - N^{\text{E1}}(E_0, E, Z)] \frac{dE}{E}. \end{aligned} \quad (2)$$

In equation (2) $\sigma_{\gamma, \text{Tn}}(E)$ is the experimental photonuclear cross section for the corresponding nucleus [6–8], $\sigma_{\gamma, \text{Tn}}^{\text{E2}(T=0)}$ and $\sigma_{\gamma, \text{Tn}}^{\text{E2}(T=1)}$ will be represented by Lorentzian shapes, with parameters that are in agreement with the systematics for the isoscalar and isovector E2 resonances and exhausting one energy-weighted sum (EWSR). For $\sigma_{\gamma, \text{Tn}}^{\text{M1}}$ we assumed a strength of $1 \mu_0^2$ uniformly distributed in a region of 2 MeV just above the (γ, n) threshold, since the M1 will be distributed into many levels and our measurement is not sensitive to the detailed structure.

We fitted our data with equation (2) to obtain the values of k_1 to k_4 , where k_1 gives the difference between our absolute scale and that of the $\sigma_{\gamma, \text{Tn}}$ measurement; k_2 and k_3 give the fraction of the isoscalar and isovector E2 sums and k_4 the M1 intensity in units of μ_0^2 . The full curves of figures 1–3 show the results of the fits and table 2 gives the parameters used for the isoscalar E2 and the strengths obtained. The isovector E2 was located at $120 A^{-1/3}$ MeV with a width of 6 MeV for ^{181}Ta and ^{209}Bi ($E_p = 21.2$ and 20.2 MeV respectively) and at 22.5 MeV with a width of 5 MeV for ^{208}Pb based on previous experimental evidence [9, 10]. The intensities obtained for M1 and the

Table 2. Parameters for isoscalar E2.

Nucleus	E_p (MeV)	(MeV)	% EWSR	Reaction	Reference
^{181}Ta	11.1 ± 0.2	3.9	42 ± 8	(p, p')	[13]
	9.54 ± 0.02	2.07 ± 0.35	8 ± 4	(e, e')	[14]
	11.47 ± 0.22	3.13 ± 0.55	22 ± 8		
	9.5 ± 0.2	1.8 ± 0.6	29 ± 5	(e, e')	[15]
	11.3 ± 0.2	2.2 ± 0.7	63 ± 8		
	12.6 ± 0.2	1.3 ± 0.8	11 ± 5		
^{208}Pb	11.3	4	95 ± 50	(e, n)	This work
	10.9 ± 0.3	3.0 ± 0.3	145 ± 30	(a, a')	[11]
	11.0 ± 0.2	2.7 ± 0.3	105 ± 254	(a, a')	[12]
	10.8	3.2	76	(c, c')	[10]
	10.9	3.0	157 ± 62	(e, n)	This work
^{209}Bi	10.7 ± 0.3	2.7 ± 0.3	0–150	(a, a')	[11]
	10.9	2.7	50 ± 30	(e, n)	This work

isovector E2 resonance are given in table 3. For k_1 we have obtained the values 0.96 ± 0.02 for ^{208}Pb , 1.08 ± 0.02 for ^{209}Bi and 1.06 ± 0.02 for ^{181}Ta .

4. Discussion

The values obtained for k_1 are close to unity, showing that our absolute scale is in agreement with the absolute scale of the corresponding photonuclear cross sections [6–8].

The M1 strength obtained for ^{208}Pb is of the same order as that observed in ^{206}Pb and is in agreement with the theoretical prediction of $20 \mu_0^2$ [7]. For ^{181}Ta there are no theoretical predictions to compare with. Since we can only observe the M1 intensity above the (γ, n) threshold, the obtained values must be understood as lower limits. In ^{209}Bi all the strength could be below the neutron emission threshold as in ^{206}Pb [16].

The isoscalar E2 intensity is in agreement with the existing systematics which shows that in heavy nuclei it exhausts 80–100% of the EWSR [18]. However the uncertainties in our results for the isoscalar E2 are large and to improve their accuracy the experiment should be extended to higher electron bombarding energies, which is not feasible with our detector system because of the higher neutron multiplicities. Since the isoscalar E2 is in the (γ, n) channel, we should be able to separate the (e, n) channel and extend the experiment up to 50–60 MeV in order to enhance its contribution relative to E1. The isovector E2 is basically undetermined. However, if we exclude it from equation (2) the fits get worse, all points higher in energy than 18 MeV being above the calculated curves.

We carried out an interesting exercise with the isovector E2 resonance. When the photonuclear cross sections are fitted with Lorentzian shapes representing the E1 absorption, the points in the high-energy tail of the resonance are always above the calculated curve [19]. We fitted the photonuclear data for $\sigma_{\gamma, \text{xn}} = \sigma_{\gamma, n} + \sigma_{\gamma, 2n}$ with Lorentzian shapes in the following way: one representing the isoscalar E2 and another representing the isovector E2, both with the corresponding parameters of our best fit, and a third Lorentzian shape with free parameters to represent E1. In the case of deformed ^{181}Ta , there were, of course, two Lorentzian shapes with free parameters to represent E1. The results obtained are shown in figures 4–6. It is also interesting to compare the parameters for the E1 resonance obtained with this procedure with those from the fit of the photonuclear data assuming only E1 absorption [19]. This comparison is shown in table 4. The parameters for the E1 are nearly unchanged. This indicates that the photonuclear cross sections are compatible with the existence of an isovector E2 strength exhausting one sum and located at about $120 A^{-1/3}$ MeV. Under this assumption the photonuclear data can be well fitted even at the high-energy tail.

Table 3. M1 and isovector E2 intensities.

Target	Isvector E2 (% EWSR)	$B(\text{M1})$ (μ_0^2)	Range of $B(\text{M1})$ (MeV)
^{181}Ta	120 ± 150	12.8 ± 4.8	7.6–9.6
^{208}Pb	110 ± 80	26.7 ± 4.0	7.5–9.5
^{209}Bi	200 ± 90	None	—

Figure 4. $\sigma_{\gamma,n}$ and $\sigma_{\gamma,2n}$ for ^{181}Ta . The curve E1 + E2 is the result of fitting the data with Lorentzian shapes representing E1 and E2 absorption.

Figure 5. $\sigma_{\gamma,n}$ and $\sigma_{\gamma,2n}$ for ^{208}Pb . The curve E1 + E2 shows the result of fitting the data with Lorentzian shapes representing E1 and E2 absorption.

Figure 6. $\sigma_{\gamma,n}$ and $\sigma_{\gamma,2n}$ for ^{209}Bi . The curve E1 + E2 is the result of fitting the data with Lorentzian shapes representing E1 and E2 absorption.

5. Conclusions

We have shown that when photoneutron cross sections are fitted with Lorentzian shapes representing both E1 and E2 absorption, the experimental data are well reproduced by the fit in the full energy region covered by the data. This is an important result, since the assumption of pure E1 absorption leaves in all cases the high-energy part of the measured cross section unexplained, because the experimental points lie considerably above the Lorentzian shapes that represent the E1. It would be interesting to fit all available photonuclear cross sections under the assumption of E1 and E2 absorption, because this could yield at least the gross features of isovector E2 absorption in nuclei and a more accurate representation of the E1.

Table 4. Parameters for E1.

Nucleus	α_p (mb)	E_p (MeV)	Γ (MeV)	Reference
^{181}Ta	259	12.30	2.43	[19]
	341	15.23	4.48	
	283 ± 9	12.36 ± 0.04	2.58 ± 0.13	This work
	336 ± 10	15.35 ± 0.06	4.01 ± 0.13	
^{208}Pb	491	13.46	3.90	[19]
	486 ± 3	13.55 ± 0.02	4.00 ± 0.04	This work
^{209}Bi	521	13.45	3.97	[19]
	516 ± 4	13.45 ± 0.01	3.89 ± 0.03	This work

The isoscalar and isovector E2 intensities obtained are consistent with depletion of one EWSR in heavy nuclei.

The result obtained for $B(M1)$ in ^{208}Pb is consistent with the value observed for ^{206}Pb [16] and also with theoretical predictions [17]. For ^{209}Bi our result indicates that the M1 intensity lies mostly below the neutron emission threshold. For ^{181}Ta the obtained $B(M1)$ is about one half of that observed for both ^{206}Pb [16] and ^{208}Pb , but it could well suggest that part of the M1 strength lies below the (γ, n) threshold.

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Proton emission cross sections of silicon isotopes

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Measurements of the electrodisintegration cross sections, (e,p) and $(e,2p)$, and their corresponding electrodisintegration-plus-photodisintegration yields for ^{29}Si and ^{30}Si are presented. The experiment covers the range from 20 to 60 MeV electron incident energies. The experimental results were analyzed using $E1$ and $E2$ distorted-wave Born approximation virtual photon spectra for a finite nucleus in order to obtain the corresponding photodisintegration cross sections for both channels and isotopes. For ^{29}Si the integrated (γ,p) cross sections is 47% bigger than (γ,n) while for ^{30}Si they are of the same order of magnitude. For both isotopes the (γ,p) decay channel is dominantly $E1$. The $(\gamma,2p)$ channel is much smaller than the $(\gamma,2n)$ channel but the $E2$ contribution to this decay mode is of the same magnitude of the $E1$.

I. INTRODUCTION

Giant resonances are a common feature of all nuclei and have been extensively studied in the last 50 years. Yet there are few measurements of multiple charged particle emission in the decay process. Measurements carried out by Dodge *et al.*¹ for (e,p) and (e,α) cross sections in ^{90}Zr and ^{92}Zr , by Tamae *et al.*² for $^{90}\text{Zr}(e,\alpha)$ and by Tanaka *et al.*³ for $^{65}\text{Cu}(e,\alpha)$ show a common feature. The corresponding photodisintegration cross sections, obtained through the virtual photon analysis, showed a nonresonant behavior, steadily increasing for increasing bombarding energies. Since it has been shown that DWBA virtual photon spectra for finite nuclei relate electrodisintegration and photodisintegration cross sections, showing good agreement between measured photodisintegration cross sections and those obtained from electrodisintegration measurements,^{4,5} a failure of the virtual photon calculations to explain this strange behavior can be disregarded. Dodge *et al.*¹ pointed out that a possible explanation could be multiple alpha and proton emission, which would be interpreted as single events in all these experiments, since they detected, on line, the number of emitted charged particles. This explanation would imply that, as an example, the $2p$ decay channel should have a magnitude of about half of that of the one-proton decay channel. Furthermore, when the $^{65}\text{Cu}(e,\alpha)$ cross section was measured by residual activity⁶ the results differed from those of Ref. 3 and the obtained (γ,α) cross section showed a resonant behavior.

The lack of information, in the literature, on the decay of the giant resonance by $2p$ emission, motivated us to measure the (e,p) and $(e,2p)$ cross sections in ^{29}Si and ^{30}Si , to obtain their corresponding photodisintegration cross sections. These isotopes were chosen because they allow the measurements to be carried out by counting residual activity, thus distinguishing completely both decay channels. Apart from the interest in obtaining the relative magnitudes of the decay by one and two proton emission, it is important to study the fraction of the dipole strength carried out by proton emission.

The decay of the dipole giant resonance by neutron emission has already been measured for these isotopes⁷ and it was found that this mode carries about half of one classical dipole sum. Since for silicon isotopes the Coulomb barrier is about 2.4 MeV for each proton emitted,⁸ their (γ,p) cross sections are expected to be responsible for a non-negligible fraction of the giant resonance decay.

II. THE EXPERIMENT

The experiment was performed using the 60 MeV linear electron accelerator of the University of São Paulo. The electrodisintegration cross sections and the electrodisintegration-plus-photodisintegration yields were obtained by bombarding natural silicon targets and measuring, off line, their residual activity, counting the gamma rays emitted by their reaction products. The charge was obtained by using a Faraday cup for the electrodisintegration measurements and a ferrite monitor (calibrated to the Faraday cup) for the radiator-in measurements. Table I summarizes the reactions studied, their threshold energies, the energy and intensity of the detected gamma rays and the half-lives of the corresponding reaction products.

In order to allow enough time for the activity of reaction products to die out, measurements were carried out using four targets with 22 mg/cm² for electrodisintegration and three targets with 41 mg/cm² for electrodisintegration-plus-photodisintegration yields. For the activation of the electro plus photodisintegration yields a tantalum radiator with 275 mg/cm² was placed in the electron beam immediately ahead of the target.

Since the experiment was performed using natural silicon targets (92.2% ^{28}Si , 4.7% ^{29}Si , and 3.1% ^{30}Si),⁹ the bombarding and counting times had to be adjusted in order to allow measurements of all reaction products by residual activity. An irradiation time of four hours was used for the reactions $^{30}\text{Si}(e,2p)$ and $^{30}\text{Si}(e+\gamma,2p)$ and the residual activity was counted for 24 hours. The other three reactions (see Table I), which yield half-lives of a few minutes, were simultaneously measured using bom-

TABLE I. Main features of the (γ, p) and $(\gamma, 2p)$ reactions studied in this work. The threshold (Ref. 13) does not include the Coulomb barrier (Ref. 8) (about 2.4 MeV for each proton emitted). The other data were taken from Ref. 9.

Reaction	Decay product	E_γ (keV)	I_γ (%)	Half-life	Threshold (MeV)
$^{29}\text{Si}(e, p)^{28}\text{Al}$					
$^{29}\text{Si}(e + \gamma, p)^{28}\text{Al}$	$^{28}\text{Al} \xrightarrow{\beta^-} ^{28}\text{Si}$	1778.7	100	2.24 min	12.3
$^{29}\text{Si}(e, 2p)^{27}\text{Mg}$					
$^{29}\text{Si}(e + \gamma, 2p)^{27}\text{Mg}$	$^{27}\text{Mg} \xrightarrow{\beta^-} ^{27}\text{Al}$	843.8	73	9.46 min	21.9
$^{30}\text{Si}(e, p)^{29}\text{Al}$					
$^{30}\text{Si}(e + \gamma, p)^{29}\text{Al}$	$^{29}\text{Al} \xrightarrow{\beta^-} ^{29}\text{Si}$	1273.2	89	6.56 min	13.5
$^{30}\text{Si}(e, 2p)^{28}\text{Mg}$					
$^{30}\text{Si}(e + \gamma, 2p)^{28}\text{Mg}$	$^{28}\text{Mg} \xrightarrow{\beta^-} ^{28}\text{Al}$	30.6	95	20.93 h	24.0

bombarding times of 15 min. and counting times of about 20 min.

The gamma rays listed in Table I were detected using two high purity germanium detectors. One of them was used for the 30.6 keV gamma ray resulting from the decay products of the reactions $^{30}\text{Si}(e, 2p)$ and $^{30}\text{Si}(e + \gamma, 2p)$, which were studied using longer bombarding times, and the other for the 843.8, 1273.2, and 1778.7 keV gamma rays corresponding to the remaining studied

reactions, which used shorter bombarding times. Both detector efficiencies were measured with calibrated standard gamma-ray sources, in the same fixed geometry used to detect gamma rays from the targets, thus avoiding the need to measure absolute efficiencies. In the counting geometry used the distance between target and detector was 4 cm. The data acquisition system consisted of an amplifier, and analog-to-digital converter and a computer-aided measurement and control system connected with a PDP 11/84 computer.

In Figs. 1 to 4 the results obtained for the electrodisintegration cross sections (solid circles) and for the electrodisintegration-plus-photodisintegration yields (solid squares) are presented. The error bars show the statistical uncertainties of the measurements when these are

FIG. 1. The measured $\sigma_{\gamma,p}(E_0)$ (solid squares) and yield $Y_{e+\gamma,p}(E_0)$ (solid circles) for ^{29}Si as a function of the incident electron energy E_0 (right-hand scale). The $Y_{e+\gamma,p}(E_0)$ was obtained when a 275 mg/cm² tantalum radiator was placed in the electron beam ahead of the target. The histogram represents the photodisintegration cross section, $\sigma_{\gamma,p}$ (left-hand scale). The smooth curves represent the best simultaneous fit to the electrodisintegration cross section [Eq. (1)] and the electrodisintegration-plus-photodisintegration yield [Eq. (2)] that originated the histogram (see text).

FIG. 2. $\sigma_{\gamma,2p}$ for ^{29}Si . The dashed portion of the histogram represents the E2 contribution to the $(\gamma, 2p)$ cross section. See caption of Fig. 1.

FIG. 3. $\sigma_{e,p}$ for ^{30}Si . See caption of Fig. 1.

larger than the points representing the measured cross sections or yields. The overall uncertainty in the absolute scale is about 20%.

II. DATA ANALYSIS AND RESULTS

The measured cross section and yields were analyzed using the virtual photon method which relates electrodisintegration and photodisintegration cross sections. According to this method, the relationship between the photodisintegration cross section $\sigma_{\gamma,x}$ and the measured electrodisintegration cross section $\sigma_{e,x}$ is

$$\sigma_{e,x}(E_c) = \int_0^{E_c} \sum_{\lambda l} \sigma_{\gamma,x}(E) N^{\lambda l}(E_c, E, Z, A) \frac{dE}{E} \quad (1)$$

where E_c is the electron kinetic energy; E is the photon

energy; $N^{\lambda l}$ is the virtual photon spectrum calculated for a nucleus with atomic number Z and mass number A (Ref. 10) and λl is the multipolarity. The yield with the radiator in is

$$Y_{e+\gamma,x}(E_c) = \sigma_{e,x}(E_c - 2\Delta) + N_r \int_0^{E_c - \Delta} \sigma_{\gamma,x}(E) N^{\text{BR}}(E_c - \Delta, E, Z_r) \frac{dE}{E}, \quad (2)$$

where Δ is the electron energy loss in half the radiator thickness, N_r is the number of atoms/cm² of the radiator, and N^{BR} is the bremsstrahlung cross section for a radiator of atomic number Z_r .

Photodisintegration produced by bremsstrahlung in the target was calculated and found to be negligible when compared with the measured electrodisintegration cross sections or the electrodisintegration-plus-photodisintegration yields.

The virtual photon spectrum used¹⁰ was calculated in the DWBA taking into account the size of the nucleus. The bremsstrahlung spectrum used is that calculated by Seltzer *et al.* and discussed in Ref. 11.

Figure 5 shows, for an electron incident energy of 30 MeV, the E1 and E2 virtual photon spectra calculated for ^{29}Si and the bremsstrahlung spectrum calculated for the Ta radiator used.

The photodisintegration cross sections were represented by histograms in Eqs. (1) and (2) and obtained by fitting simultaneously the electrodisintegration and electrodisintegration-plus-photodisintegration measurements, for each of the studied reactions, using the VIRLIB code.¹² It was assumed that only E1 and E2 multipoles contribute to the photonuclear processes in the energy region covered by the measurements. For the (γ, p) cross sections, in both isotopes, it was not possible to fit an E2

FIG. 4. $\sigma_{e,2p}$ for ^{30}Si . The dashed portion of the histogram represents the E2 contribution to the $(\gamma, 2p)$ cross section. See caption of Fig. 1. FIG. 5. The bremsstrahlung (BR) and the virtual (E1 and E2) photon spectra calculated for a kinetic electron energy of 30 MeV. The DWBA virtual photon spectra are those calculated by Onley *et al.* (Ref. 10), while the bremsstrahlung spectrum used is that calculated by Seltzer *et al.* (Ref. 11) (see text).

TABLE II. Integrated cross sections for (γ, p) and $(\gamma, 2p)$ reactions on silicon isotopes.

Reaction	EL	χ^2 red.	Integrated cross section (MeV mb)	Energy range (MeV)	% of sum- rule
$^{29}\text{Si}(\gamma, p)$	E1	1.03	269.3 (40.1)	14.7–45.0	70.0 (9.2)
$^{29}\text{Si}(\gamma, 2p)$	E1	0.85	5.8 (2.1)	26.6–60.0	1.3 (0.5)
$^{29}\text{Si}(\gamma, 2p)$	E2		2.5 (1.3)	30.0–40.0	15.3 (8.0)
$^{30}\text{Si}(\gamma, p)$	E1	0.76	150.8 (17.3)	15.9–40.0	33.7 (3.9)
$^{30}\text{Si}(\gamma, 2p)$	E1	1.84	0.9 (0.4)	28.7–55.0	0.20 (0.09)
$^{30}\text{Si}(\gamma, 2p)$	E2		0.6 (0.3)	33.5–50.0	2.6 (1.2)

component. The measurements were compatible with pure E1. For the $(\gamma, 2p)$ cross sections the inclusion of an E2 component resulted in a much better fit than that obtained under the assumption of pure E1.

In Figs. 1 to 4, the histograms show the obtained photodisintegration cross sections for each reaction. The E2 contribution to the $(\gamma, 2p)$ cross sections is shown by the dashed portions of the corresponding histograms. In these figures the full lines through the measured points show the results of the best simultaneous fits, which resulted in the presented histograms for each reaction. For all four reactions studied the compatibility of electrodisintegration and photodisintegration is excellent.

It has to be pointed out that the energy range of the fit was chosen in order to yield the best fit. Any other tentative fit to the $\sigma_{\gamma, p}$ or the $\sigma_{\gamma, 2p}$ below or above the energies shown in Figs. 1 to 4 impoverished the quality of the fit, producing cross sections compatible with zero.

Table II summarizes the obtained photonuclear cross sections, presenting for each channel and multipole the integrated cross section and the corresponding percentage of the sum rule exhausted.

IV. DISCUSSION

Figures 6 and 7 show the photoproton cross sections $(\gamma, p + \gamma, 2p)$ and the photoneutron cross sections $(\gamma, n + \gamma, pn)$ for ^{29}Si and $(\gamma, n + \gamma, pn + \gamma, 2n)$ for ^{30}Si (Ref. 7). Table III shows the integrated cross sections for photoprotons and photoneutrons and the sum of these decay channels for both isotopes. The results show that the photoproton decay channel is as important as the photoneutron channel, being bigger than the latter for ^{29}Si and smaller for ^{30}Si . The sum of photoproton and photoneutron decays exhausts about one E1 sum.

Comparing Figs. 6 and 7 one sees that at the peak position (~ 20 MeV) (γ, p) is bigger than (γ, n) for ^{29}Si , while for ^{30}Si it is the reverse. The decay mechanism certainly depends on the number of available levels in the residual nucleus (density of final states). It happens that for ^{29}Si , at an excitation energy of 20 MeV, the residual nucleus reached through one-proton emission (^{28}Al) has about 180 levels⁹ in the energy interval from ground state to 7.7 MeV, while the nucleus reached through one-neutron emission (^{28}Si) has only 50 levels in the energy interval from ground state to 11.5 MeV of excitation energy. For

FIG. 6. The photoproton $(\gamma, p + \gamma, 2p)$ cross section obtained in this work (histogram) and the photoneutron $(\gamma, n + \gamma, np)$ cross section (Ref. 7) (circles) for ^{29}Si .

FIG. 7. The photoproton $(\gamma, p + \gamma, 2p)$ cross section obtained in this work (histogram) and the photoneutron $(\gamma, n + \gamma, np)$ cross section (Ref. 7) (circles) for ^{30}Si .

TABLE III. Results obtained for $\sigma_{\gamma,p}$ (this work) and $\sigma_{\gamma,n}$ (Ref. 7) for silicon isotopes.

Reaction	Threshold (MeV)	Integrated cross section (MeV mb)	% of sum rule (E1)
$^{29}\text{Si}(\gamma,p)$	12.3	269.3	62.0
$^{29}\text{Si}(\gamma,2p)$	21.9	5.8	1.3
$^{29}\text{Si}(\gamma,n)$	8.50	183.0	42.1
$^{29}\text{Si}(\gamma,n + \gamma,np + \gamma,p + \gamma,2p): 458.1 \text{ MeV mb}$			
$^{30}\text{Si}(\gamma,p)$	13.5	150.8	33.7
$^{30}\text{Si}(\gamma,2p)$	24.0	0.9	0.2
$^{30}\text{Si}(\gamma,n)$	10.6	181.0	40.4
$^{30}\text{Si}(\gamma,2n)$	19.1	67.0	15.0
$^{30}\text{Si}(\gamma,n + \gamma,np + \gamma,2n + \gamma,p + \gamma,2p): 399.7 \text{ MeV mb}$			

^{30}Si , at 20 MeV of excitation energy, the situation is reversed, since there are three times more levels in the residual nucleus that can be reached by neutron decay (^{29}Si) than in that reached by proton decay (^{29}Al).⁹

In both isotopes the $(\gamma,2p)$ cross section is much smaller than the (γ,p) . The $(\gamma,2p)$ integrated cross section for ^{29}Si is five times bigger than that for ^{30}Si . For ^{30}Si the $(\gamma,2p)$ threshold is 4.9 MeV higher than the $(\gamma,2n)$, while for ^{29}Si the $(\gamma,2p)$ threshold is 3.8 MeV lower than the $(\gamma,2n)$. These differences in threshold may be re-

sponsible for the large differences in the $(\gamma,2p)$ integrated cross sections for ^{29}Si and ^{30}Si . Yet the $(\gamma,2p)$ cross sections in both isotopes present a common feature, they have an $E2$ component located in the energy range of the isovector $E2$ and they both exhaust a larger fraction of the $E2$ isovector sum than of the $E1$ sum.

While for both isotopes $E1$ dominates the one proton emission, $E2$ has an important participation in the two proton emission channel. Our results, however, cannot exclude a small $E2$ component in the one proton emission channel.

Both (γ,p) and $(\gamma,2p)$ cross sections present a resonant behavior and, because $(\gamma,2p)$ is much smaller than (γ,p) , even in an experiment which would measure all protons and interpret all $2p$ events as two one- p events the resulting photoproton cross section would still have a resonant shape. This indicates that either the $(\gamma,2p)$ cross section is more important for Zr isotopes than it is for Si isotopes, or the results obtained in Refs. 1 and 2 have an explanation other than multiple particle emission.

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One neutron photoemission cross section of ^{238}U

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The (e,n) cross section and the $(e+\gamma,n)$ yield have been measured for ^{238}U in the electron energy range 12–60 MeV. The experimental results have been analyzed using distorted-wave Born approximation $E1$ and $E2$ virtual photon spectra for a finite nucleus in order to derive the $E1$ and $E2$ components of the (γ,n) cross section. The $E1$ component exhausts 40% of the $E1$ energy weighted sum rule while the isoscalar $E2$ exhausts 70–80 % of its sum.

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I. INTRODUCTION

The decay of the isoscalar electric quadrupole resonance (IS $E2$) in ^{238}U has been extensively studied by means of several experiments using different methods, [1–4] but a few points still remain unclear. The integrated isoscalar $E2$ strength present in the two open channels (fission and one neutron emission) is below the expected value of about one $E2$ sum rule. The electron scattering coincidence $(e,e'f)$ experiments, by far the more sophisticated ones, are unfortunately subject to ambiguities in the separation of $E2$ and $E0$ components. The largest set of results comes from inclusive experiments, namely disintegration or fission induced by electrons or positrons, using the virtual photon analysis to derive the photonuclear cross sections [2–4]. It is important to notice that all the available data were analyzed using virtual photon spectra (VPS) calculated in DWBA (distorted-wave Born approximation) for a point nucleus or using *ad hoc* corrections for nuclear size effects based on plane-wave Born approximation models. The DWBA calculation for a finite nucleus from Zamani-Noor and Onley [5] has been shown to yield consistent results for heavy nuclei [6]. Since for heavy nuclei the difference between $E2$ spectra for finite and point nuclei is large even for electron energies around 10 MeV, the multipole decomposition performed in these previous experiments must be revised. In this paper we report the measurement of the ^{238}U (e,n) cross section, between 12 and 60 MeV, to extend the energy range of a previous experiment [3]. We also measured the electro-plus-photodisintegration yield (bremsstrahlung induced) to use as an additional constraint on the multipolarity assignments in the analysis. The data were analyzed with the virtual photon technique in order to separate the $E1$ and $E2$ contributions to the (γ,n) channel, using VPS calculated in DWBA for finite nuclei [5].

The experiment was performed using the residual activity technique, thus selecting uniquely the decay channel to be studied. This is one of the few situations where inclusive experiments can have advantage over exclusive ones, since $(e,e'n)$ measurements have severe difficulties in separating the neutrons arising from different channels (i.e., fission, fission products, etc.).

II. THE EXPERIMENT

The experiment was performed using the 60 MeV linear electron accelerator of the University of São Paulo. The electrodisintegration cross sections and the electro-plus-photodisintegration yields were obtained by bombarding very thin uranium targets (typically $150\text{ }\mu\text{g}/\text{cm}^2$ of ^{238}U) placed in a vacuum chamber and measuring, off line, their residual activity. The charge was integrated by a Faraday cup for the electrodisintegration measurements and a secondary emission monitor (calibrated to the Faraday cup) for the radiator-in measurements.

The cross section was obtained by measuring the activity of the 59.54 keV gamma-ray line from the 6.75 day β^- decay of ^{237}U , using a HPGe low energy photon spectrometer system. The detector efficiency was measured with a ^{241}Am calibrated gamma-ray source [7] in the same fixed geometry used to detect gamma rays from the targets. Since the 59.54 keV level of ^{237}Np is populated from both decays (^{237}U by β^- and ^{241}Am by α emission) with the same branching ratio, our efficiency measurement actually determined the product of the detector efficiency by the solid angle and the 59.54 keV gamma-ray line intensity, thus avoiding several sources of uncertainties. The data acquisition system consisted of an amplifier, an analog-to-digital converter, and a computer-aided measurement and control system connected to a PDP 11/84 computer.

In order to allow enough time for the activity of ^{237}U to die out, measurements were carried out using six UO_3 targets. For the activation of electro-plus-photodisintegration yields a tantalum radiator with $275\text{ mg}/\text{cm}^2$ was placed in the electron beam immediately ahead of the target.

The electrodisintegration results are shown in Fig. 1 by the full circles. The open circles are the results of a previous experiment by Martins *et al.* [3]. This data set had its energy scale corrected due to a small change in the energy calibration function of the São Paulo linac performed in 1979. After this correction the two data sets agree within 4% without any difference in shape. So, after normalization, the old data were used in the analysis in order to complete the low energy region. Also shown in Fig. 1 (by the triangles) are the (e,n) data of Shott

FIG. 1. Electrodisintegration cross sections: Martins *et al.* [3] (open circles), Shotter *et al.* [8] (triangles), this work (full circles). Electro-plus-photodisintegration yield: this work (squares).

FIG. 2. Electrodisintegration cross sections: Martins *et al.* [3] (open circles); this work (full circles). Electro-plus-photodisintegration yield: this work (squares). The lines represent the fits to the data.

et al. [8], without normalization. The electro-plus-photodisintegration data are also shown in Fig. 1 by the squares. The lowest energy point of this data set is at 25 MeV, well above the peak of the photonuclear cross section, so that the correction for energy losses of the beam in the radiator could be done with a very small contribution due to the tip of the spectrum, where the correction by the average energy loss is not valid.

The error bars show the statistical uncertainties when they are larger than the points. The overall uncertainty is 20%.

III. ANALYSIS AND RESULTS

The electrodisintegration cross section $\sigma_{e,x}(E_0)$ may be obtained from the photonuclear cross section $\sigma_{\gamma,x}(E)$ through a: integral over the virtual photon intensity spectrum $N^{\lambda L}(E_0, E, Z, A)$:

$$\sigma_{e,x}(E_0) = \int_0^{E_0 - m_e} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z, A) \frac{dE}{E} \quad (1)$$

In Eq. (1), E_0 stands for the total electron energy and E stands for the energy of the photon of multipolarity λL . In the same spirit, the yield with the radiator in is

$$Y_{e,x}(E_0) = \sigma_{e,x}(E_0 - 2\Delta E_0) + N_r \int_0^{E_0 - \Delta E_0 - m_e} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) K(E_0 - \Delta E_0, E) \frac{dE}{E} \quad (2)$$

where N_r is the number of nuclei/cm² in the tantalum radiator, $K(E_0, E)$ is the bremsstrahlung cross section for tantalum, and ΔE_0 is the electron energy loss in half of the radiator thickness.

The cross sections $\sigma_{e,x}(E_0)$ and yields $Y_{e,x}(E_0)$ measured in this work, plus the data of Martins *et al.* [3], shown in Fig. 2, were simultaneously fitted [9] in order to derive the $E1$ and $E2$ photonuclear cross sections. These

cross sections were represented by coarse histograms, thus not showing any detailed feature of their shapes. The fitting procedure used $E1$ and $E2$ VPS calculated in DWBA for a finite nucleus [5] in Eq. (1) and the bremsstrahlung cross section of Seltzer *et al.* [10] for $K(E_0, E)$ in Eq. (2). The lines shown in Fig. 2 represent the best fit (reduced $\chi^2 = 0.8088$ with 48 degrees of freedom) to the data. The resulting $E1$ and $E2$ cross sections are shown in Fig. 3 by the histograms, the hatched part corresponding to the $E2$ component. The $E1$ component exhausts $(41 \pm 5)\%$ of the $E1$ sum rule, while the $E2$ component exhausts $(86 \pm 22)\%$ of its energy weighted sum rule. It is important to note that, with this method, we are not able

FIG. 3. Photodisintegration cross sections derived by the VIRLIII code: $\sigma_{\gamma,e}^{E1}$ (histogram) and $\sigma_{\gamma,e}^{E2}$ (hatched histogram). Photodisintegration cross sections measured by monochromatic photons: Caldwell *et al.* [12] (open circles) and Dickey and Axel [11] (triangles). There is no normalization between the different sets of data.

to separate isoscalar and isovector $E2$, thus interpreting all the $E2$ strength found as isoscalar. Also shown in Fig. 3 (by the points) is a combination of (γ, n) results obtained from different laboratories using monochromatic photons from positron annihilation in flight. The agreement is quite good and we must emphasize that there is no normalization between the different sets of data.

As an alternate procedure, we analyzed our data employing the (γ, n) cross sections obtained with monochromatic photons as a shape constraint. This was done assuming that the total photodisintegration cross section can be written as $\sigma_{\gamma, n} = \sigma_{\gamma, n}^{E1} + \sigma_{\gamma, n}^{E2(T=0)} + \sigma_{\gamma, n}^{M1}$. So we can rewrite Eq. (1) as

$$\sigma_{\gamma, n}(E_0) = \int_0^{E_0 - m_n} [K_1 \sigma_{\gamma, n}(E) N^{(E1)} + K_2 \sigma_{\gamma, n}^{E2}(E) (N^{E2} - N^{E1}) + K_3 \sigma_{\gamma, n}^{M1}(E) (N^{M1} - N^{E1})] \frac{dE}{E}, \quad (3)$$

where K_i are constants to be determined from the fit. For the $\sigma_{\gamma, n}$ cross section we used a combination of the Dickey and Axel [11] data (between 6.5 and 7.8 MeV) and the Caldwell *et al.* [12] data, in the intervals 6.1–6.4 MeV and 7.9–18.7 MeV. We decided to use the cross section measured in Livermore [12] instead of Saclay [13] because in the work of Wolyneć *et al.* [14] about the differences of the cross sections from these two laboratories, it is shown that the Saclay measurements may have problems in the multiplicity sorting procedure. The (γ, n) cross section was measured up to 18.7 MeV only, so we had to make an *ad hoc* assumption about its behavior above this energy. We made a linear extrapolation, from the average value of 26.5 mb at 18.7 MeV to zero at 25 MeV. This is a reasonable assumption, since the integrated $\sigma_{\gamma, n} + \sigma_{\gamma, pn}$ cross section from 18.3 to 140 MeV should be around 136 mb MeV [15] (which gives an average cross section of 1.1 mb in this region). The combined cross section is shown in Fig. 3.

In Eq. (3), K_1 is called "normalization constant" because it allows to compensate for differences in the absolute scales of the photodisintegration and electrodisintegration cross sections. The value of K_1 is determined independently, by means of Eq. (2) and the radiator-in data:

$$K_1 = \left(\frac{Y_{e, n}(E_0) - \sigma_{e, n}(E_0 - 2\Delta E_0)}{N_r \int_0^{E_0 - m_n - \Delta E_0} \sigma_{\gamma, n}(E) K(E_0 - \Delta E_0, E) dE/E} \right). \quad (4)$$

The average value of K_1 (over all the radiator-in data) is 0.96 ± 0.08 , showing a very good agreement between the sets of data.

To represent the other cross sections in Eq. (3) we followed the systematics of the literature. For the isoscalar $E2$ we used a Lorentz line shape with peak at 10.5 MeV ($E_p = 10.5$ MeV), 3.0 MeV wide and with 8.25 mb at the peak ($\sigma_p = 8.25$ mb), thus exhausting one energy weight-

ed sum rule. By fitting our data with Eq. (3), K_2 then gives the fraction of the energy weighted sum exhausted by the isoscalar $E2$ cross section. For the $M1$ cross section we assumed a uniform strength of 1 mb distributed between 6.1 and 7.1 MeV, since the $M1$ strength, if present, should be distributed over several individual levels.

Fitting our data keeping K_1 fixed at the value determined by Eq. (4), we obtained $K_2 = 0.75 \pm 0.09$ and K_3 compatible with zero, with a reduced χ^2 of 1.159 (with 43 degrees of freedom). The fit is visually identical to the one shown in Fig. 2. As in the previous analysis, any attempts to include an isovector $E2$ component as a free parameter to be determined by the fit had to be given up because they yielded uncertainties larger than the parameters. We decided to estimate an upper limit for the isovector $E2$ component, assuming a Lorentz line shape with peak position at 21.0 MeV and 5.0 MeV wide. To exhaust one isovector $E2$ energy weighted sum, this cross section should have $\sigma_p = 27.9$ mb. This is clearly incompatible with the total $\sigma_{\gamma, n}$ cross section measured with monochromatic photons, which has a strength of 26.5 mb at 18.7 MeV and should be dropping. From our linear extrapolation, at 21 MeV the cross section should be 8.5 mb high. So we tried a fit fixing $\sigma_p^{E2(T=1)} = 5.6$ mb, thus exhausting 20% of the energy weighted sum. In this fit (still keeping $K_1 = 0.96$) we got $K_2 = 0.68 \pm 0.09$ and K_3 again compatible with zero. The reduced χ^2 of the fit was 1.039 with 43 degrees of freedom.

It is reassuring to notice that the isoscalar $E2$ strength determined is not strongly dependent either on the method of analysis or on the presence of an isovector $E2$ component.

Figure 4 summarizes our results for the isoscalar $E2$ resonance and shows a comparison with other results available. Both our results agree well with the strength predicted by the quasiparticle random phase approximation (QRPA) calculation [16] but, as expected, the information on the shape is not very relevant. The QRPA cal-

FIG. 4. Isoscalar $E2$ cross sections. QRPA prediction [16] (histogram), $\sigma_{E2}^{[1]}$ [1] (points), $\sigma_{E2}^{[2]}$ derived using the VIR13 CODE [this work—dashed histogram], and $\sigma_{E2}^{[3]}$ (Lorentz line shape—this work, see text).

ulation, shown by the histogram, exhausts 68.5% of the $E2$ energy weighted sum rule (EWSR), very close to both our results. The points in Fig. 4 represent the $E2$ component of $\sigma_{\gamma,f}$ obtained from the $E2$ strength of Ref. [1]. The transformation of the strength function $dB/d\omega$ to cross section was done by means of the following expression:

$$\sigma(EL, \omega) = \frac{8\pi^3 \alpha}{(\hbar c)^2 L^2 - 2} \frac{(L+1)\omega^{2L-1}}{L[(2L+1)!!]^2} \frac{dB}{d\omega}(EL, \omega),$$

where EL stands for electrical transitions of order L , α is the fine structure constant, $\hbar$ is Planck's constant, and c the velocity of light.

Our results corroborate the $E2/E0$ strength separation for ^{238}U suggested by the authors of Ref. [1]. They suggested that the $E2$ component should dominate the strength distribution up to 12 MeV, with an integrated strength exhausting $(19 \pm 2)\%$ of the $E2$ EWSR. Our result with the Lorentz line shape exhausts $(75 \pm 9)\%$ of the $E2$ EWSR, meaning that the two dominant decay channels (fission and one neutron emission) exhaust $(94 \pm 9)\%$ of the $E2$ EWSR, very close to the total exhaustion expected in this mass region.

IV. CONCLUSIONS

The absolute cross sections for one neutron emission by absorption of real and virtual photons have been measured in ^{238}U up to 60 MeV. The $E1$ and $E2$ photonuclear cross sections have been derived from these data. The multipole decomposition was done by two different

methods, yielding compatible results.

The one neutron emission channel is the dominant decay mode for the isoscalar $E2$ resonance, exhausting around 70–80% of the energy weighted $E2$ sum rule. The $E1$ component decaying through this channel exhausts 40% of its sum. Our analysis procedure was unable to detect any $M1$ contribution to the cross section and could only define an upper limit for the isovector $E2$ component of about 20% of its energy weighted sum.

The experimental results obtained in this work are in good agreement with previous measurements, but the photonuclear cross sections derived from them changed somewhat due to the use of virtual photon spectra calculated for finite size nuclei in the analysis. Our results, together with the results for the fission channel available from the literature, indicate that the two dominant channels for the decay of the isoscalar $E2$ resonance (fission and one neutron emission) exhaust about 90–100% of the $E2$ sum, as expected.

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Two proton and two neutron photoemission cross sections of ^{63}Cu

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The $(e,2p)$ and $(e,2n)$ cross sections and the $(e+\gamma,2p)$ and $(e+\gamma,2n)$ yields have been measured for ^{63}Cu in the electron energy range 28–60 MeV. The experimental results have been analyzed using distorted-wave Born approximation $E1$ and $E2$ virtual photon spectra in order to derive the $E1$ and isovector $E2$ components of the $(\gamma,2p)$ and $(\gamma,2n)$ cross sections. In the $2n$ channel the $E1$ component exhausts 12.4(11)% of the $E1$ energy weighted sum, while the isovector $E2$ component exhausts 52(15)% of its sum. The $2p$ channel exhausts 1.03(8)% of the energy weighted sum for $E1$ and we were able to establish an upper limit for the isovector $E2$ contribution of 0.58% of its sum.

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I. INTRODUCTION

The decay of the giant resonances through the multiple emission of charged particles is poorly known in the literature. Measurements carried out by Tanaka *et al.* [1] for the (e,α) cross sections in ^{63}Cu and ^{65}Cu , by Dodge *et al.* [2] for the (e,p) and (e,α) cross sections in ^{90}Zr and ^{92}Zr , and by Tamae *et al.* [3] for $^{90}\text{Zr}(e,\alpha)$ show a common and intriguing feature. The corresponding photodisintegration cross sections, obtained through a virtual photon analysis, show a nonresonant behavior, steadily increasing for increasing bombarding energies. A failure of the virtual photon calculations can be disregarded as a possible explanation for this behavior, since photodisintegration cross sections obtained from electrodisintegration data through the virtual photon analysis showed good agreement with cross sections measured with real photons [4,5]. Dodge *et al.* [2] pointed out that a possible explanation for this increase could be multiple alpha or proton emission, which would be interpreted as multiple single events in all these experiments, since the charged particles were detected, on line, by a system unable to distinguish different multiplicities. This explanation would imply that, as an example, the $2p$ decay channel should have a magnitude of about one-half of that of the one-proton decay channel. In a previous experiment [6] we showed that, at least for the silicon 29 and 30 isotopes, this is not the case, since the $(\gamma,2p)$ cross section is much smaller than the (γ,p) . Furthermore, when the $^{63}\text{Cu}(e,\alpha)$ cross section was measured by residual activity [4], the results differed from those of Ref. [1] and the obtained (γ,α) cross section showed a peak, decreasing at higher energies.

Continuing our effort to increase knowledge of the multiple particle decay of the giant resonances, we measured the disintegration cross sections of ^{63}Cu by emission of two protons or two neutrons induced by electrons and bremsstrahlung photons in order to obtain the corresponding photonuclear cross sections by means of the virtual photon analysis. Since the measurement was done by counting the residual activity induced in the target, it completely distinguishes the decay channels studied.

II. EXPERIMENT

The experiment was performed using the 60-MeV linear electron accelerator of the University of São Paulo. The elec-

trodisintegration cross sections and the electro-plus-photodisintegration yields were obtained by bombarding 99.89% enriched ^{63}Cu targets [7] and measuring, off line, their residual activity, counting the gamma rays emitted by the reactions products. The charge was integrated using a Faraday cup for the electrodisintegration measurements and a ferrite monitor (calibrated to the Faraday cup) for the radiator-in measurements.

To obtain the $(e,2n)$ cross section, we measured the 67.4-keV gamma ray which follows the decay of $^{61}\text{Cu} \rightarrow ^{61}\text{Ni} + \beta^+$, with a half-life of 3.41(1) h [8]. For the $(e,2p)$ cross section, we measured the same 67.4-keV gamma ray, since ^{61}Co also decays to ^{61}Ni (by β^- emission), with a half-life of 1.650(5) h [8]. The 67.4-keV gamma rays were detected using a HPGe low-energy photon spectrometer system. The data acquisition system consisted of an amplifier, an analog-to-digital converter, and a computer-aided measurement and control system connected to a PDP-11/84 computer. The 67.4-keV line appeared in the spectrum isolated and over a low and smooth background, allowing an accurate determination of the photopeak area. As a check of our data, we followed the decay of the 67.4-keV gamma ray activity for 9 h after irradiation with 45-MeV electrons. Using the ^{61}Co and ^{61}Cu decay constants from the literature [8], we fitted our data with two exponentials, obtaining a reduced χ^2 of 1.07 (with 30 degrees of freedom). Fits with just one exponential or more than two exponentials were statistically unacceptable.

In order to allow enough time for the activity of reaction products to die out, measurements were carried out using seven different targets with 10.06(10) mg/cm². For the activation of electro-plus-photodisintegration yields, a 275-mg/cm² tantalum radiator was placed in the electron beam immediately ahead of the target. Irradiation times were of 3 h for both electrodisintegration and electro-plus-photodisintegration measurements. After irradiation at least seven spectra were taken, with times ranging from 30 min to 4 h each. From these data the contribution of each decaying nuclide to the initial activity was determined by a multiple linear regression [9].

The relative cross sections obtained in this experiment were normalized to an absolute measurement of the $(e,2n)$ cross section performed previously at this laboratory [4]. We

FIG. 1. Measured cross section $\sigma_{e,2n}(E_0)$ (squares, this work; triangles, Martins *et al.* [4]) and yield $Y_{e+\gamma,2n}(E_0)$ (circles) for ^{63}Cu as a function of the incident electron energy E_0 (right-hand scale). The histogram represents the photodisintegration cross section $\sigma_{\gamma,2n}(E)$ (left-hand scale). The dashed part of the histogram represents the isovector E2 contribution to the $(\gamma,2n)$ cross section. The smooth curves through the points represent the best simultaneous fit to the electrodisintegration cross section [Eq. (1)] and the electro-plus-photodisintegration yield [Eq. (2)] that originated the histogram (see text).

used three points in the interval covered by both experiments (28–32 MeV). The ratio of cross sections (old/new) varies by about $\pm 2\%$ from energy point to energy point, showing good compatibility between them. The normalization constant is set by the average of the ratio. The electro-plus-photodisintegration data were normalized by the same constant. The $(e,2p)$ cross section and the $(e+\gamma,2p)$ yield were normalized with a normalization constant obtained by the product of the $(e,2n)$ normalization constant times the ratio of the 67.4-keV line intensity from the decay of ^{61}Cu to that of ^{61}Co , since the line intensity was the only parameter changed from one measurement to the other.

The results obtained for the $(e,2n)$ cross section are shown in Fig. 1 (squares) together with the data of Martins *et al.* [4] (triangles) and the $(e+\gamma,2n)$ yield (circles). Both sets of $(e,2n)$ data (present work and from Ref. [4]) were used in the analysis in order to increase the energy range to be fitted and improve the accuracy of the derived $(\gamma,2n)$ cross section. The $(e,2p)$ cross section (squares) and the $(e+\gamma,2p)$ yield (circles) are shown in Fig. 2. The error bars show the statistical uncertainties of the measurements when these are larger than the points representing the measured cross section or yields. The overall uncertainty in the absolute scale is 25% (mainly due to the uncertainty in the 67.4-

FIG. 2. Cross sections for the $2p$ channel. See caption of Fig. 1.

keV line intensity from ^{61}Cu).

III. DATA ANALYSIS AND RESULTS

The electrodisintegration cross section $\sigma_{e,x}(E_0)$ may be obtained from the photonuclear cross section $\sigma_{\gamma,x}^{\lambda L}(E)$ through an integral over the virtual photon spectrum $N^{\lambda L}(E_0, E, Z)$:

$$\sigma_{e,x}(E_0) = \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z) \frac{dE}{E}. \quad (1)$$

In Eq. (1), E_0 stands for the total electron energy, E stands for the virtual photon energy of multipolarity λL , and m is the electron's rest energy. In the same spirit, the yield with the radiator in is

$$Y_{e,x}(E_0) = \sigma_{e,x}(E_0 - 2\Delta E_0) + N_r \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,x}^{\lambda L}(E) \times K(E_0 - \Delta E_0, E, Z_r) \frac{dE}{E}, \quad (2)$$

where N_r is the number of nuclei/cm² in the tantalum radiator, $K(E_0, E, Z_r)$ is the bremsstrahlung cross section in tantalum, and ΔE_0 is the electron energy loss in half the radiator thickness.

We have used the virtual photon spectra of Onley and co-workers [10,11], calculated with the distorted wave Born approximation (DWBA) assuming a point nucleus. We did not use the more realistic spectra calculated taking into account the finite size of the nucleus [12] because this would

TABLE I. Integrated cross sections for the reactions $(\gamma,2n)$ and $(\gamma,2p)$ in ^{63}Cu . The uncertainties shown arise from the statistical uncertainties of the data.

Reaction	EL	χ^2 (reduced)	Integrated cross section (MeV mb)	Energy range (MeV)	Fraction of sum rule (%)
$(\gamma,2n)$	E1	1.163	116(13)	19.7–55	12.4(11)
$(\gamma,2n)$	E2		20(6)	24–35	52(15)
$(\gamma,2p)$	E1	1.270	9.7(8)	17.2–50	1.03(8)
$(\gamma,2p)$	E2	upper limit	0.27	27–37	0.58

FIG. 3. Comparison between the $(\gamma, 2n)$ cross sections measured with monochromatic photons by Fultz *et al.* [15] (points) and obtained in this work (histogram).

demand a lot of computer time for very small corrections (about 1% for $E1$ and less than 5% for $E2$).

Photodisintegration produced by bremsstrahlung in the target was calculated and found to be negligible when compared with the measured electrodisintegration cross sections or the electro-plus-photodisintegration yields. The bremsstrahlung cross section used in Eq. (2) is that tabulated by Seltzer and Berger [13].

The photodisintegration cross sections were represented by histograms in Eqs. (1) and (2) and obtained by fitting simultaneously the electrodisintegration and electro-plus-photodisintegration measurements, for each of the studied reactions, using the VIRLIB code [14]. It was assumed that only $E1$ and $E2$ multipoles contribute to the photonuclear processes in the energy region covered by the measurements. For the $(\gamma, 2p)$ cross section, it was not possible to fit an $E2$ component. The measurements were compatible with pure $E1$, but we were able to establish an upper limit for an isovector $E2$ component within the uncertainties of our data. For the $(\gamma, 2n)$ cross section, the inclusion of an isovector $E2$ component resulted in a much better fit than that obtained under the assumption of pure $E1$.

In Figs. 1 and 2 the histograms show the obtained $(\gamma, 2n)$ and $(\gamma, 2p)$ cross sections, respectively. The isovector $E2$ contribution to the $(\gamma, 2n)$ cross section is shown by the dashed portion of the corresponding histogram. In both figures the solid lines through the experimental points show the results of the best simultaneous fits, which resulted in the

presented histogram for each reaction. In both cases the compatibility between the electrodisintegration and the radiator in data is excellent.

It has to be pointed out that the energy bins and range of the histograms were chosen in order to yield the best fit, without interfering with the reactions thresholds. Any other tentative fit below or above the energies shown in Figs. 1 and 2 impoverished the quality of the fit, producing cross sections compatible with zero. The upper limit for the isovector $E2$ contribution was evaluated by adding, to the result obtained from the best fit (pure $E1$), a small isovector $E2$ cross section (fixed), and recalculating the chi square. We increased the $E2$ contribution until the chi square exceeded the 95% confidence level limit. The isovector $E2$ strength at this point was considered the upper limit.

Table I summarizes the obtained photonuclear cross sections, presenting for each channel and multipole the integrated cross section and the corresponding fraction of the sum rule exhausted.

IV. DISCUSSION

Figure 3 shows a comparison between the $(\gamma, 2n)$ cross section measured with monochromatic photons at Livermore by Fultz *et al.* [15] (points) and the one obtained in this work (histogram). There is a good agreement in shape and an excellent agreement in the integrated strength up to 27.8 MeV (highest energy reached by the Livermore measurement): 80(8) MeV mb in Ref. [15] and 75(4) MeV mb in this work.

We calculated the branching ratio for the $2n$ channel, using the following expression:

$$R(E) = \frac{\sigma_{\gamma, 2n}(E)}{\sigma_{\gamma, n}(E) + \sigma_{\gamma, 2n}(E) + \sigma_{\gamma, 2p}(E) + \sigma_{\gamma, np}(E) + \sigma_{\gamma, 2np}(E)},$$

where E is the photon energy and $\sigma_{\gamma, n}, \sigma_{\gamma, 2n}, \sigma_{\gamma, np}, \sigma_{\gamma, 2np}$ are the cross sections measured by Fultz *et al.* [15] for the several open channels. The one-proton-emission channel has no data available, and the $2p$ channel was measured in this work. Table II shows the values for $R(E)$ up to 27 MeV [above this energy there were no data available for the (γ, n) cross section]. One can notice from Table II that the relative importance of the $2n$ channel increases with energy up to 27 MeV. This dominance of the $2n$ channel above 25 MeV could be an explanation for the high fraction of the isovector $E2$ sum rule strength exhausted by this channel, since the isovector $E2$ strength is concentrated in an energy region (24–35 MeV) where the emission of two neutrons is the most probable decay channel. The rest of the isovector $E2$ resonance strength should decay through the $1n$ and $1p$ channels, but we were not able to measure them in this experiment.

Unlike the $2n$ channel, which has an important contribution to the decay of both $E1$ and isovector $E2$ resonances, the $2p$ channel exhausts only 1% of the $E1$ energy weighted sum rule (EWSR) and less than 0.58% of the isovector $E2$ EWSR. This result corroborates the results obtained previously [6] for ^{29}Si and ^{30}Si , where the $2p$ channel had a small contribution to the $E1$ resonance decay (1.3% and 0.2%, respectively). Both $(\gamma, 2n)$ and $(\gamma, 2p)$ cross sections

TABLE II. Branching ratio for the $(\gamma, 2n)$ reaction.

E (MeV)	$R(E)$
19.7	0.17(2)
20.4	0.19(2)
21.0	0.24(2)
22.5	0.36(4)
24.0	0.29(5)
27.0	0.6(3)

obtained in this work present a maximum and then decline with increasing energy. Because of the small magnitude of ($\gamma, 2p$), even an experiment measuring protons and interpreting all $2p$ events as two $1p$ events would probably obtain a photoproton cross section showing a maximum. This indicates that the results obtained in Refs. [1–3] should have an explanation other than multiple particle emission for the steadily increasing behavior of the cross sections.

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NUCLEAR INSTRUMENTS & METHODS IN PHYSICS RESEARCH

Section A

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Tagged bremsstrahlung polarization: A Monte-Carlo approach

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Tagged bremsstrahlung polarization: A Monte-Carlo approach

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Abstract

We present a Monte-Carlo approach for the determination of the degree of polarization of tagged bremsstrahlung. Our method allows the evaluation of the Bethe–Heitler expression in the photon-oriented reference system, while defining the photon and electron collimators in the spectrometer coordinate system. This avoids the need to perform fourfold integrals of complicated and bulky expressions. We reproduced the published calculations and measurements, and studied the polarization behavior as a function of several parameters in different geometric conditions.

1. Introduction

The study of polarization observables of photonuclear reactions, with recent theoretical [1,2] and experimental [3,4] interest, increases the possibility of extracting fine details of nuclear structure significantly and thus, the understanding of the contribution of small amplitudes. Experiments in this area need intensive and quasi-monochromatic polarized photon beams with accurate knowledge of the degree of polarization.

For low-energy accelerators (5–30 MeV), the technically most simple way of obtaining linearly polarized photon beams consists in selecting the direction of the bremsstrahlung photons produced in a thin radiator by unpolarized electrons. CW accelerators allow quasi-monochromatic photon beams to be produced by the tagging method [5], and the degree of linear polarization and intensity can be enhanced, if the scattered electrons are also detected in an off-axis direction [6,7]. This is the so-called *double off-axis collimation method*.

The basis of the method is that there is an asymmetry in the photon polarization vector distribution, which is correlated to the photon and scattered electron emission angles. Therefore, although the bremsstrahlung beam is unpolarized as a whole, by selecting small solid angle

acceptances for photons and electrons, a polarized photon beam with a high degree of polarization may be obtained.

This method is simple in principle, but involves a rather large number of difficulties for practical realization, since the photon intensity and polarization spectra depend strongly on the photon emission angles [8]. As a result, the specific experimental conditions, such as position and dimension of electron and photon collimators and the alignment of the whole system (radiator, spectrometer, collimators and target), influence the photon beam characteristics strongly. On the other hand, because of the sharp behavior of the dependence, a detailed evaluation of the approximations used in the calculation is necessary. The complex behavior of the photon polarization as a function of the angular variables makes the interpretation of results obtained for the double collimation method rather complicated. Let us take, as an example, the case of photon and post-bremsstrahlung electron collimators opposing each other in relation to the incident electron beam. Moving the electron collimator to the same side of the photon causes the calculated polarization to change sign in Refs. [6, 11], but not in Ref. [8]. These references are in fact compatible; the apparent disagreement is due to the sensitive behavior of the polarization on differences in the conditions of the geometry (see Section 5). Another interesting behavior is that the polarization can change sign along the photon energy spectrum for the double collimation case.

A detailed study of polarization and intensity of tagged bremsstrahlung photons has been done in Ref. [8] in the framework of high-energy, small-angle approximation in

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the photon-oriented coordinate system. This analysis, however, cannot be applied directly to our conditions since it must be transformed (i) to the low electron energy regime, and (ii) to a real, space-extended collimation condition, which can be constructed only in the spectrometer coordinate system, with space-fixed axes.

The aim of this work is to carry out an analysis of the photon intensity and polarization behaviors for low incident electron energies (5–30 MeV) and space-extended photon and electron collimators, free of any approximation other than first-order Born.

The attempt to transform analytically the expression for the fivefold differential cross section [8] from the photon-oriented system to the spectrometer system and then perform a fourfold integration (over the acceptances of both collimators) leads to huge difficulties and ambiguities, arising mainly from the bulky character of the expressions.

A more reasonable way to do this is by Monte-Carlo simulation (MCS) in the photon system, transforming the collimators boundary conditions from the spectrometer system. This approach has some advantages:

- (i) Collimator shapes are defined in the spectrometer coordinate system.
- (ii) Cross sections are calculated in the photon coordinate system (thus, avoiding bulky expressions).
- (iii) Allows simple elimination of poles (see text further on).

2. The fivefold differential cross section

2.1. Notation

Here, we define the various quantities that are used in this work. We assume $\hbar = c = m_e = 1$ and, unless otherwise stated, energies are in units of the electron rest energy, $m_e c^2$, and momenta in units of $m_e c$:

e_1 and p_1	energy and momentum of the incident electron
e_2 and p_2	energy and momentum of the scattered electron
k and k	energy and momentum of the emitted photon
$q = p_1 - p_2 - k$	momentum transferred to the target nucleus
$u = p_{1\perp}$	component of p_1 perpendicular to k ($u = p_1 \sin \theta_1$)
$v = p_{2\perp}$	component of p_2 perpendicular to k ($v = p_2 \sin \theta_2$)
$\theta_1(\theta_2)$	angle between $p_1(p_2)$ and k
ϕ	azimuthal angle between the planes containing, respectively, p_1 and k , and p_2 and k

e	polarization vector of the photon
ψ	polarization angle, defined as the angle between the plane containing p_1 and k and the plane containing e and k ($\psi = 0$, when the two planes coincide)
$d\Omega_k$	solid angle in which the photon is emitted
$d\Omega_{p_1}$	solid angle in which the scattered electron is emitted
$d\sigma_\psi \equiv d^5\sigma_\psi$ $\equiv (d^5\sigma)/(d\Omega_{p_1} d\Omega_k dk)$	fivefold differential bremsstrahlung cross section for an angle of polarization ψ
$d\sigma_\perp \equiv d^5\sigma_\perp$ $d\sigma_\parallel \equiv d^5\sigma_\parallel$ $d\sigma_0 = d\sigma_\perp + d\sigma_\parallel$ $P = (d\sigma_\perp - d\sigma_\parallel)/(d\sigma_\perp + d\sigma_\parallel)$	value of $d\sigma_\psi$ for $\psi = \pi/2$ value of $d\sigma_\psi$ for $\psi = 0$ (or π) degree of polarization using as reference the production plane (k, p_1),
$P_\psi =$ $(d\sigma_\psi - d\sigma_{(\psi+\pi/2)})/(d\sigma_\psi + d\sigma_{(\psi+\pi/2)})$	degree of polarization using as reference an arbitrarily oriented plane regarding the production plane (ψ is the angle between the planes)

2.2. The bremsstrahlung cross section

The fivefold differential cross section for emission of polarized photons by unpolarized electrons, calculated in first-order Born approximation, can be written, according to Maximon et al. [8], in the form

$$\frac{d^5\sigma}{d\Omega_k d\Omega_{p_2} dk} = \frac{Z^2 e^2}{\hbar c} \left(\frac{e^2}{m_e c^2} \right)^2 \frac{[1 - F(q)]^2}{q^4} \frac{p_2}{k p_1 (2\pi)^2} \{ \}, \quad (1)$$

where Ze is the nuclear charge, $e^2/\hbar c = \alpha = 1/137$ is the fine structure constant, $(e^2)/(m_e c^2) = r_e$ is the classical electron radius, and $F(q)$ the atomic form factor. We use the Bethe-Heitler expression, as given by Maximon et al. in Ref. [8], in order to avoid losing details of the fine structure of the differential cross section and polarization that might occur because of cancellations [9].

Assuming a coordinate system where the z -axis coincides with the direction of the photon momentum (see Fig. 1), then

$$\{ \} = A \cos^2 \Psi + 2B \cos \Psi \sin \Psi + C \sin^2 \Psi, \quad (2)$$

where

$$A = 4k^2 e_1 e_2 (u - v)^2 \xi \eta + 16e_1^2 e_2^2 (u \xi - v \eta \cos \phi)^2 - 4q^2 (e_1 u \xi - e_2 v \eta \cos \phi)^2,$$

$$B = [-16e_1^2 e_2^2 (u \xi - v \eta \cos \phi) v \eta$$

Fig. 1. Photon-oriented coordinate system: polar coordinate system, where the z -axis coincides with the photon momentum direction.

$$\begin{aligned}
 & + 4q^2(\varepsilon_1 u \xi - \varepsilon_2 v \eta \cos \phi) + \varepsilon_2 v \eta] \sin \phi, \\
 C = & 4k^2 \varepsilon_1 \varepsilon_2 (u - v)^2 \xi \eta + 16\varepsilon_1^2 \varepsilon_2^2 v^2 \eta^2 \sin^2 \phi \\
 & - 4q^2 \varepsilon_2^2 v^2 \eta^2 \sin^2 \phi, \\
 \xi = & \frac{1}{1 + u^2} - \frac{1}{2\varepsilon_1(\varepsilon_1 + p_1 \cos \theta_1)}, \\
 \eta = & \frac{1}{1 + v^2} - \frac{1}{2\varepsilon_2(\varepsilon_2 + p_2 \cos \theta_2)}, \text{ and} \\
 q^2 = & p_1^2 + p_2^2 + k^2 - 2p_1 k \cos \theta_1 + 2p_2 k \cos \theta_2 \\
 & - 2p_1 p_2 (\cos \theta_2 \cos \theta_1 + \sin \theta_2 \sin \theta_1 \cos \phi).
 \end{aligned}$$

The effect of screening, which can be significant in the low-energy region (low momentum transfer) for large collimator acceptance, was taken into account in this work through the atomic form factor $F(q)$, in accordance with [10, 11]

$$\frac{1 - F(Z, q)}{q^2} = \frac{0.1}{b_1^2 + q^2} + \frac{0.55}{b_2^2 + q^2} + \frac{0.35}{b_3^2 + q^2}, \quad (3)$$

where

$$\begin{aligned}
 b_1 = & \frac{1}{1.21} 6Z^{1/3}, \quad b_2 = \frac{1}{1.21} 1.2Z^{1/3}, \quad \text{and} \\
 b_3 = & \frac{1}{1.21} 0.3Z^{1/3}.
 \end{aligned}$$

3. Description of the computer code

The code was written in FORTRAN-77, version 5.3. As a first step, four independent values (0–1) are generated randomly. ϕ and ϕ_e are obtained by multiplying two of them by 2π , so that these angles will span the range $0-2\pi$ with a uniform distribution. The other two random numbers are used to determine θ_1 and θ_2 , by calculating the arccos, so that the angles will span the range $0-\pi$, with the distribution uniform in $\cos \theta$. Then the angles θ_1 , θ_2 and ϕ are recalculated in the laboratory system (see Fig. 2), as

$$\begin{aligned}
 \theta_1 &= \theta_1, \\
 \theta_e &= \arccos(\cos \theta_1 \cos \theta_2 + \sin \theta_1 \sin \theta_2 \cos \phi), \\
 \omega' &= \arccos\left(\frac{\cos \phi \sin \theta_2 - \sin \theta_1 \cos \theta_e}{\cos \theta_1 \sin \theta_e}\right), \\
 \phi_1 &= \omega' + \pi + \phi_e.
 \end{aligned} \quad (4)$$

As a result, we have 4 angles in the laboratory system: $\theta_1(0-\pi)$, $\theta_e(0-\pi)$, $\phi_e(0-2\pi)$, and $\phi_1(0-2\pi)$. Then these angles are recalculated in the spectrometer coordinate system (see Fig. 3), as angles:

Fig. 2. Laboratory coordinate system: Polar coordinate system, where the z -axis coincides with the direction of the momentum of the incoming electron.

Fig. 3. Spectrometer coordinate system: Cartesian coordinate system, where the z -axis coincides with the direction of the incoming electron and the x - y plane is parallel to the spectrometer entrance.

$$\begin{aligned}\chi_e^x &= \arcsin[\sin(\theta_e)\cos(\phi_e)], \\ \chi_e^y &= \arcsin[\sin(\theta_e)\sin(\phi_e)], \\ \chi_r^x &= \arcsin[\sin(\theta_r)\cos(\phi_r)], \\ \chi_r^y &= \arcsin[\sin(\theta_r)\sin(\phi_r)],\end{aligned}\quad (5)$$

The angles ϕ_e and ϕ_r are measured from the x -axis, see Figs. 2 and 3.

The boundary conditions are then analyzed:

- $\chi_{e,r}^{x_1, y_1} \leq \chi_{e,r}^{x, y} \leq \chi_{e,r}^{x_2, y_2}$ for the case of rectangular collimators, where $\chi_{e,r}^{x_1, y_1}, \chi_{e,r}^{x_2, y_2}$ are the collimators boundaries;
- $d^2(\tan \chi_e^x - \tan \chi_r^{x_0})^2 + d^2(\tan \chi_e^y - \tan \chi_r^{y_0})^2 \leq r^2$, for the case of a circular photon collimator, where $\chi_r^{x_0, y_0}$ define the center of the collimator, r its radius and d the distance from the radiator to the collimator wall.

If the angles $\chi_{e,r}^{x, y}$ are inside the boundary conditions, the corresponding initial angles θ_1, θ_2 , and ϕ are used to calculate $d\sigma$ and $d\sigma_1$, and the values

$$\langle d\sigma_0 \rangle = \frac{\sum_{i=1}^N (d\sigma_1 + d\sigma)_i}{N}, \quad (6)$$

$$\langle P \rangle = \sum_{i=1}^N \left[W_i \frac{C_i - A_i}{C_i + A_i} \right], \quad (7)$$

where the weight W_i is defined as

$$W_i = \frac{(d\sigma_1 + d\sigma)_i}{\sum_{i=1}^N (d\sigma_1 + d\sigma)_i}. \quad (8)$$

Here N is the number of events which are inside the boundary conditions. $\langle d\sigma_0 \rangle$ is the average fivefold differential cross section over the collimators areas and $\langle P \rangle$ is the average polarization.

Photon and electron solid angles are defined as $\Delta\Omega_k = 4\pi(N_k^{\text{MC}}/N_{\text{NOR}})$ and $\Delta\Omega_{p_2} = 4\pi(N_{p_2}^{\text{MC}}/N_{\text{NOR}})$, where N_k^{MC} is the number of Monte-Carlo events which are inside of the boundary conditions of the photon collimator, $N_{p_2}^{\text{MC}}$ is the same, but for the electron collimator, and N_{NOR} is the number of Monte-Carlo starts.

The photon flux is given by $N_\gamma = \langle d\sigma_0 \rangle \Delta\Omega_k \Delta\Omega_{p_2} \Delta k N_e N_n$, where N_e is the primary electron flux, N_n the radiator thickness in nuclei per cm^2 , and Δk the photon energy bin.

In the simulations done for point collimators (very small acceptances, 10^{-6} rad, to simulate the evaluation of the fivefold differential cross section for specific values of the angular variables), we used a procedure to compress the sorting close to the collimator acceptance in order to increase statistics. This compression procedure was done in such a way as to maintain the shape of the distribution in $\cos \theta$ intact. So, if, instead of sorting in the $[0, \pi]$ interval, we wanted an interval $[\theta_1, \theta_2]$, the following expression was used:

$$\theta = \arccos[\cos \theta_2 + (\cos \theta_1 - \cos \theta_2)\text{Random}],$$

where Random represents the random number generator subroutine. In such cases, we used a sorting interval three times wider than the acceptance.

Another feature of our code is to eliminate poles. They may occur for some values of variables θ_1 and θ_e , causing Eq. (4) to diverge. In those cases, where $\theta_1 = \pi/2$ or $\theta_e = 0$, the event is discarded without loss of the simulation.

4. Simulation results

To check the computer code, we reproduced polarization results calculated in Ref. [8] in the photon system. Events were generated in a small angular acceptance (10^{-6} rad) around the nominal values. Fig. 4 presents the results of such simulation. It shows that the polar angular dependence of the polarization has a complicated character, with a sharp dip for $\phi = 0$, in the region where $q \cong 0$. The simulation reproduces well the results of Ref. [8].

Results from Refs. [6,11] were also reproduced. Table 1 shows the result of a Monte-Carlo simulation (MCS) reproducing conditions of Ref. [11]. The simulation results for both cross section and polarization agree within statistical uncertainties with Ref. [11]. Fig. 5 compares, for the same conditions, results from MCS (full circles, joined by solid curves) and from Ref. [6]: calculation (dashed curves) and measurement (data points, open circles with error bars). For this comparison, the results from Ref. [6] had their signs inverted, since they use the spectrometer mid-plane as reference, which is perpendicular to the production plane. The agreement is quite good. We also reproduced the theoretical curves calculated according to the Bethe-Heitler formula

Table 1

Cross sections and polarizations for $\varepsilon_1 = 100$ MeV, $k = 50$ MeV, $Z = 78$ and $-8.61 \leq \chi_e^x \leq 1.61$ mrad, $-3.5 \leq \chi_e^y \leq 3.5$ mrad, $-25 \leq \chi_e^z \leq 25$ mrad and different acceptances χ_e^z

Case	χ_e^z (mrad)	$1/Z^2 d\sigma/dk$ (fm ² /MeV)		P (%)	
		This work	Ref. [11]	This work	Ref. [11]
"O"	0–10	$1.44(7) \times 10^{-3}$	1.456×10^{-3}	52.2(2.5)	48.7
	0–25	$1.61(7) \times 10^{-3}$	1.528×10^{-3}	49.5(2.5)	46.6
"S"	–10–0	$4.48(26) \times 10^{-4}$	4.373×10^{-4}	–11.0(5)	–11.0
	–25–0	$5.68(26) \times 10^{-4}$	5.532×10^{-4}	–17.9(5)	–18.4

Fig. 4. Polarization versus θ_2 for $\varepsilon_1 = 140$ MeV, $k = 95$ MeV, $\theta_1 = 1$. Upper part corresponds to $\phi = 0$, lower part to $\phi = 30$. In both figures, the points represent results of this work, while the curves represent results from Ref. [8].

presented in Ref. [12], for $\varepsilon_1 = 300$ keV, $\theta_1 = 35^\circ$, and $\theta_2 = 55^\circ$. Therefore, with our Monte-Carlo approach, we were able to reproduce results in a wide range of initial electron energies (300 keV–140 MeV) and angles of the emitted electron and photon (mrad–rad).

With the reliability of our code assessed, we proceeded to study the influence of several parameters on the polari-

Fig. 5. Polarization versus k/ε_1 for $\varepsilon_1 = 18.4$ MeV; $0.052 \leq \chi_e^x \leq 0.078$ mrad; $-0.25 \leq \chi_e^y \leq 0.25$ mrad; $-0.04 \leq \chi_e^z \leq -0.013$ mrad; and $-0.013 \leq \chi_e^z \leq 0.013$ mrad (case "O", upper part of the figure) and $0.052 \leq \chi_e^x \leq 0.078$ mrad; $-0.25 \leq \chi_e^y \leq 0.25$ mrad; $0.013 \leq \chi_e^z \leq 0.04$ mrad; and $-0.013 \leq \chi_e^z \leq 0.013$ mrad (case "S", lower part of the figure). Full circles joined by solid curves are the results from MCS; results from Ref. [6] are shown by the dashed curves (calculation) and open circles with error bars (measurement).

zation behavior of tagged bremsstrahlung, for 5 MeV initial electron energy, which corresponds to the first stage of the São Paulo Microtron.

Fig. 6 shows the polar angular dependence of the polarization, differential cross section and momentum transfer, calculated in the photon coordinate system, for two characteristic cases:

- $\phi = 0$, photon and electron point collimators are located on opposite ($\theta_2 > \theta_1$, case "O") and same ($\theta_2 < \theta_1$, case "S") sides relative to the initial electron beam (upper part of the figure).
- $\phi = 180^\circ$, $\theta_2 < \pi - \theta_1$, case "S", and $\theta_2 > \pi - \theta_1$, case "O" (lower part of the figure).

It can be seen from the figures that the polarization for cases "O" and "S" have the same sign (negative) and that case "O" is characterized by

- a minimum in the q^2 versus θ_2 function at $\theta_2 \sim \theta_{e,c}$;
- a peak in the cross section in the region of this minimum, and

Fig. 6. The polarization (squares), differential cross section (in arbitrary units, circles) and the square of the momentum transferred ($m_e^2 c^2$, triangles) as a function of θ_2 for $\varepsilon_1 = 5.1$ MeV, $k = 2$ MeV, $\theta_1 = 5^\circ$. In the upper part of the figure, $\phi = 0$ ($\theta_2 < 5^\circ$ corresponds to case "S", $\theta_2 > 5^\circ$ to case "O"). In the lower part of the figure, $\phi = 180^\circ$ (case "S"). The curves are just to guide the eye.

(c) a sharp dip in the polarization in the same region.

Fig. 7 presents polarization and cross section versus azimuthal angle ϕ for $\theta_1 = 0.1$ rad, $\theta_2 = 0.2$ rad, and $k/\varepsilon_1 = 0.2$ and 0.5 . It can be seen that, for these photon energies and ϕ close to 0 (case "O"), the polarization is negative (with a sharp dip, as discussed in Ref. [8]), then it changes sign, reaches its maximum positive value at $\phi \sim 0.2$ rad, equals zero at $\phi \sim \pi/2$ rad and reaches its peak negative value at $\phi \sim \pi$ rad (case "S"). It is clear that, for ϕ close to zero, the polarization sign will depend strongly on the ϕ -acceptance of the collimators. So the effect of change in signs in "O" and "S" cases, observed in Refs. [6] and [11], is not related to the changes in momentum transfer, but associated with the ϕ -acceptance of the collimators. The ϕ -dependence of the cross section shows two maxima at ϕ close to 0 and 2π , but presents a sharp dip at $\phi = 0$.

Fig. 7. Polarization versus azimuthal angle ϕ for $\varepsilon_1 = 5.1$ MeV, $Z = 13$, $\theta_1 = 0.1$ rad, $\theta_2 = 0.2$ rad and $k/\varepsilon_1 = 0.2$ (solid curve, squares) and 0.5 (dotted curve, circles). Also shown is the cross section (in arbitrary units) for $k/\varepsilon_1 = 0.5$ (dashed curve, triangles, see text). The curves are just to guide the eye.

Fig. 8 presents polarization as a function of photon energy for some angles ϕ close to zero. In this region, the polarization depends strongly on both the photon energy and the angle ϕ . The ϕ -dependence is particularly important because of the collimator dimensions. As the figure clearly shows, small variations in the collimator width can cause dramatic changes in the average polarization obtained.

In the next figures, we present MCS results for collimation conditions specified in the spectrometer coordinate system.

Fig. 9 presents the linear polarization, for the case of no restriction in the electron arm and a 10^{-6} rad photon acceptance, as a function of $\chi_{\gamma}^2/\theta_{\gamma c}$, for two photon energies. The polarization presents a maximum near the angle $\theta_{\gamma c}$ ($\theta_{\gamma c} = m_e/\varepsilon_1$ is the characteristic electrodynamic angle) in agreement with Ref. [11], as expected for the off-axis bremsstrahlung radiation.

Fig. 8. Polarization as a function of photon energy for $\varepsilon_1 = 5.1$ MeV, $\theta_1 = 0.1$ rad, $\theta_2 = 0.2$ rad and $\phi = 0.06$ rad (solid curve, squares), $\phi = 0.15$ rad (dashed curve, triangles), and $\phi = 0.20$ rad (dotted curve, circles). The curves are just to guide the eye.

Fig. 9. Polarization as a function of $\chi_\gamma^x / \theta_{\gamma c}$ for the case of no restriction in the electron arm for $\epsilon_1 = 5.1$ MeV, $Z = 4$ and $k/\epsilon_1 = 0.2$ (solid curve, squares) and 0.4 (dashed curve, triangles). The curves are just to guide the eye.

Fig. 10. Polarization as a function of photon energy, calculated with screening correction ($Z = 4$ squares, $Z = 82$ triangles) for $\chi_\gamma^x = \theta_{\gamma c} \pm 0.5\theta_{\gamma c}$ and $\chi_\gamma^y = 0 \pm 0.5\theta_{\gamma c}$ with $\epsilon_1 = 5.1$ MeV and no restriction in the electron arm. The curves are just to guide the eye.

Fig. 10 presents the polarization as a function of photon energy for $\chi_\gamma^x = \theta_{\gamma c}$, to illustrate the effect of screening (see caption). This case is calculated with no restriction in the electron arm, for two different radiators, $Z = 4$ and 82 (the result without screening correction coincides visually with the result for $Z = 4$). It is seen that a lower Z radiator is slightly better for obtaining a higher degree of polarization. It is important to notice that, even though the screening correction is applied as the same multiplicative factor in both σ and σ_1 , the degree of polarization is sensitive to this correction, because of the sum in the denominator of Eq. (7), which avoids the cancellation. Nevertheless, this effect is only relevant for extended collimators, being negligible in the case of small acceptances.

Fig. 11 presents the polarization as a function of photon energy for realistic electron and photon collimators for the São Paulo tagger (see caption). In this case, $\Delta\chi_\gamma^{x,y} = \theta_{\gamma c}$. We can conclude that (i) collimation of the scattered electrons produces a large increase in the

Fig. 11. Polarization as a function of photon energy for $\epsilon_1 = 5.1$ MeV. Dotted curve (circles) corresponds to $0.017 \leq \chi_\gamma^x \leq 0.07$ rad, $-0.14 \leq \chi_\gamma^y \leq 0.14$ rad, $-0.131 \leq \chi_\gamma^z \leq -0.044$ rad, and $-0.044 \leq \chi_\gamma^x \leq 0.044$ rad. This configuration corresponds to case "O", with selection of electrons in the upper half of the spectrometer acceptance and the photon collimator placed in the lower hemisphere. Solid curve (dark squares) corresponds to the same collimators dimensions, but with the electrons being selected in the lower half of the spectrometer, case "S". Dashed curve (triangles) corresponds to the same photon collimator but full spectrometer acceptance: $-0.07 \leq \chi_\gamma^x \leq 0.07$ rad; $-0.14 \leq \chi_\gamma^y \leq 0.14$ rad; $-0.131 \leq \chi_\gamma^z \leq -0.044$ rad; and $-0.044 \leq \chi_\gamma^x \leq 0.044$ rad. The curves are just to guide the eye.

polarization at the low-energy part of the spectrum, as compared with what would be expected without residual electron selection (Fig. 10) and (ii) polarization changes sign and is close to zero at the high-energy part of the spectrum. So this part cannot be used for practical purposes.

5. Conclusions

We have described a method to calculate the bremsstrahlung radiation cross section and linear polarization based on Monte-Carlo simulation. We checked our code by reproducing published calculations and measurements. The flexibility of the Monte-Carlo method allowed us to study the polarization behavior as a function of several parameters in different geometric conditions. This study showed that different published calculations [6,8,11], which presented apparently conflicting results, were compatible. The differences arose from the complex behavior of the polarization function and differences in the geometrical conditions.

Our analysis shows that a 5 MeV continuous electron beam can be used to obtain a linearly polarized tagged photon beam with a high degree of polarization.

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Experimental test of the bremsstrahlung cross section

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The bremsstrahlung cross section has been studied by measuring the activity induced in ^{63}Cu by electrodisintegration and when thin radiators of Cu, Mo, Ta, and Th were placed in the electron beam just ahead of the target. The electron energies were varied from 13.5 to 60 MeV for the electrodisintegration and from 20 to 60 MeV for the radiator-in measurements; the (γ, n) cross section for ^{63}Cu was determined using virtual photon theory; the radiator data were fitted using various bremsstrahlung cross sections. The best fit is obtained using the synthesized spectrum of Seltzer which differs from the Davies-Bethe-Maximon cross section as given by equation (3CS) of Koch and Motz.

INTRODUCTION

In the past ten years, since accurate virtual photon spectra calculated in the distorted wave Born approximation (DWBA) have been available, the multipole composition of photonuclear cross sections has been studied from measurements of both electrodisintegration cross sections and photodisintegration yields. The analysis technique relies on the strong multipole dependence of the virtual photon spectra, manifested through the electrodisintegration cross section, whereas the bremsstrahlung spectrum contains all multipoles.

In two previous experiments the (e, p) and (e, α) cross sections of several pf -shell nuclei were studied.^{1,2} In these experiments a tantalum radiator was placed ahead of the target and the photo-plus-electrodisintegration yields were used to determine the magnitude of the $E1$ and $E2$ photodisintegration cross sections of the nuclei under study. During the course of the data analysis it became apparent that the results obtained were extremely sensitive to the choice of bremsstrahlung cross section used. The present experiment was designed to study the bremsstrahlung cross section for incident electron energies between 20 and 60 MeV where the extreme relativistic approximation is not expected to hold.

The activation of ^{63}Cu by the (e, n) or (γ, n) reactions was used as the detector in this study. The ten-minute half-life identifies the reaction, avoiding the multiparticle reactions which are sometimes confusing when outgoing particles are detected. Furthermore, the thin radiators of Cu, Mo, Ta, or Th, could be placed ahead of and in actual physical contact with the ^{63}Cu targets so that all of the bremsstrahlung would be detected.

THE EXPERIMENT

The $^{63}\text{Cu}(e, n)$ cross section was measured both at the Center for Radiation Research (CRR) of the National Bureau of Standards (from 20 to 60 MeV) and at the

Linear Accelerator Laboratory (LAL) of the University of São Paulo (from 13.5 to 30.5 MeV). At LAL the targets (99.9% enriched ^{63}Cu foils $10.0 \pm 0.2 \text{ mg/cm}^2$ thick, determined by weighing) were irradiated in a vacuum chamber located at the end of the beam pipe, while the irradiated charge was measured with a Faraday cup. The charge was recorded with a multichannel scalar in ten-second intervals.

A clock was started at the end of the irradiation, and the target, after removal from the irradiation chamber, was placed in a target holder fastened to a 27 cm^3 coaxial

FIG. 1. The electrodisintegration cross section $^{63}\text{Cu}(e, n)$ measured in the LAL (open circles) and in the CRR (closed circles). The line is the result of the best fit.

Ge(Li) detector. The target holder included a 3.10 ± 0.01 mm thick copper plate placed between the target and the detector to increase the β^+ annihilation rate. For irradiations with energies below the $(\gamma, 2n)$ threshold we made typically three 5 min counts of the target activity, and for energies above 20 MeV the target was counted eight to ten times, spanning a time interval of approximately three hours. The β^+ emission rate of the target was monitored by detecting the 511 keV annihilation gamma.

At the CRR the experiment was performed using the electroactivation facility, which includes a chamber at the end of the beam pipe in the electron scattering measurement room. The chamber can be vacuum isolated from the rest of the system and has a remote motor driven target holder, which is readily accessible for changing targets.

The electron beam energy was established by elastic electron scattering and its resolution set at one percent.³ During the runs the beam was monitored with a nonintercepting ferrite monitor which was in turn compared with the standard Faraday cup.⁴ Fluctuations in the beam current were recorded by scaling the measured charge in ten-second time intervals. The irradiation times were about 1.5 half-lives of ^{62}Cu , with beam currents ranging

from 3 to 10 μA . The beam spot size at the irradiation site was not larger than 0.7 cm in diameter, as observed on a view screen.

The targets consisted of 99% enriched ^{63}Cu foils approximately 5 mg/cm^2 thick. We used twelve targets. Eight of them had a thickness determined by weighing, the other four by normalizing the cross section at 22 MeV for the average value obtained with the known targets. A clock was started at the instant the irradiation was terminated. The target was then removed from the vacuum system, placed in the detector system, and counted for about four half-lives of ^{62}Cu . The elapsed time between the end of irradiation and the start of counting was around five minutes.

The detection system consisted of a copper annihilator source sandwich mounted a nominal 25 cm from a high purity Ge detector and a computer-based pulse-height analyzer. The annihilator sandwich consisted of two rectangular pieces of copper ($50 \times 38 \times 3.2$ mm), positioned above and below the target foil. The coaxial detector was 4.46 cm long and 3.96 cm in diameter, with a dead layer of 0.5 mm Be and 0.3 μm of inactive Ge, being mounted vertically. The detector had a 1.75 keV resolution and a relative efficiency of 10.9% for the 1332 keV gamma ray

FIG. 2. (a) Electro-plus-photodisintegration yield measured with a copper radiator. The line represents the best fit, obtained with the Seltzer spectrum (see the text). (b) Electro-plus-photodisintegration yield measured with a molybdenum radiator. The line represents the best fit, obtained with the Seltzer spectrum (see the text). (c) Electro-plus-photodisintegration yield measured with a tantalum radiator. The line represents the best fit, obtained with the Seltzer spectrum (see the text). (d) Electro-plus-disintegration yield measured with a thorium radiator. The line represents the best fit, obtained with the Seltzer spectrum (see the text).

in ^{60}Co . The annihilation emission rate from the target foils was determined by measuring the peak detection rate of the 511 keV annihilation radiation. The setup was calibrated for energy and annihilation efficiency. The energy and efficiency calibration was obtained by means of the long lived mixed radionuclide SRM 4275.⁵ The annihilation efficiency was estimated to 1.8% by measurement of ^{22}Na , $^{44}\text{Ti}/^{44}\text{Sc}$, and $^{82}\text{Sr}/^{82}\text{Rb}$ β^+ sources.

The data from the induced activity were acquired in five sequential four minute counting periods, synchronized with the end of the irradiation. After each counting period, the 511 keV annihilation peak counting rate was determined and corrected for decay to the end of the irradiation and for decay within the counting period. The annihilation source size correction was estimated to be less than 0.2% and ignored, as the source to detector distance was about 25 cm.

Besides the electrodisintegration cross section we also measured the electro-plus-photodisintegration yield, by placing a radiator immediately ahead of the target, to produce the bremsstrahlung. The radiators used were copper [0.485 g/cm², 3.62% radiation length (RL)], molybdenum [0.399 g/cm², 4.04% RL], tantalum [0.274 g/cm², 4.24% RL], and thorium [0.168 g/cm², 3.08% RL]. The results are shown in Figs. 1 (electrodisintegration cross section) and 2 (electro-plus-photodisintegration yields). The error bars shown represent statistical uncertainties only. There is a 4.2% uncertainty in the absolute scale due to the composite systematic uncertainties of the ^{62}Cu β^+ annihilation efficiency (1.2%), 511 keV detection efficiency (0.5%), distributed source correction (0.2%), ^{62}Cu β^+ branching ratio (4%),⁶ and Faraday cup calibration (0.1%).

DATA ANALYSIS

The number of ^{62}Cu nuclei produced and still present at the end of the irradiation is given by

$$N_0 = p\lambda^{-1}\sigma_{e,n} \left[\sum_i^m q_i \{1 - \exp[-\lambda(t_i - t_{i-1})]\} \times \exp[-\lambda(t_m - t_i)] \right], \quad (1)$$

$$\sigma_{e,n}(E_0) = \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,n}^{\lambda L}(E) N^{\lambda L}(E_0, E, Z, A) E^{-1} dE. \quad (3)$$

In Eq. (3) E_0 is the total electron energy and E is the excitation energy of multipolarity λL . In the same spirit, the yield obtained with a radiator placed in front of the target is

$$Y_{e,n}(E_0) = \sigma_{e,n}(E_0 - 2\Delta E_0) + N_r \int_0^{E_0-m} \sum_{\lambda L} \sigma_{\gamma,n}^{\lambda L}(E) K(E_0 - \Delta E_0, E, Z_r) E^{-1} dE, \quad (4)$$

where N_r is the number of atoms per square centimeter in the radiator of atomic number Z_r , $K(E_0, E, Z_r)$ is the bremsstrahlung cross section, and ΔE_0 is the electron energy loss in half the radiator thickness.

The experimental electrodisintegration cross section was corrected for the contribution due to the bremsstrahlung generated in the target. This was done by computing

where p is the number of ^{62}Cu nuclei per square centimeter, λ the ^{62}Cu decay constant, t_i the time at the end of the i th interval, q_i the number of electrons delivered during the i th interval, and t_m the total irradiation time. The factor in the large parentheses accounts for the fluctuation in the beam current and corrects for the decay of the ^{62}Cu formed during this irradiation. Since the time interval for recording the charge, $t_i - t_{i-1}$, was ten seconds and $\lambda = 0.0712 \text{ min}^{-1}$,⁷ the cross section is given by

$$\sigma_{e,n} = N_0 \lambda / (0.0118 p Q),$$

where

$$Q = \sum_i^m q_i \exp[-\lambda(t_m - t_i)].$$

The number of ^{62}Cu atoms present at the end of irradiation, N_0 , is related to the disintegration rate by

$$N_0 = (dN/dt|_0) / \lambda = (dC/dt|_0) / (\lambda \epsilon I),$$

where $(dC/dt|_0)$ is the counting rate at $t=0$, ϵ is the efficiency for detecting the annihilation radiation (geometric factor times the annihilation coefficient), and I is the branching ratio of the β^+ decay of ^{62}Cu (97.4%).⁶ Then

$$\sigma_{e,n} = (dC/dt|_0) / (0.0118 p Q I \epsilon). \quad (2)$$

For the data obtained at energies above the $(\gamma, 2n)$ threshold the counting rate was determined by a least squares fit of two exponentials to the decay curve, so that the contribution from the 194.7 min (Ref. 7) decay of ^{61}Cu [produced in the reaction $^{63}\text{Cu}(e, 2n)^{61}\text{Cu}$] could be separated. The LAL relative data were normalized to match the NBS absolute measurements. We used nine points in the energy covered by both experiments (20 to 30 MeV), making a linear interpolation between the LAL data when the energies did not coincide exactly. The average ratio between the NBS and the LAL data has a 3% dispersion, showing good compatibility between the two measurements.

The electrodisintegration cross section $\sigma_{e,n}(E_0)$ may be obtained from the photonuclear cross section $\sigma_{\gamma,n}^{\lambda L}(E)$, through an integration over the virtual photon spectrum $N^{\lambda L}(E_0, E, Z, A)$,

the second term on the right-hand side of Eq. (4), substituting half of the target thickness for N_r , and using for $\sum_{\lambda L} \sigma_{\gamma,n}^{\lambda L}(E)$ the data of Fultz *et al.*⁸ This contribution was then subtracted from the measured cross section. The correction amounted typically to 4% of the cross section in the LAL measurements and to 2% in the NBS data.

The cross section $\sigma_{e,n}(E_0)$ and the yields $Y_{e,n}(E_0)$ ob-

TABLE I. Integrated cross sections.

Bremsstrahlung spectrum	E1 fit		E1+E2 fit		
	Integrated cross section ^a E1 (MeV mb)	χ^2_{red} b	Integrated cross section ^a E1 (MeV mb)	E2 (MeV mb)	χ^2_{red} c
Seltzer	529±24	1.13	515±25	3.9±2.7	1.12
DBM	518±24	1.49	504±25	3.6±2.7	1.48
Schiff	502±23	1.76	443±24	16.5±2.6	1.13

^aUp to 30 MeV.^b60 degrees of freedom.^c59 degrees of freedom.

tained with the four different radiators have been fitted simultaneously. In Eq. (3) we used either $E1$ or $E1$ and $E2$ virtual photon spectra^{9,10} calculated in DWBA assuming a point nucleus and corrected for nuclear size effects by the method described in Ref. 2 (a very small correction). The results are compared in Table I. The accuracy of the $E1$ virtual photon spectrum calculated with nuclear size corrections has been tested experimentally for a ^{90}Zr target nucleus for electron energies between 17 and 105 MeV by Dodge *et al.*¹¹ This same experiment also assessed the compatibility of the $E1$ virtual photon spectrum and the Davies, Bethe, and Maximon bremsstrahlung spectrum for energies above 60 MeV. For the bremsstrahlung cross section in Eq. (4) we used the Schiff,^{12,13} the Davies, Bethe, and Maximon (DBM),^{13,14} and the Seltzer¹⁵ cross sections. The $E1$ photonuclear cross section was represented by a histogram, as shown in Fig. 3, and the $E2$ strength by a single bin 2 MeV wide located at the energy of the isoscalar $E2$ giant resonance.

THE BREMSSTRAHLUNG CROSS SECTIONS

Both the Schiff and the DBM spectra have already been discussed extensively in the literature.^{13,14} The Schiff formula, calculated in the Born approximation and with an approximate screening, has been shown to be about 10% higher than more accurate calculations,¹⁴ but is nevertheless commonly used in the analysis of photodisintegration

experiments, probably because it is a single and compact expression. We used the formula 3BS(e) of Ref. 13, substituting $Z(Z+1)$ for Z^2 to take into account the electron-electron bremsstrahlung.

For the DBM spectrum we used the expression 3CS of Ref. 13. This corresponds to the Bethe-Heitler extreme relativistic formula, with the Coulomb correction of Davies, Bethe, and Maximon¹⁶ and intermediate screening (with Thomas-Fermi screening functions). This expression was used for photon energies not exceeding 95% of the kinetic energy of the incoming electron. In the tip region (where the outgoing electron cannot be treated as extreme relativistic) we used the Schiff expression, making a suitable normalization to ensure a continuous spectrum. The electron-electron bremsstrahlung was taken into account in the same way as for the Schiff expression.

The Seltzer spectrum was obtained by a quadratic interpolation (for both electron and photon energies) from the data tabulated by Seltzer. For electron energies above 50% MeV, these data are obtained basically from the Bethe-Heitler unscreened expression (formula 3BN in Ref. 13) with added screening and Coulomb corrections; below 2 MeV it is based on the numerical phase shift results of Pratt *et al.*;¹⁷ and the region from 2 to 50 MeV is bridged by a least squares cubic spline interpolation.

In the high energy region Seltzer calculates the screening correction in the high energy approximation by the difference between the Bethe-Heitler screened and unscreened extreme relativistic cross sections [essentially formulas 3BS and 3BN(b) in Ref. 13], using screening functions calculated with Hartree-Fock atomic form factors. His Coulomb correction is given by a combination of factors. It agrees with the DBM calculation for high energy electrons which have not lost a substantial fraction of their original energy. Near the tip of the spectrum the correction changes sign and then behaves like a modified Elwert correction.¹⁸ The Elwert correction was modified to match the expected value of the cross section at the high frequency limit (when the incoming electron loses all its kinetic energy), based on calculations of Jabbur and Pratt¹⁹ and Pratt *et al.*¹⁷ Near the tip, the difference between the Seltzer and the Elwert corrections is small for nuclei with Z up to 60, becoming larger beyond that.

The electron-electron bremsstrahlung in the Seltzer formulation is calculated using the Haug cross section²⁰ plus a screening correction, which is calculated in the high energy approximation by the difference between the

FIG. 3. $^{63}\text{Cu}(\gamma, n)$ cross sections. Present work (histogram), Fultz *et al.* (Ref. 8) (continuous line), and Sund *et al.* (Ref. 21) (dotted line).

FIG. 4. Intensity spectra for bremsstrahlung photons produced by 25 MeV electrons incident on a 0.168 g/cm² Th ($Z=90$) radiator. Seltzer spectrum (continuous line), DBM spectrum (dashed line), and Schiff spectrum (dashed-dotted line).

Wheeler-Lamb screened and unscreened cross sections.

Figure 4 compares the intensity spectra for 25 MeV electrons incident on a 0.168 g/cm² Th ($Z=90$) radiator, calculated with the three different cross sections discussed above. Energy losses were not taken into account. The Seltzer spectrum is very close to the DBM for photon energies up to about 60% of the incoming electron kinetic energy, becoming larger for higher photon energies. The Schiff spectrum is about 10–20% larger than the other two far from the tip region, but turns out to be much closer to the Seltzer values at the tip than the DBM.

RESULTS AND DISCUSSION

The smooth curves in Figs. 1 and 2 show the result of the best fit, corresponding to the $E1$ photonuclear cross section shown as a histogram in Fig. 3. Table I summarizes the results obtained with the different bremsstrahlung cross sections with and without the inclusion of $E2$ strength and also shows the quality of fit, by the reduced chi-square (chi-square per degree of freedom), obtained in each case. The best fits were obtained with the Seltzer spectrum.

The fits using the Seltzer and the DBM cross sections did not require any $E2$ strength. The fit obtained using the Schiff spectrum, on the other hand, was greatly improved by the inclusion of the $E2$ strength but required an unreasonable amount, exhausting $148 \pm 24\%$ of an energy weighted $E2$ sum. We tried extending the $E1$ photonuclear cross section to energies above 30 MeV, but the cross sections obtained for the higher energy bins were compatible with zero, yielding approximately the same integrated cross section up to 30 MeV but increasing the overall uncertainty of the result.

The DBM and Seltzer spectra are the same (within 2%) for electron energies above 30 MeV and photon energies not close to the tip. This ensures that, for transitions in the giant resonance region (where most of the cross section is located), the yield calculated using either Seltzer or DBM will not differ significantly for electron energies

FIG. 5. Ratio of DBM/Seltzer yield of 17 MeV photons as a function of electron energy for copper ($Z=29$), molybdenum ($Z=42$), tantalum ($Z=73$), and thorium ($Z=90$) radiators.

above 30 MeV. Figure 5 shows the ratio of DBM/Seltzer yield of 17 MeV photons as a function of electron energy for different radiators. It is clear that above 28 MeV the number of 17 MeV photons calculated using either cross section is virtually the same. Below 28 MeV, on the other hand, the DBM cross section becomes smaller as the electron goes down (since the photon energy is kept constant that means getting close to the tip), due to the excessive Coulomb correction. The statistically unacceptable chi-square obtained with the DBM fit is due mainly to the photodisintegration data in the 20–30 MeV region that could not be properly fitted.

Figure 3 shows a comparison of our derived photonuclear cross section with two previous measurements.^{8,21} Our rough histogram gives little information about the detailed shape of the cross section, but the magnitude of the integrated cross section is in good agreement with that of Sund *et al.*,²¹ as shown in Table II. The uncertainty includes statistical and systematic errors.

TABLE II. Integrated cross sections.

Experiment	Ref.	Upper limit (MeV)	Integrated (γ, n) (MeV mb)
Sund ^a	21	25	490 ± 40
Fultz ^{a,b}	7	28	523 ± 52
This work		25	490 ± 25
		30	529 ± 33

^aQuasimonochromatic photons.

^bIncludes (γ, pn).

CONCLUSIONS

The cross sections for emission of one neutron by absorption of real and virtual radiation in ^{63}Cu have been measured. The $E1$ photonuclear cross section has been derived, yielding an integrated cross section of 529 ± 33 MeV mb up to 30 MeV. Our experiment was not sensitive to the presence of $E2$ strength and the derived $E1$ photonuclear cross section does not change significantly with the inclusion of an $E2$ strength of up to 35% of the $E2$ energy weighted sum rule.

Our results indicate that the DBM cross section should not be used for electron energies below 30 MeV or even above that if photons with energies above 60% of that of

the incoming electron account for a significant fraction of the transitions involved. The Schiff spectrum was found to be unacceptable almost in the whole energy region studied, being good only in a narrow band near the tip.

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" MÉTODO DE SELEÇÃO DE FÓTONS QUASE
MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

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"MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO"

Refere-se a presente invenção a um método de seleção de fótons quase monocromáticos e polarizados (obtenção de raios gama -
5 fótons de alta energia), através do controle dos parâmetros da radiação de bremsstrahlung por seleção temporal e angular e respectivo dispositivo.

Esse método pode ser utilizado em experimentos fotonucleares, microtomografia com raios gama e outras aplicações que envolvam a
10 detecção de raios gama e que requeiram o uso de fótons quase monocromáticos e polarizados para melhorar a qualidade da imagem. Atualmente são largamente utilizados em vários campos da ciência, da indústria e da medicina. A monocromatização e a polarização constituem os principais caminhos para se conseguir imagens com melhor definição
15 e maior riqueza de detalhes.

A radiação de bremsstrahlung, normalmente utilizada como fonte de raios gama é não cromática e não polarizada como um todo, o que faz com que diferentes componentes interajam distintamente com a matéria, prejudicando a definição das imagens.

20 Métodos de seleção de uma fração da radiação emitida fora do eixo de simetria do feixe permitem a obtenção de fótons polarizados com grau de polarização moderado.

O método conhecido como "marcação" ou "rotulação" de fótons, que consiste em medir a energia dos elétrons espalhados em

coincidência com os fótons (ou produtos da interação destes fótons),
permite a determinação da energia desses fótons (desprezando a
pequena fração da energia que é transferida ao núcleo), pois:

$$E_{\gamma} = E_{0e} - E_{fe} - E_N, \text{ onde } E_{\gamma}, E_{0e}, E_{fe} \text{ e } E_N \text{ são respectivamente as}$$

5 energias do fóton, inicial e final do elétron e transferida ao núcleo.

A combinação do método de marcação de fótons com a
colimação dos fótons e elétrons pós-bremsstrahlung em ângulos fora do
eixo de simetria do cone de bremsstrahlung, permite aumentar o grau de
polarização linear dos fótons. As dificuldades desse método estão
10 associadas à necessidade de se usar um complexo espectrômetro
magnético para a determinação das energias dos elétrons espalhados,
além de que a necessidade da colimação dos fótons e elétrons reduz a
intensidade do feixe final. Essas dificuldades fazem com que método de
marcação de fótons seja usado em experimentos em física fundamental,
15 não apresentando aplicações práticas.

O objetivo da invenção ora proposta, é obtenção de método e
respectivo dispositivo, que proporcionarão o aumento do grau de
polarização linear do feixe de fótons e o aumento da intensidade da
componente quase monocromática da radiação extraída.

20 O método desenvolvido representa uma evolução do método de
marcação de fótons usual, dada a utilização de uma característica
especial da seção de choque de bremsstrahlung quádrupla diferencial,
que apresenta um máximo localizado para determinadas combinações
das variáveis cinemáticas. Quando os ângulos de emissão do fóton e do

elétron e a energia do fóton correspondem a uma situação em que o momento transferido ao núcleo espalhador seja mínimo, a seção do choque e o grau de polarização linear apresentam máximos. Esta propriedade essencialmente diferencial (que corresponde a uma
5 situação de colimação acentuada nos ângulos de fótons e elétrons) foi detectada em um estudo detalhado do comportamento da seção do choque diferencial de bremsstrahlung, mas tende a desaparecer quando a seção de choque é integrada. Essa propriedade pode ser conservada, no entanto, se do fluxo total de fótons e elétrons espalhados forem
10 selecionados apenas aqueles que são correlacionados entre si e com as características desejadas de monocromatização e grau de polarização. Colimadores (ou detectores) de elétrons e fótons em geometrias específicas que selecionam a condição de momento transferido mínimo, permitem a obtenção de um espectro de coincidência de fótons quase
15 monocromáticos sem a necessidade de se determinar a energia final do elétron.

O método e dispositivo ora propostos, permitem a utilização de detectores planares de silício para a detecção dos elétrons e seleção dos ângulos de emissão dos fótons e elétrons, para a obtenção de
20 imagens com maior definição e melhor qualidade; além de permitir uma redução nas dimensões do dispositivo; tornando-o mais simples.

O método e respectivo dispositivo, são representados através das figuras:

Figura - 1: Vista esquemática do dispositivo.

Figura - 2: Representação gráfica, em função da energia do fóton das seguintes grandezas : valor médio do momento transferido ao quadrado; intensidade espectral; polarização.

5 Figura - 3: Representação gráfica da intensidade espectral em função da energia dos fótons.

Conforme a figura 1, o dispositivo consiste em feixe de elétrons (1), Conversor de fótons (2), mosaicos detectores (3 e 4), eletroimã (5), absorvedores (6 e 8) e objeto (7), o qual permite a realização prática do
10 método sugerido, a exemplo da microtomografia.

O feixe de elétrons (1) será incidido sobre o conversor (2); os fótons e elétrons secundários são detectados por dois mosaicos detectores planares de silício (3 e 4); elétrons que não interagem no conversor e a parte central do feixe de fótons são absorvidos nos
15 absorvedores (6 e 8).

Um dos mosaicos, em especial (3) é utilizado para detectar elétrons e é apenas dividido em setores, enquanto o outro mosaico (4) serve para detecção dos fótons após sua passagem pelo objeto de estudo e seus setores apresentam uma estrutura interna celular.

20 A aceitação no ângulo ϕ de cada setor é 0,3 rad., enquanto que a aceitação no ângulo θ deve ser determinada pelos requisitos de monocromatização (na maioria dos casos ela deve ser da ordem de $m_e/2 E_{0e}$). Os setores para detecção de elétrons e fótons têm as mesmas dimensões angulares.

O método consiste em: o feixe de elétrons primário interage com conversor (2), produzindo fótons de bremsstrahlung [γ] e elétrons espalhados [e']. Elétrons são desviados pelo eletroímã (5) de desvio para o absorvedor (6) do feixe.

O sistema de aquisição seleciona as coincidências entre setores cujas posições angulares correspondem ao mínimo de momento transferido para aquela energia de fóton. Isso significa que as seguintes condições são satisfeitas:

$$\frac{\theta^i}{\theta^j_\gamma} = F(E_0, E_\gamma) \text{ e } |\phi^i - \pi - \phi^j_\gamma| = 0,3 \text{ rad.}, \text{ onde } F(E_0, E_\gamma)$$

é a função calculada.

Os fótons selecionados do espectro total pela coincidência de tempo e pelas condições acima, são quase - monocromáticas e polarizados linearmente.

O método foi verificado, por simulação pelo método de Monte Carlo, para radiação de bremsstrahlung produzida por um feixe de elétrons com energia inicial $E_0 = 15 \text{ MeV}$. Na simulação foi utilizada a expressão de Bethe-Heitler para a seção de choque de bremsstrahlung quíntupla diferencial e as condições geométricas do equipamento proposto. A figura 2, mostra o momento transferido ao quadrado médio, a intensidade espectral e a polarização em função da energia do fóton, para um par de detectores com $\theta_0 = 0,017 \text{ rad.}$, $\phi_0 = 0,3 \text{ rad.}$, $\theta_\gamma = 0,034 \text{ rad}$

e $\phi_\gamma = \pi$ rad, que seleccionam fótons quase monocromáticos com energias em torno de 5 MeV. A figura 2 mostra que o pico de intensidade corresponde ao mínimo do momento transferido e que o grau de polarização linear no pico chega a cerca de 90 %, o que é
5 aproximadamente 50% acima do que é obtido pelos métodos convencionais e próximo do máximo possível (100 %). A figura 3 mostra três espectros, para três diferentes pares de setores dos detectores de elétrons e fótons, com posições que seleccionam fótons quase monocromáticos com diferentes energias. A intensidade da
10 componente quase monocromática é limitada pela taxa aceitável de coincidências acidentais entre os setores dos detectores de elétrons e fótons. O método ora proposto, permite elevar a taxa de produção em cerca de 10 vezes em relação aos métodos convencionais, porque permite a utilização de toda a aceitação angular dos fótons sem
15 prejudicar o grau de polarização obtido. A intensidade de fótons para um par ($E_\gamma=9,4$ MeV) é aproximadamente 10^5s^{-1} , considerando a resolução temporal em 100 ps, relação pico/fundo igual a 5 e nível superior de discriminação no canal de elétrons em 13 MeV.

O método e o respectivo dispositivo ora propostos, têm utilização
20 nos experimentos no campo da física fotonuclear e na obtenção de imagens precisas, quando de diagnósticos médicos através de imagens, de raios gama, a exemplo da microtomografia e outros.

REIVINDICAÇÕES

1. " MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

caracterizado pelas seguintes etapas e respectivo dispositivo.

- 5 a) Incidir um feixe de elétrons sobre o conversor.
- b) Interação do feixe de elétrons com o conversor.
- c) Detecção dos elétrons.
- d) Detecção dos fótons.
- e) Desvio dos elétrons pelo eletroimã.
- 10 f) Seleção dos fótons.
- g) Obtenção da Imagem de alta definição.

2. " MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

de acordo com a reivindicação 1, caracterizado por ser o dispositivo
15 composto por feixe de elétrons (1), conversor de fótons (2), mosaicos
detetores (3 e 4), eletroimã (5), absorvedores (6 e 8).

3. " MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

de acordo com a reivindicação 2, caracterizado por serem os detetores
20 (3) e (4) de estado sólido, semicondutores, planos, divididos em setores
e subdividido em células.

4. " MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

de acordo com as reivindicações 2 e 4, caracterizado por terem os detectores (3 e 4) setores de ângulos idênticos.

5. " MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

5 de acordo com a reivindicação 2, caracterizado por serem os detectores (3 e 4) posicionados coaxialmente à direção do feixe de elétrons incidente.

6. " MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

10 de acordo com a reivindicação 2, caracterizado por terem os sinais em coincidência temporal entre os setores dos detectores (3) e (4), cujas posições angulares correspondam ao mínimo do momento transferido.

FIGURA 1

FIGURA 2

FIGURA 3

RESUMO

" MÉTODO DE SELEÇÃO DE FÓTONS QUASE MONOCROMÁTICOS, POLARIZADOS E RESPECTIVO DISPOSITIVO "

Trata-se de um método de seleção de fótons quase
5 monocromáticos e polarizados (obtenção de raios - gama), através da
produção de fótons de bremsstrahlung pela irradiação de um alvo fino
com um feixe de elétrons e na seleção de uma componente
monocromática e polarizada desses fótons, pela detecção desses fótons
em coincidência com os elétrons espalhados por meio de detectores de
10 estado sólido, planos e organizados em mosaicos, com dimensões
angulares específicas e um sistema de aquisição de dados que
seleciona as coincidências entre setores específicos e respectivo
dispositivo, desenvolvido para o uso em experimentos de física
fotonuclear e na obtenção de imagens precisas, quando de diagnósticos
15 médicos através de imagens de raios gama, a exemplo da
microtomografia e outros.

RESUMO

" KIT PARA RECUPERAÇÃO DE MERCÚRIO EM AMÁLGAMAS ODONTOLÓGICOS "

Trata-se de um kit para recuperação de mercúrio em amálgamas
5 odontológicos, composto por câmara de aquecimento (1), tampa da
câmara (2), resistências (3), isolante (4), bulbo (5), tampa do bulbo (6),
conexão (7), haste de condensação (8), juntas (9) dilatações (10)
alimentação (11) e orifício (12); para uso em consultórios, laboratórios
de prótese e faculdades de odontologia, na recuperação de mercúrio em
10 amálgamas, tanto em pequena escala, no próprio ambiente de trabalho
do profissional, como em escala semi - industrial ou industrial.

A recuperação do mercúrio através do Kit, é superior a 97%.

RESUMO

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Disintegration of the Deuteron by Tagged, Linearly-Polarized Photons: Sensitivity of the Differential Cross Sections

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We present results of theoretical calculations together with simulations for some observables of the reaction $\gamma d \rightarrow np$ at energies close to the threshold. This approach allows the selection of the set of observables and the best kinematical conditions in order to probe different potentials. The simulations used the parameters of the forthcoming São Paulo Microtron accelerator and photon tagging facilities.

I. Introduction

The photodisintegration of the deuteron is a powerful tool for studying the nucleon-nucleon interaction, helping to probe the various NN potentials. The advantage of this experiment is that we combine the well-known electromagnetic interaction with the study of the simplest nuclear aggregate. The disintegration of the deuteron is usually a test case for new experimental techniques, also bringing up more information about the strong interaction.

The interpretation of low-energy breakup data presents problems when one tries to make use of realistic potentials, as those associated to (a) the choice of observables sensitive to the NN -force properties and their description; (b) the introduction of non-nucleonic degrees of freedom, like meson-exchange currents (MEC); and (c) the isobar configurations and relativistic corrections to the one- and two-body currents. Effects like final-state interaction (FSI) must also be

taken into account [1].

The near-threshold region is particularly noticeable, due to simplifications on both theoretical and experimental aspects, as (a) the presence of a dominant multipole (usually $E1$ or $M1$), with negligible contributions from the others at the entrance channel; each multipole takes few partial waves in the exit channel; and (b) the absence of background from other channels or competing reactions. All this simplifies the interpretation of low-energy data.

The use of polarized tagged photons is of particular interest, because the polarization observables contain interference terms between reaction amplitudes and are sensitive to small amplitudes belonging to contributions of non-nucleonic degrees of freedom. Although the latter are generally small, this is compensated by the high experimental accuracy attainable in this energy range. Besides, FSI effects are strong close to the threshold

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and, thus, convenient to be studied.

This work presents an analysis of the sensitivity of differential cross sections for different theoretical predictions in the photon low-energy region. Predictions of different calculations are analyzed with the expected experimental results obtained from Monte Carlo simulation, including an analysis of the experimental uncertainties expected for the experimental conditions of the São Paulo Microtron Accelerator and photon tagging facilities.

II. General definitions

The differential cross section of the reaction $\gamma d \rightarrow np$ (see Fig.1.a) can be expressed in terms of the helicity amplitudes [2] for γ -quanta, with any polarization, incident on a non-polarized target, and in the center of mass system (CM), as follows

$$\frac{d\sigma}{d\Omega} = \frac{1}{(8\pi W)^2} \frac{|\vec{p}|}{|\vec{K}|} Sp\{R\} \quad (1)$$

where

$$Sp\{R\} = \sum_{\{\lambda\}} T_{\lambda_p, \lambda_n}^{\lambda_\gamma, \lambda_d} \rho_{\lambda_\gamma, \lambda_\gamma'} T_{\lambda_p, \lambda_n}^{*\lambda_\gamma', \lambda_d}, \quad (2)$$

$T_{\lambda_p, \lambda_n}^{\lambda_\gamma, \lambda_d}(W, \cos\theta_p)$ is the CM helicity amplitude; λ_q , $q = \gamma, d, p, n$, are the helicities of the photon, deuteron, proton and neutron, respectively; $W = D_0 + K_0 = 2E_0$ is the energy of the pn pair, and θ_p is the angle between

the momenta of the photon and the outgoing proton (see Fig.1.b). In the case of a coincidence experiment, $\Delta\Omega$ is the solid angle subtended by the detector pair, according to the reaction kinematics; $\rho_{\lambda_\gamma, \lambda_\gamma'}$ is the polarization density matrix of the photon [3]. For the case of a completely linearly polarized photon, their elements are

$$\begin{aligned} \rho_{1,1} &= \rho_{-1,-1} = \frac{1}{2}, \\ \rho_{1,-1} &= -\frac{1}{2}e^{-2i\phi}, \\ \rho_{-1,1} &= \rho_{1,-1}^*, \end{aligned} \quad (3)$$

where ϕ is the angle between the polarization vector of the photon and the x axis (see Fig.1.b). Using (2) and (3), the cross section can be written, in the case of partial polarization, as:

Figure 1. General notation.

$$\frac{d\sigma}{d\Omega}(\theta, \phi, E_\gamma) = \frac{d\sigma_0}{d\Omega}(\theta, E_\gamma) + P_\gamma \frac{d\sigma_1}{d\Omega}(\theta, E_\gamma) \cos(2\phi), \quad (4)$$

where

$$\frac{d\sigma_0}{d\Omega}(\theta, E_\gamma) = \frac{1}{2} \left(\frac{d\sigma_{\parallel}}{d\Omega}(\theta, E_\gamma) + \frac{d\sigma_{\perp}}{d\Omega}(\theta, E_\gamma) \right). \quad (5)$$

Here we have used the following notation for the differential cross sections

$$\frac{d\sigma_{\{\parallel\}}}{d\Omega} = \frac{d\sigma}{d\Omega}(\phi = \left\{ \begin{array}{l} 0 \\ \frac{\pi}{2} \end{array} \right\}), \quad (6)$$

where the upper (lower) subscript on the left corresponds to the upper (lower) choice of ϕ on the right hand side. P_γ is the degree of linear polarization of the photon and $d\sigma_1/d\Omega$ characterizes the azimuthal asymmetry. We can define the azimuthal cross section asymmetry, $\Sigma(\theta, E_\gamma)$, as

$$\Sigma(\theta, E_\gamma) = \left(\frac{d\sigma_1/d\Omega}{d\sigma_0/d\Omega} \right) = \frac{1}{P_\gamma} \frac{(d\sigma_{\parallel}/d\Omega) - (d\sigma_{\perp}/d\Omega)}{(d\sigma_{\parallel}/d\Omega) + (d\sigma_{\perp}/d\Omega)}. \quad (7)$$

The differential cross section can then be written as

$$\frac{d\sigma}{d\Omega}(\theta, \phi, E_\gamma) = \frac{d\sigma_0}{d\Omega}(\theta, E_\gamma) [1 + P_\gamma \Sigma(\theta, E_\gamma) \cos(2\phi)]. \quad (8)$$

For complete linear polarization, $P_\gamma = 1$, and with a convenient choice of directions, the differential cross sections can be written as

$$\frac{d\sigma_{\parallel, \perp}}{d\Omega}(\theta, E_\gamma) = \frac{d\sigma_0}{d\Omega}(\theta, E_\gamma) [1 \pm \Sigma(\theta, E_\gamma)], \quad (9)$$

where the sign $+$ ($-$) on the right corresponds to $\parallel$ ($\perp$).

The statistical uncertainties of the differential cross sections are

$$\Delta \left(\frac{d\sigma_{\parallel, \perp}}{d\Omega} \right) = \left\{ \left[\Delta \left(\frac{d\sigma_0}{d\Omega} \right) \right]^2 (1 \pm \Sigma)^2 + \left(\frac{d\sigma_0}{d\Omega} \right)^2 (\Delta \Sigma)^2 \right\}^{\frac{1}{2}}, \quad (10)$$

where $\Delta(d\sigma_0/d\Omega)$ and $\Delta \Sigma$ are the statistical uncertainties of the $d\sigma_0/d\Omega$ cross section and of the Σ -asymmetry, respectively. The uncertainty of the Σ -asymmetry is

$$\Delta \Sigma = \frac{1}{2P_\gamma N_0} \left\{ (\Delta N_{\parallel})^2 + (\Delta N_{\perp})^2 + (N_{\parallel} - N_{\perp})^2 \left[\left(\frac{\Delta P_\gamma}{P_\gamma} \right)^2 + \left(\frac{\Delta N_0}{N_0} \right)^2 \right] \right\}^{\frac{1}{2}}, \quad (11)$$

where $\Delta P_\gamma/P_\gamma$ is the relative uncertainty of the degree of linear polarization and $N_{0, \parallel, \perp}$ and $\Delta N_{0, \parallel, \perp}$ are the measured experimental yields and their uncertainties for each polarization situation, with

$$\Delta N_{0, \parallel, \perp} = [(1 + R_{0, \parallel, \perp}) N_{0, \parallel, \perp}]^{\frac{1}{2}}, \quad (12)$$

where $R_{0, \parallel, \perp}$ is the ratio of the number of random to true coincidences, for each polarization situation.

The number of counts for one tagging energy interval (energy bin) can be evaluated by the following expression:

$$N_{0, \parallel, \perp} = \frac{d\sigma_0}{d\Omega} N_n N_\gamma^{tag} t (4\pi\epsilon) \left\{ \begin{array}{c} 1 \\ 1 + P_\gamma \Sigma \\ 1 - P_\gamma \Sigma \end{array} \right\} \quad (13)$$

where N_n is the number of nuclei per cm^2 in the target, N_γ^{tag} is the number of photons per second per tagger channel, t is the acquisition time and ϵ is the absolute efficiency of the detector pair. The parameter $4\pi\epsilon$ includes the geometrical solid angle defined by the detectors at the target position and the intrinsic efficiency of the detectors. The absolute efficiency ϵ , determined by Monte Carlo simulation [4], corresponds to a weighted average of the intrinsic efficiency over the solid angle defined by the detector's surface. Since this is a coincidence experiment, the solid angle is actually a dynamic one, because its evaluation includes the kinematics of the reaction [4].

III. Theoretical Model

The theoretical analysis was carried out in the framework of a relativistic, gauge-invariant theory of the photo- and electrodisintegration of light nuclei. This approach was developed during the last few years [5, 6, 7, 8] and is based on quantum field theory with methods which allow to introduce nuclear structure, final-state interaction and MEC effects, although these last are present non-explicitly, via gauge-invariance constraints.

Fig. 2 shows two sets of diagrams used in the calculation. Set number I corresponds to the relativistic pole approximation. The last diagram takes MEC into account. The contact current does not have pole-type peculiarities, but must be taken into account in the pole approximation, because the contact amplitude is the one that provides nuclear current conservation for this set of diagrams.

Set number II is characterized by the nuclear vertex function $d \rightarrow pn$ and by the half off-shell T matrix of $pn \rightarrow pn$ scattering. Diagrams (5) and (6) correspond to the usual final-state interaction; loop diagrams (7) and (8) are necessary for current conservation, and account for both final-state interactions and meson-exchange currents.

FSI was taken into account through the 1S_0 resonance, its contribution to the amplitudes of the pn rescattering being calculated via a multipole expansion as

$$f_l = A_l \frac{e^{2i\delta_l} - 1}{2i}, \quad (14)$$

where

$$\delta_l = \delta_l^{\text{background}} + \delta_l^R, \quad (15)$$

with

$$\delta_l^R = \arctan \left[\frac{G_R}{2(M_R c^2 - W)} \right]. \quad (16)$$

The parameters of the resonance are $G_R c^2 = 300 \text{ keV}$, $M_R c^2 = m_p + m_n + I_R$, and $I_R c^2 = 67 \text{ keV}$. The background phase shift was calculated as in ref. [9]. The resonance forces, A_l , were calculated by the one-loop diagrams of Fig. 2 [5].

Figure 2. Two gauge invariant sets of diagrams used for the description of the deuteron photobreakup.

We have used two types of realistic deuteron wave functions: the first given by the Paris potential [10] (called DWF₁), and the other is the relativistic Gross' wave function [11] (DWF₂). They adequately describe not only the ground state properties, but also the elastic and quasielastic form factors.

IV. Results and discussion

One of the objectives of this paper is to find out the conditions where discrepancies, arising from different theoretical predictions, could be identified by experiment. In this section we first describe the experimental conditions that will be attained at the Linear Accelerator Laboratory, with the use of the São Paulo Microtron and the photon tagging facilities. Next, we present the theoretical calculations and Monte Carlo simulations, and describe the conditions where the discrepancies can be experimentally checked, as well.

A. Experimental conditions

Table 1 presents some of the important parameters entering the choice of the experimental conditions. The values are those expected for the forthcoming São Paulo Microtron beam and photon tagging facilities [12, 13].

Table 1 - Characteristics of the experiment.

E_γ (MeV) (mean tagging energy)	3.0	5.0
N_γ^{tag} (s^{-1})	1.4×10^8	
P_γ	0.7	
$\Delta E_\gamma^{\text{tag}}$ /channel (MeV)	0.2	0.3
N_γ^{tag} /channel (s^{-1})	2×10^7	
N_n (cm^{-2})	9×10^{19}	
ϵ	0.013	
ΔE_p (keV)	30	7

The Bremsstrahlung beam intensity and polarization were obtained by Monte Carlo simulation [14] for an incident electron energy of 15 MeV. Photons and post-Bremsstrahlung electrons are collimated off axis. The geometry of the collimators (dimensions and positioning) was chosen in order to optimize the degree of polarization and to concentrate the photon spectral intensity in the near-threshold energy region. In this case, random coincidences, which restrict the photon beam intensity [13], arise mainly from the partial angular acceptance of the electron collimator. Signals of

reaction products from untagged photons (produced by electrons that hit the collimator) will cause random coincidences with signals from the focal plane detectors. The value of N_7^{tag} , shown in table 1, assumes a ratio of the number of true to random coincidences $R = 1$, summed over all (seven, in this case) tagger channels. The target thickness, N_n , corresponds to a low pressure (0.5 atm), 7-cm long (along the beam direction), gaseous deuterium target. The reaction products detection system is composed of two detector planes placed at opposite sides of the target. The protons are detected by silicon strip detectors, described in ref. [15], while the neutrons are detected by plastic scintillators coupled to photomultipliers (0.025 MeV threshold), and this will be the subject of a forthcoming publication. The absolute detection efficiency, ϵ , is typical for opposing detector pairs seeing the target at 90° . Proton energy losses in the target (ΔE_p , table 1) were calculated for 90° emission and averaged over the photon beam dimension ($\sim 1 \text{ cm}^2$).

B. Model predictions and simulation results

In this section we present the results obtained for the four different calculations performed, namely using DWF₁ or DWF₂, with or without the inclusion of final state interactions (FSI).

Figure 3. Differential cross sections $d\sigma_{||}/d\Omega(E_\gamma, \theta)$, calculated with DWF for Paris potential and including FSI.

The cross sections $d\sigma_{||}/d\Omega$ and $d\sigma_0/d\Omega$ (shown in Figs. 3 and 4 for DWF₁ with FSI) have a similar angu-

lar behavior: a maximum at $\theta = 90^\circ$, which is characteristic of the E1 multipole. Calculations using other models (DWF₁ without FSI or DWF₂ with or without FSI) give results that are visually identical to those shown in Figs. 3 and 4.

Figure 4. Differential cross sections $d\sigma_0/d\Omega(E_\gamma, \theta)$, calculated with DWF for Paris potential and including FSI.

The Σ -asymmetry is practically not sensitive to different choices of DWF, giving very close results for DWF₁ or DWF₂. It also presents very similar angular behaviors for calculations including or not FSI, as shown respectively in Figs. 5 and 6. The absolute value of the Σ -asymmetry drops for angles close to 0 and π , especially at low photon energies and in the calculations including FSI.

Figure 5. Σ -asymmetry, calculated with DWF for Paris potential and including FSI.

Figure 6. Σ -asymmetry, calculated with DWF for Paris potential and without FSI.

Figure 7. Differential cross sections $d\sigma_{\perp}/d\Omega(E_{\gamma}, \theta)$, calculated with DWF for Paris potential and without FSI.

Figure 8. Differential cross sections $d\sigma_{\perp}/d\Omega(E_{\gamma}, \theta)$, calculated with DWF for Paris potential and including FSI.

As the Σ -asymmetry, the $d\sigma_{\perp}/d\Omega$ cross section is not sensitive to the choice of DWF, but is strongly sensitive to the inclusion of FSI. Figs. 7 and 8 show the

angular and energy dependence of $d\sigma_{\perp}/d\Omega$ calculated using DWF₁, without and with FSI, respectively. It is interesting to note that the inclusion or not of FSI gives different results not only for the absolute value of the cross section, but it is also responsible for radical changes in the angular and energy dependence.

Figure 9. Differential cross sections $d\sigma_0/d\Omega(E_{\gamma})$ for $\theta = 90^\circ$. The experimental uncertainties are shown as bands. The upper band corresponds to calculations with DWF for Paris potential (with and without the inclusion of FSI, visually the same), the lower band uses Gross potential (with and without the inclusion of FSI, visually the same). (a) corresponds to $L = 10$, and (b) to $L = 100$ ($L = 4\pi\epsilon N_n N_p^{\text{tag}} \times 10^{-30} \text{ sr}/\mu\text{b}$, $L = 100$ corresponds to the data from table 1 and 100 h). The energy resolution for both cases is 0.2 MeV at the beginning of the range and 0.3 at the end.

Figs. 9–12 respectively show the energy dependence of $d\sigma_0/d\Omega$, $d\sigma_{\parallel}/d\Omega$, Σ -asymmetry and $d\sigma_{\perp}/d\Omega$. This form was chosen because the most convenient way of resolving the discrepancies seems to be the detailed measurement of the energy dependence of the cross sections at energies close to the threshold, at $\theta = 90^\circ$. The ability to resolve different model calculations depends strongly on the uncertainties that can be achieved in the experiments, which, in turn, are determined by the experimental setup. The main sources of uncertainties

in the proposed experiment are the count rate statistics (due to expected low count rates) and the achievable accuracy in the determination of the degree of polarization of the photon beam. The bands (instead of curves) presented in Figs. 9–12 represent the experimental uncertainties, evaluated for the expected conditions of the São Paulo Microtron and tagger, for a typical tagger channel. The evaluation assumed the data presented in table 1 and beam times of 10, 100 or 200 h (see figure captions). The experimental uncertainties shown are only statistical. Possible systematic errors arising from the absolute values of the cross sections or the degree of polarization, were not considered. The energy resolution is about 0.2 MeV in the beginning of the range and about 0.3 MeV at the end. It corresponds to the width of the tagger channel.

Figure 10. Differential cross sections $d\sigma_{||}/d\Omega(E_\gamma)$ for $\theta = 90^\circ$. All designations are the same as in Fig. 9. The uncertainties of the experimental data are only statistical.

As it can be seen from Figs. 9 and 10, $d\sigma_{0,||}/d\Omega$, in the energy range 2.5–6.0 MeV, is sensitive to the DWF choice (the bands are separated) but not to FSI effects. On the other hand, Σ and $d\sigma_{\perp}/d\Omega$ (shown in Figs. 11

and 12) are not sensitive to the choice of DWF, but pinpoint FSI effects. From Figs. 9–12 one can also select the energy ranges where the bands do not overlap, which should be favoured in the measurements.

Figure 11. Σ -asymmetry for $\theta = 90^\circ$ versus photon energy, calculated with DWF for Paris and Gross potentials (not distinguished) with FSI (lower band) and without FSI (upper band). Bands in (a) correspond to $L = 10$ and in (b) to $L = 200$ (see caption of Fig. 9). The uncertainties of the experimental data are only statistical.

Figs. 10–12 also present some preliminary data, obtained with photons from channeling and similar equipment (gaseous target and silicon detectors). The experimental details are presented in ref. [16]. It can be seen from Fig. 10 that the experimental uncertainties (about 1%) would be enough to distinguish between Paris and Gross potentials if a better energy resolution is available. On the other hand, the data for Σ and $d\sigma_{\perp}/d\Omega$ (Figs. 11 and 12), even though favouring the calculation including FSI, could certainly be improved with better statistics and energy resolution. Moreover, since the photons were obtained by the channeling process, the degree of the polarization of the beam is subject to uncertainties and, as shown below, this strongly affects the confidence on the data.

Equations (7) and (9) define how sensitive is the analysis to systematic errors in the cross sections and degree of polarization of the photon beam. A 1% change in $d\sigma_0/d\Omega$ or P_γ shifts the bands of $d\sigma_{\parallel}/d\Omega$ of 1% or 0.5% respectively, in the range where $\Sigma \approx 0.98$ (see Fig. 11), which are much smaller than the 5% difference between the curves. The situation with $d\sigma_{\perp}/d\Omega$ is quite different. The same 1% change in $d\sigma_0/d\Omega$ or P_γ shifts the bands of $d\sigma_{\perp}/d\Omega$ of 1% and 100% respectively, which should be compared with the $\sim 100\%$ difference between both theoretical predictions. So, in order to be able to distinguish between the different calculations, the systematic errors in $d\sigma_0/d\Omega$ and P_γ must be within 5% and 1%, respectively.

Figure 12. Differential cross sections $d\sigma_{\perp}/d\Omega(E_\gamma)$ for $\theta = 90^\circ$ versus photon energy, calculated with DWF for Paris and Gross potentials (not distinguished) with FSI (upper band, below 3 MeV) and without FSI (lower band, below 3 MeV). Bands in (a) correspond to $L = 10$ and in (b) to $L = 200$ (see caption of Fig. 9). The uncertainties of the experimental data are only statistical.

V. Conclusions

This work shows that, in the near-threshold energy region, where discrepancies between different models

do exist, there are no experimental data with enough accuracy to resolve those discrepancies. Angular distributions of differential cross sections, which contain information about the multipolarity of the transitions, are also completely absent.

The considerations presented previously show that, within the experimental conditions that will be available at the São Paulo Microtron and tagging facilities, it will be possible to assess the assumptions of the relativistic gauge-invariant model of nuclear photodisintegration and, in particular, to differentiate between the deuteron wave functions based on Paris or Gross potentials by measuring $d\sigma_0/d\Omega$ or $d\sigma_{\parallel}/d\Omega$; also, the inclusion or not of final state interaction effects can be resolved by the measurement of $d\sigma_{\perp}/d\Omega$ or the Σ -asymmetry.

The analysis presented here can be considered as a study of (i) the necessary conditions to be achieved by the experimental equipment and acquisition system to be used in the measurement and (ii) the necessary accuracy to be achieved in the determination of the absolute value of the cross section $d\sigma_0/d\Omega$ (better than 5%) and the degree of polarization of the photon beam (better than 1%).

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NUCLEAR INSTRUMENTS & METHODS IN PHYSICS RESEARCH

Section A

Nuclear Instruments and Methods in Physics Research A 376 (1996) 455–461

Silicon strip detectors for fission

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Silicon strip detectors for fission

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Abstract

The performance of a silicon strip detector especially designed for photofission experiments and a particular arrangement for low count rate photofission measurements are presented. Other applications for the same detection system are discussed.

1. Introduction

The study of photofission of actinide nuclei, especially at the region close to the barrier, has provided much interesting information about the fission process. In particular, by investigating the fission fragment angular distributions it was possible to obtain important information about the multi-humped structure of the fission barrier [1]. Moreover, the study of the fission fragment mass, energy and charge distributions revealed important aspects of the dynamical properties of the fissioning system [2]. By measuring both the angular and mass distributions of the fragments it is possible to obtain information about the correlation between the multipolarity of the excitation and the mass division. First experiments in this direction have been carried out by Wilke et al. [3] and Steiper et al. [4] using position sensitive parallel plate avalanche detectors (PPADs).

It is clear that the kind of investigation mentioned above requires accurate measurement of the fission fragment angular, energy and mass distributions together with a precise knowledge of the bombarding photon energy. Current high duty-factor electron accelerators, associated with the tagged photon technique [5,6] and high quality fission fragment detectors constitute a powerful tool for

this kind of study. However, the tagged photon technique limits the photon flux to a maximum of about 10^7 s^{-1} [7,8], which appears to be incompatible with measurements involving low cross sections as the ones mentioned above. Moreover, in order to achieve this flux, the photon beam has to be non-collimated. This circumstance leads, especially in the case of low initial electron energy (thus large divergence), to large beam spots on the target.

To take into account the specific needs of this kind of measurement, the detection system should match the following requirements:

- high energy resolution for heavy charged particles;
- good angular resolution;
- high time resolution (comparable to the difference in the time of flight of the particles);
- stability over large periods of time;
- possibility of large illuminated target area;
- detector array as close as possible to the target.

All these requirements, in particular the high energy resolution and good stability over time, could be fulfilled by e.g. position sensitive silicon detector arrangements, which were widely used in the past in high energy physics for precise determination of charged particle coordinates [9–13]. Other possible choices like PPADs [14] and multiwire proportional parallel plate detectors (M3Ps) [15] are not adequate to our particular conditions. Indeed, experiments with low count rates, large illuminated target area, and limited available room for detectors are incompatible with detectors such as PPADs and M3Ps because of the long time-of-flight distances needed for achieving accurate mass determination.

This work presents the development of a detecting system suitable for low energy experiments involving low count rate statistics. This system consists of a new silicon strip detector arranged in an optimized geometry especially chosen to fulfill our purposes. In Section 2 we describe the

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Table 1
Main features of the silicon strip detectors

Description	Value
Sensitive area	$52 \times 52 \text{ mm}^2$
Thickness	$300 \mu\text{m}$
Number of Si crystals	2
Number of strips	16
Strip width	$2900 \mu\text{m}$
Interstrip distance	$100 \mu\text{m}$
Working voltage	10–100 V
Breakdown voltage	>100 V
Thickness of upper Al contact	$0.3 \mu\text{m}$
Thickness of lower Al contact	$0.3 \mu\text{m}$
depth of p^+ layer	$0.2 \mu\text{m}$

detectors characteristics and the tests carried out. In Section 3 the proposed detection system is discussed and in Section 4 we summarize the main conclusions achieved in this work.

2. The silicon strip detector

2.1. Manufacturing details

In order to accomplish all the requirements of the intended experiments, silicon strip detectors were especially manufactured at Kiev SPA Detector (Ukraine). These detectors use an n-Si substratum with resistivity of 2–5 k Ω cm. The p^+ -n junctions were made with implanted boron using planar technology. The beam energy was 15 keV and the implanted dose 5×10^{14} ions/cm 2 . Such low ion energies make small depth junctions ($0.2 \mu\text{m}$) and allow the use of low temperatures for the activation. These characteristics reduce the thickness of the input windows. The n $^+$ contact on the back side of the detector was made with implanted phosphorus. The phosphorus dose was 5×10^{15} ions/cm 2 at 30 keV. The p^+ -n junctions and the detector back side were covered by a thin aluminum layer. The two silicon substrata with their output connectors were

mounted on a plastic board. The strips were welded to the board connectors using gold strips. In Table 1 the main features of the silicon strip detectors are summed up.

2.2. Performance

Tests were carried out in order to assess the detector performance in measuring conditions. Initially, we checked the leakage current and depleted depth variations of the detectors as a function of the bias voltage. Table 2 summarizes the results of this study for four detectors. It is seen in this table that for a bias voltage of 20 V the depletion depth is more than $100 \mu\text{m}$, which is enough to detect fission fragments. Other important parameters related to the detector performance are discussed in the following subsections.

2.2.1. Energy resolution

The fission fragment spectra of actinide nuclei are known to exhibit a double-humped structure, which is the signature of asymmetric fission. For these nuclei the contribution of symmetric fission is small (for low excitation energies) and corresponds to the valley of the double-humped structure. Therefore, the clear separation between the two humps allows the determination of the mass distribution in terms of light and heavy fragment groups. Thus, the energy resolution and spectrum distortion problems for the detection of fission fragments are very important issues. This means that the detector performance must be designed to minimize the pulse height defect caused by the high energy loss along the ion track. Moreover, the energy resolution of detector and electronics should give a small contribution to the final energy spectrum.

The measurement of ^{252}Cf fission fragments is a standard test to assess the performance of heavy ion spectroscopy detectors. For that purpose, a ^{252}Cf source was placed inside a small reaction chamber together with the silicon strip detectors. The system was evacuated and kept at a pressure of about 10^{-3} Pa. Conventional electronics consisting of preamplifier, amplifier, and ADC, connected

Table 2
Relevant strip parameters obtained for four different strip detectors. Each detector consists of two silicon crystals with 8 strips each. DN: detector number; CN: crystal number

DN	CN	Current [nA]				Depleted depth [μm]				
		Bias (V) $\Rightarrow$	20	40	60	100	20	40	60	100
1	1		70–100	100–150	100–200	200–300	110	140	165	200
1	2		50–60	70–80	100–150	200	105	140	165	210
2	1		60–80	70–100	100–200	200–300	105	130	160	200
2	2		80–100	100–120	200–300	400–500	105	140	160	210
3	1		80–200	100–250	180–300	300–400	110	145	175	220
3	2		50–130	70–200	160–250	300–500	110	140	165	210
4	1		100–120	130–240	200–400	500–600	110	140	170	210
4	2		10–20	20–30	100	150–200	115	150	175	220

to a personal computer, were used in these tests. In Fig. 1 we present the ^{252}Cf energy spectrum obtained using a single strip of one particular detector. As we can see in this figure, this spectrum exhibits two prominent peaks (due to asymmetric fission) with an intermediate valley (due to symmetric fission). The two peaks correspond to the light (higher channels) and heavy (lower channels) fission fragments. Following the procedure of Weissenberger and collaborators [16], we fitted two Gaussians to the total spectrum and obtained the spectrum shape parameters in order to check the quality of the detector. For instance, the results obtained for the peak to valley ratio for the heavier and lighter fission fragments were 2.1 and 2.8 respectively, while the peak (lighter fragments) to peak (heavier fragments) ratio was about 1.3. The overall agreement of these parameters with the values tabulated by Schmitt and Pleasonton [17] and recently confirmed by Weissenberger et al. [16] are very good, indicating the good quality of the detector.

2.2.2. Time resolution

In order to assess the timing features of the detectors, we used a light emitting diode (LED) driven by a fast pulser with adjustable amplitude [8]. With this arrangement we could simulate the passage of alpha particles or fission fragments through the strips by adjusting proper amplitudes in the pulse generator. Since all the strips were

illuminated simultaneously, we used the generated signals to simulate coincidence spectra from two strips in different combinations. Fig. 2 presents two coincidence spectra added to appear in a single plot. We used different external delays in each spectrum to set the peaks 4 ns apart. The FWHM of each peak is about 0.4 ns. This width is mainly due to two factors: i) time variation of the emitted light; ii) time resolution of the strip detector [8]. Therefore, the real time resolution of the detector should be less than 0.4 ns.

2.2.3. Strip uniformity

The uniformity of the strips was checked through irradiation of the detector by an ^{241}Am alpha source. The pulse height spectrum was independently measured for each of the 16 strips of a particular detector (bias voltage: 20 V). These spectra were normalized in order to compensate for the change in solid angle from the center to the border strips (up to 12% for extreme strips). Finally these spectra were put together in one single spectrum as depicted in Fig. 3. The strips present different leakage currents and the ones having higher currents appear

Fig. 1. Pulse height spectra of ^{252}Cf fission fragments detected by a single strip (no. 11) of one particular detector. (A) without encoding system; (B) with encoding system. See text for further information.

Fig. 2. Time coincidence spectra in three specific cases: LED light flash (A) simulating α particles without encoding system; (B) simulating α particles with the encoding system; and (C) simulating fission fragments with the encoding system. See text for further information.

Fig. 3. Pulse height spectra of ^{241}Am alpha source measured by different strips. In order to present the spectra in the same scale, the gain was changed for each strip. The areas under the peaks are the same within $\pm 1.7\%$.

broader than the others, but all the peaks have the same area within 1.7%.

The uniformity along each strip was measured with an ^{241}Am alpha source and a 2 mm movable collimator. Fig. 4 shows the pulse height spectra obtained at four different positions along one single strip. The areas under the peaks are compatible within 5.0%, which shows that the strip response as a function of the incident position of the particle is quite uniform.

3. Detector arrangement

As mentioned before, the silicon strip detectors arrangement will be used in experiments involving i) very low cross sections; and ii) large illuminated target area needing, nevertheless, good angular resolution. The detector arrangement described here was designed for the specific case in which two reaction products are emitted in opposite directions (requiring a thin, self-supported target). For this reason, we opted for a *sandwich* configuration where the

Fig. 4. Pulse height spectra of ^{241}Am alpha source measured in different positions along one of the strips of a particular detector. See text for further information.

Fig. 5. Schematic drawing depicting the target position and the detector arrays. f_1 and f_2 correspond to fission fragments emitted in opposite directions.

target is tilted with respect to the photon beam and lies between two planes of detectors (see Fig. 5). The associated electronics select any pair of opposing strips, thus determining the fission fragment emitted angles. These features will be discussed in the following sub-sections.

3.1. Geometry optimization

The geometry of the detection system was optimized by a Monte Carlo simulation [18]. The photon coordinates on the beam spot and the emission angles θ_i and φ_i of one of the fission fragments were randomly generated. The angles of the second fission fragment were then calculated according to the reaction kinematics. Finally, the case where both particles were detected in coincidence by an opposing pair of strips was analyzed. The dependence of the coincidence rate on the beam diameter, on the distance from target to detector planes, and on the target angle with respect to the photon beam direction was also calculated [18]. For a beam diameter of 3 cm, the best target angle with respect to the photon beam direction is about 17° . A schematic drawing of the detector array including the target is shown in Fig. 5. The arrangement consists of two parallel planes of detectors, each of them with two strip detectors (thus having a total of 32 strips, see Table 1). The strips are oriented in a direction perpendicular to the photon beam direction and the distance between the detector planes is 40 mm. This setup covers polar angles $22^\circ < \theta < 158^\circ$. The dependence of the polar angle θ on the azimuthal angle φ (which varies along the strip) was checked for limiting cases (i.e., fragments emitted from the central and border regions of the target at 90° with respect to the photon beam; and fragments hitting far opposing strips). The results show that for $\theta = 90^\circ$ the polar angular resolution is $\pm 4^\circ$ while for $\theta = 22^\circ$ and 158° the polar angular resolution is $\pm 5^\circ$. In these limiting cases the φ acceptance varies from $\pm 50^\circ$ to $\pm 30^\circ$ respectively.

3.2. Electronics

The forthcoming experiments will require knowledge of parameters like the energy of both fission products and the polar angle θ , the latter being related to the identification of which pair of strips fired. Therefore, the electronics setup should be conceived to fulfil such requirements.

In order to reduce the number of electronic modules like preamplifiers (PA), analog to digital converters (ADC), time to digital converters (TDC), discriminators (D) and others, we opted for the method of time codification of the strip number based on delay lines. Fig. 6 shows a detailed scheme for one of the detectors.

The common backside output of the detectors provides a signal that is sent to a preamplifier (PA). This PA generates fast and charge sensitive output signals. The charge sensitive signal goes to an ADC (E signal) for energy determination. For the strip identification, two fast signals are used: one of them comes from the PA (T signal) and is sent straight to a discriminator (D) and then to a TDC module (start input). The other one comes from the particular strip which fired. This signal is delayed and, after the PA module, is sent to the TDC (stop input). Actually, all strip outputs are connected to each other by delay lines of 5 ns each, which means that the total delay of a particular strip can be only assigned to that very strip, thus providing its identification. We will refer to this system as the *encoding system*. Note that a further delay line (50 ns) is added in order to allow the identification of the crystal to which that strip belonged (each detector has two crystals with 8 strips each).

Pulses from two detectors observing the same target side are sent to OR modules. Pulses coming from the OR modules are sent to a coincidence/anti-coincidence module. In case two pulses arrive within the time resolution in this module, an AND pulse is generated and sent to the Master Trigger (MT), which is responsible for activating the data acquisition system. Gates are then generated for the ADCs and the bit pattern module (BP). The BP module is used for the identification of the detectors that fired. A pulse is sent to the crate controller (CC) and then the ADCs, TDCs and the BP are read out and transferred to the computer and saved in tape for off-line analysis. The events are saved in an *event by event* mode. In case only one pulse arrives in the coincidence/anti-coincidence module within the time resolution, and AND pulse is generated, which serves as a veto to the TDC. Scalers are used to keep track of the total number of counts.

Finally, it is important to mention that this electronic setup could be easily coupled to a larger setup including the photon tagging electronics.

3.2.1. Performance

Since the encoding system is placed right after the detectors, it is necessary to use high capacitance preamplifiers in order to keep electronic noise at acceptable levels to achieve good time and energy resolutions. The charge-sensitive and high-capacitance preamplifiers (PA) used in this work were manufactured at Kiev Institute of Nuclear Physics. With this PA it is possible to have precise determination of the energy loss in the detector with an open-loop gain of $K = 9000$. The full energy resolution

Fig. 6. The electronics setup. PA: preamplifier; D: discriminator; ADC: analog to digital converter; TDC: time to digital converter; CC: crate controller; BP: bit pattern; MT: master trigger. The dashed box depicts the electronics for one detector and is the same for all four detectors used.

Fig. 7. Equivalent energy of the electronic noise at the fast output channel of the preamplifier as a function of the external capacitance. The dotted lines connecting the points were drawn to guide the eye.

achieved in this case is $R = 2.06 \text{ keV} + 0.018 \text{ keV/pF}$ with an output pulse rise time of $t = 25 \text{ ns} + 0.035 \text{ ns/pF}$ at a feedback capacitance of $C_{fb} = 0.24 \text{ pF}$. The fast output channel used for the timing has a pulse width of about 30 ns.

In the proposed scheme, the value of the external capacitance is limited by the electronic noise of the fast output channel. Fig. 7 shows the electronic noise at the fast output versus the external capacitance. From this figure it can be seen that an encoding system with capacitance of about 700 pF still allows a clear discrimination of the electronic noise from the pulses of interest (fission fragments have energies of, at least, 5 MeV).

In order to assess the effects of the high capacitance (700 pF) on the energy resolution for the detection of fission fragments, tests were carried out using a ^{252}Cf source. The results of these tests are shown in Fig. 1. From this figure, it is clear that the capacitance generated by the encoding system affects the energy resolution, but this effect is small and does not prevent good energy resolution data to be obtained.

The effects of such high capacitances on the time resolution were evaluated by simulating alpha particles and fission fragments with the LED (Fig. 2). In the case of alpha particles, the FWHM is about 1 ns, due to the influence of the increasing electronic noise. However, for the case of fission fragments, the FWHM is 0.8 ns. Here, the influence of the electronic noise is less important because the signal pulse amplitudes are far higher than the noise amplitudes.

4. Final remarks and conclusion

In this paper we report on the working characteristics of a silicon strip detector which is intended to be used in photofission experiments. This detector proved to be suitable for such experiments because of the excellent

spectroscopy and time characteristics for heavy ion detection. In particular, a special arrangement (large solid angle and large illuminated target spot) is proposed aiming experiments involving low counting rates.

As a matter of fact, any experiment involving the emission of two charged particles with well defined kinematics can be carried out with the proposed arrangement. Of particular interest is the study of photodisintegration of light nuclei such as lithium isotopes. Indeed, the study of reactions like $^6\text{Li}(\gamma, ^4\text{He}d)$, $^6\text{Li}(\gamma, ^3\text{He}t)$ and $^7\text{Li}(\gamma, ^4\text{He}t)$ are suitable to be investigated with such experimental setup because of the very low cross sections expected for these processes [19]. The study of such reactions at photon energies below 30 MeV could enhance our understanding of the giant resonance process in light nuclei where collective motion of the nucleons is unlikely [20] because of the strong cluster effects in such nuclei.

Finally, this work must be understood in a broad context, aiming the future experiments to be carried out at the 31 MeV cw microtron together with a photon tagging facility under construction in São Paulo.

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NUCLEAR INSTRUMENTS & METHODS IN PHYSICS RESEARCH

Section A

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Geometrical parameters of a detection system: A Monte Carlo approach

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Geometrical parameters of a detection system: A Monte Carlo approach

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Abstract

This work presents a method for the evaluation of dynamic coincidence solid angles for 2π -detector arrangements with large illuminated target area or volume, based on Monte Carlo simulation. Particular constructions useful for photonuclear experiments, using film and gas targets, are considered.

1. Introduction

New theoretical approaches [1,2] for the description of polarization observables of two-particle photodisintegration of light nuclei motivate the interest to make measurements at low photon energies, where experimental results are practically absent [3]. Such kind of experiments require quasimonochromatic photon beams and can be carried out on high duty-factor electron accelerators, associated with the photon tagging technique [4,5]. However, the tagged photon flux is limited to a maximum of about 10^7 s^{-1} . Low photon flux and small cross sections require a 2π -detection system and large illuminated target area or volume in order to maximize the photoreaction yield. Other requirements of this kind of experiment are

- (i) measurement of energy spectra of both disintegration fragments;

- (ii) use of very thin film or low-density gas targets in order to allow the coincidence detection of two particles emitted in opposite directions.

The procedure for the extraction of the differential cross section, $d\sigma/d\Omega$, from the measured yields of two-particle photodisintegration, involves, in the case of extended targets, a sophisticated evaluation of the average solid angles [6,7], even for single-arm measurements. Some particular cases for such calculations were considered in Refs. [8–10].

Two-arm measurements make the problem yet more complicated because it is necessary to take into account the reaction kinematics, that may drastically change the efficiency of the detection system. Moreover, the theoretical analysis of the data is usually done in the center of mass system (CM). The attempt to calculate analytically the average dynamic coincidence solid angles leads to ambiguities in the solutions and huge difficulties, which arise from multiple integrations over the spatial boundaries of the target and detectors [11]. A more reasonable way is to obtain the required data by Monte Carlo simulation [12] of the two-particle photodisintegration events. For this purpose we developed a computer code, which involves not only the experimental setup geometry, but also the physical parameters of the reaction that allows to obtain the average dynamic coincidence solid angles in the CM system, even for complicated detector arrangements.

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2. Definitions

The number of coincidences for the two-particle photodisintegration can be written in the following form:

$$N = \int \left\langle \frac{d\sigma^{cm}}{d\Omega}(\theta, \varphi, E_\gamma) \right\rangle N_\gamma(x, y) N_n \Delta\Omega_{cm}^{i,j}(x, y, z) dx dy dz, \quad (1)$$

where $\langle (d\sigma^{cm}/d\Omega)(\theta, \varphi, E_\gamma) \rangle$ is the differential cross section in the center of mass system, averaged over the detectors acceptances, $\Delta\Omega_{cm}^{i,j}(x, y, z)$ is the dynamic solid angle subtended by the detector pair (i, j) , from the point with coordinates (x, y, z) , considering the reaction kinematics, $N_\gamma(x, y)$ is the photon flux density, and N_n is the target density.

On the other hand,

$$N = \left\langle \left\langle \frac{d\sigma^{cm}}{d\Omega} \right\rangle \right\rangle N_\gamma^0 N_n^0 \langle \Delta\Omega_{cm}^{i,j} \rangle, \quad (2)$$

where $\langle (d\sigma^{cm}/d\Omega) \rangle$ is the differential cross section in the center of mass system, averaged over the detectors acceptances and over the target volume, $\langle \Delta\Omega_{cm}^{i,j} \rangle$ is the dynamic solid angle subtended by the detector pair, averaged over the target volume; N_γ^0 is the photon flux; and N_n^0 is the number of nuclei per cm^2 of the target. It is important to notice that

$$\left\langle \left\langle \frac{d\sigma^{cm}}{d\Omega} \right\rangle \right\rangle = \frac{1}{V} \int_V \left\langle \frac{d\sigma^{cm}}{d\Omega} \right\rangle \frac{\Delta\Omega_{cm}^{i,j}(x, y, z)}{\langle \Delta\Omega_{cm}^{i,j} \rangle} dV,$$

so that the average of the cross section over the target volume, V , is weighted on the solid angle subtended by the detector pair for each point of the target volume. In this way, the average is actually done over an effective target volume, where the solid angle is not zero. Since photonuclear cross sections usually present smooth angular dependences, this average is a good representation of the cross section in that small region.

The aim of the Monte Carlo simulation is to obtain $\langle \Delta\Omega_{cm}^{i,j} \rangle$ using the correlation between the emission angles, according to the reaction kinematics. The statistical uncertainties of the Monte Carlo simulation results have to be much less than the required accuracy of the differential cross sections, in order to have a negligible contribution to the overall uncertainties.

3. Description of the computer code

The code was written in FORTRAN-77, version 5.3. The input parameters are: beam spot shape and dimensions and tilt angle of the target (film target case) or length of the target (gas target case). In the first step three coordinates are randomly generated, determining the photon-nucleus interaction position in the lab system; see Fig. 1 (in the case of a film target, only two are needed). Then, in the center of mass system, we generate randomly and uniformly, for each event, the polar ($\cos \theta^{cm}$) and azimuthal (φ^{cm}) angles of one of the reaction products and calculate the angles of the other. So, now we have the angles θ^{cm} and φ^{cm} for both particles. The next step is the transformation of these emission angles from the CM to the lab system, according to the reaction kinematics

$$\begin{aligned} \theta_a^{lab} &= \arctg \left(\frac{1}{\gamma_c} \frac{\sin \theta_a^{cm}}{\gamma_c \gamma_b + \cos \theta_a^{cm}} \right), \\ \theta_b^{lab} &= \arctg \left(\frac{1}{\gamma_c} \frac{\sin \theta_a^{cm}}{\gamma_c \gamma_b - \cos \theta_a^{cm}} \right), \end{aligned} \quad (3)$$

where

$$\begin{aligned} \gamma_c &= \frac{S + M_A^2}{2M_A \sqrt{S}}, \quad \gamma_{a,b} = \frac{\beta_c}{\beta_{a,b}}, \\ \beta_{a,b} &= \frac{p^*}{E_{a,b}^*} \quad \text{and} \quad \beta_c = \frac{S - M_A^2}{S + M_A^2}. \end{aligned}$$

Fig. 1. Schematic view of the strip detector arrangement with a film target.

M_A , M_a and M_b are the masses of the target nucleus and of fragments a and b , respectively, and

$$P^* = \left[\frac{(S - M_a^2 - M_b^2)^2 - 4M_a^2 M_b^2}{4S} \right]^{1/2}$$

$$S = M_A^2 + 2E_\gamma M_A \quad \text{and} \quad E_{a,b}^2 = \sqrt{(P^*)^2 + M_{a,b}^2}$$

Having the reaction position and the emission angles of each particle in the lab system, the last step is the determination of the coordinates where the particles trajectories hit the planes where the detectors are located, and the analysis of the boundary conditions. As a result, the code delivers four countings: N_i and N_j are the number of particles detected, respectively, by detectors i and j ; $N_{i,j}$ the number of coincidences; and N the number of undetected particles.

So, the sorting of the emission angles of the particles is done in the CM system, but the analysis of the boundary conditions (to determine whether the particles hit or not the detectors) is done in the lab system, where it can be done in a simple way. In such an approach, the CM system average coincidence solid angle is

$$\Delta\Omega_{cm}^{i,j} = \frac{N_{i,j}}{N_{ms}} 4\pi,$$

where N_{ms} is the number of Monte Carlo starts.

4. Testing the computer code with the film target case

To check the code and to understand the influence of the main geometrical parameters in the coincidence efficiency, we reproduced the result of an analytical calculation [6,7] for the single-arm average solid angle and calculated a simple coincidence case: photofission of ^{238}U .

4.1. Comparison with the analytical calculation

Table 1 presents the results for the single-arm average solid angle subtended by a circular 10 mm diameter

Table 1
Comparison between analytical calculation and Monte Carlo results

Source diameter (mm)	Solid angle (sr)	
	Monte Carlo (this work)	Analytical calculation (Ref. [7])
2.0	0.659 (1)	0.6599
3.0	0.634 (1)	0.6341
6.0	0.558 (1)	0.5589
9.0	0.465 (1)	0.4653
12.0	0.376 (1)	0.3766
15.0	0.302 (1)	0.3030

detector, located 10 mm away from a circular, uniform and isotropic source, as a function of source diameter. Source and detector are centered over the same axis.

In the analytical calculation [7], the determination of the single-arm average solid angle of the disc is approximated using the analytical expression for the solid angle of an n -sided regular polygon of the same area. In the Monte Carlo simulation the circular source is used directly. The estimated accuracy of the analytical calculation is 0.1%. Statistical uncertainties of our simulation are about 0.2%. There is an excellent agreement between both results.

4.2. Simulation of the coincidence case

The simulation of the photofission of ^{238}U , monitored by two opposite circular detectors (10 mm diameter) located 30 mm away from each side of the target, is discussed in this section. In this case the CM and lab systems coincide.

Fig. 2(a) presents the count rates for each individual detector and the coincidence count rate versus the distance from a 45°-tilted target to one of the detectors, which is movable. The beam diameter in this case is

Fig. 2. Results of Monte Carlo simulations for the count rate dependences: (a) versus the target detector distance (dashed curve: coincidence counts; solid curve: movable detector counts; dotted curve: fixed detector); (b) versus beam radius (dashed curve: coincidence counts; solid curve: singles counts of one of the detectors).

5 mm. The target tilt angle is defined as the angle between the beam direction and the plane of the target (see Fig. 1). For asymmetric positioning, the coincidence count rate is determined by the count rate of the farthest detector and, beyond a certain distance (40 cm in our case, see Fig. 2(a)), both count rates approximately coincide. So, when we take away one of the detectors to infinity, both the singles rate of this detector and the coincidence rate coincide and become zero, but the count rate of the other detector does not change (see dotted curve in Fig. 2(a)). This result is in itself a positive test of the computer code.

Fig. 2(b) shows the singles count rate for each of the detectors and their coincidence count rate as a function of the photon beam radius, for a 45° tilted target. We assume a constant beam intensity per unit area and the singles count rate (solid line in Fig. 2(b)) is roughly proportional to the square of the beam radius. Fig. 2(b) also shows that the coincidence count rate saturates at a beam radius of about 4 mm. Illuminated target areas beyond this radius will contribute only to random coincidences, since only the single count rates increase.

5. Geometry analysis of a strip detector arrangement

The geometry optimization of a detection/target system can be described as an improvement of the geometrical parameters of the detecting array and target in order to minimize the uncertainties of the measured quantities (like cross section, Σ -asymmetry, and others), obtained for a fixed acquisition time, taking into account the physical requirements and characteristics of the experiment (angular resolution, beam, target and detectors dimensions and so on).

It is clear that the space-extended character of the target restricts the angular definition that can be achieved in a single-arm angular distribution experiment. On the other hand, for a reaction with a well-defined kinematics, it is interesting to install a second detection arm and make a coincidence measurement. In this case the cross section is determined from experimental data obtained in double-average view (see item 3), including kinematical effects, allowing a much better angular definition, since it is possible to trace back the reaction position in the target and to define which angles (θ , φ) contribute to the cross section measured by each detector pair.

The aim of this item is to optimize the detector/target geometry and to obtain (i) polar and azimuthal angular distributions on phase plane, and (ii) the average dynamic coincidence solid angle for each detector pair. Monte Carlo simulation is the only practical way to get this information.

The main characteristics of the strip detectors and associated electronics were presented in a previous paper [13]. Here we will be concerned with the relevant geo-

metrical parameters of the arrangement as a whole. We will divide our discussion in two parts, dealing with the film and gas target cases, respectively.

5.1. Film target case

Here we study the expected behavior of an experiment involving a simple angular distribution of photofission fragments, supposing a film target allowing both fragments to be emitted in opposite directions. In this case, the CM and lab systems coincide. The optimized beam-target-detector geometry (for unpolarized photons) corresponds to (see Fig. 1): $H_1 = H_2 = 16$ mm, beam dimensions are 30×30 mm² and the target tilt is 17° .

The final results for the angular distribution simulation are shown for some extreme cases of strip positioning: (i) opposite strips (without shift along beam direction), corresponding to $\theta_{lab} = 90^\circ$, Figs. 3(a)–(c) and (ii) oblique strips (maximum shift along beam direction), corresponding to $\theta_{lab} = 22^\circ$ and 158° , Figs. 4(a) and (b). All other cases lie between these ones. Figs. 3 and 4 show that the polar angle subtended by the detector pair spans a range of $\pm 4^\circ$ around its mean value, while the azimuthal angle lay inside a $\pm 45^\circ$ interval.

As mentioned before, the knowledge of the average dynamic coincidence solid angle of each strip pair is necessary for the experimental data analysis. The results obtained from the simulations are presented in Table 2. We can see that the calculated efficiencies of the strip pairs vary about 40 times through the range of strip coordinates. The maximum value is reached for opposite strips facing the middle of the target, while opposing strips that face the limits of the target are about 15% less efficient. It is interesting to note that this difference becomes much larger if the target is illuminated by a narrow (along the target length) beam. Strip pairs close to the middle of the target are much more sensitive to the dimensions of the illuminated target area than the strips close to the extremes of the target, as shown in Fig. 5. The largest solid angle difference occurs between the strip pairs (1–64) and (32, 33), corresponding to the two extreme oblique cases. This difference arises mostly because of the effective target area that corresponds to the whole target in the first case and is very small (a thin transversal strip in the middle region of the target) in the second one. This difference will not occur in an actual experimental situation, because the large solid angle of the first case comes from events in which the particles were emitted almost parallel to the target surface, thus subject to high energy losses.

The singles count rates are about 20 times higher than the coincidence rates for opposite strips and do not vary so much with strip position.

The total dynamic coincidence solid angle can be evaluated as $\Omega_{total} = 4\pi(\sum_{ij} N_{ij})/N_{mc}$, where N_{ij} is the number of coincidence counts between strips i and j and

Fig. 3. θ and φ phase plane for film target and (a) strip pair 32–64 ($\theta_{lab} = 90^\circ$, target working area is close to the upper detector plane); (b) strip pair 16–16 ($\theta_{lab} = 90^\circ$, target working area is located midway between the detector planes); (c) strip pair 1–33 ($\theta_{lab} = 90^\circ$, target working area is close to the lower detector plane).

Fig. 4. θ and φ phase plane for film target and (a) strip pair 32–33 ($\theta_{lab} = 22^\circ$, particle emission direction is almost parallel to the target. Effective target area is very close to the whole target surface); (b) strip pair 1–64 ($\theta_{lab} = 158^\circ$, very small effective target area).

Table 2

Angular acceptances and average dynamic coincidence solid angles for some strip pairs of a 2π -detector arrangement, working with a film target and detecting fragments from photofission

Detector pair	$\theta_{min}-\theta_{max}$ (degree)	$\varphi_{min}-\varphi_{max}$ (degree)	$\Delta\Omega$ (10^{-3} sr)
1–33	86.5–93.5	130–225	7.78 (5)
16–48	86.5–93.5	140–215	9.36 (6)
32–64	86.5–93.5	130–225	7.77 (5)
1–64	148–158	130–230	8.139 (5)
32–33	22–32	130–230	0.20 (1)

Fig. 5. Average dynamic coincidence solid angles for detector pairs (16–48), squares, and (1–33), dark squares, as a function of beam height. Taking as reference the point where the line joining detectors 1 and 33 crosses the target, the beam spot was centered 1 mm above this point in the first case and 16 mm above in the second one. The target tilt angle is 17° .

N_{ms} is the number of Monte Carlo starts. For the case of photofission, the setup discussed above presents a total dynamic coincidence solid angle of about 20% of 4π .

5.2. Gas target case

In this case we study the behavior of a photodisintegration experiment performed with polarized photons, namely $\gamma(^3\text{He}, \text{pd})$. The geometrical parameters are $H_1 = H_2 = 31$ mm, beam size equals 30×60 mm² and target length is 100 mm. The detector–target distance was increased in this case to improve the azimuthal angular definition, since the ability to extract the Σ -asymmetry depends on this parameter. The increase in the detector–target distance, on the other hand, decreases the statistics, so the final distance is obtained by a compromise between the two parameters.

In this case there is a significant difference between the expected angular distributions observed in the CM and lab systems, especially in the near-threshold region. Consequently, all the characteristics that are being discussed, like phase plane angular distributions, average dynamic

Fig. 6. Average dynamic single-arm solid angle as a function of strip number for the reaction $\bar{\nu}(\text{He}, \text{pd})$ and $E = 5.5$ MeV (solid curve), 6 MeV (dotted curve) and 14 MeV (dashed curve).

Fig. 7. (θ, φ) phase plane for gas target and strip pair 1–64 ($\theta_{\text{cm}} = 47^\circ$).

Fig. 8. (θ, φ) phase plane for gas target and strip pair 1–33 ($\theta_{\text{cm}} = 84^\circ$).

coincidence solid angle and even dynamic single-arm solid angle depend on the incident photon energy. Fig. 6 shows the dynamic single-arm solid angle, as a function of strip number, for three photon energies. The dependence with strip number is clearly asymmetric for $E_\gamma = 5.5$ MeV (close to the threshold), becoming symmetric at $E_\gamma = 14$ MeV.

Figs. 7 and 8 show the (θ, φ) phase plane for the two extreme combinations of detector pairs, namely pairs (1–33) and (1–64). All other combinations lie between these. The results show that (i) polar angles are within a $\pm 4^\circ$ interval for opposing pairs and within a $\pm 8^\circ$ interval for the extreme oblique case; and (ii) azimuthal angles for all combinations are within a $\pm 20^\circ$ interval. The total dynamic coincidence solid angle corresponds to 14% of 4π , for $E_\gamma = 6$ MeV.

6. Conclusions

We presented a Monte Carlo method for the determination of average dynamic solid angles and also for the study of the relevant geometrical parameters of 2π -detector arrangements for coincidence measurements.

A detailed consideration of the particular constructions for film and gas targets shows that moderate angular resolution can be achieved even with large illuminated target area or volume and large total solid angle. This is a very important characteristic to be considered in the design of two-particle photodisintegration experiments where the particles are measured in coincidence. Accurate determination of the dynamic solid angles is of fundamental importance for the experimental data analysis of such experiments.

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Triton angular distributions from the ${}^7\text{Li}(\gamma, t)\alpha$ reaction near threshold

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Angular distributions of tritons emitted in the reaction ${}^7\text{Li}(\gamma, t)\alpha$ have been measured for two photon energy intervals (6.4–6.7 and 8.5–9.0 MeV), and θ , between 30° and 150° . Monochromatic photons from neutron capture reactions of the IPEN/CNEN-SP IEA-R1 research reactor were used. The measured data were analyzed in the framework of an α -cluster model for the ${}^7\text{Li}$ structure. Results are in good agreement with a recent measurement of the azimuthal cross section asymmetry.
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The ground state of ${}^7\text{Li}$ ($J^\pi = 3/2^-$) is a typical clustered nuclear system, consisting mostly of an α -particle and a triton [1], with other components of the wave function playing secondary roles [2]. The study of cluster properties in nuclei gained a new impetus due to microscopic calculations based on the method of resonating group (MRG) [3,4] and also of multicluster models [5–8]. Those calculations adequately describe the main static characteristics of nuclear systems and achieved remarkable progress in the description of the radiative capture of clusters (like $t + \alpha \rightarrow {}^7\text{Li} + \gamma$ [9]).

According to these models the $E1$ multipole dominates the differential cross section of the ${}^7\text{Li}(\gamma, t)\alpha$ reaction at low photon energies. $E1$ transitions give rise to the S and D waves of the final state continuum:

$$\left(\frac{3}{2}^-\right)_{E1} \rightarrow S_{(1/2^+)} + D_{(3/2^+, 5/2^+)},$$

where the values in parentheses indicate total angular momentum and parity, and the D wave is split by the spin-orbit interaction into $3/2^+$ and $5/2^+$ states. Near the threshold the process is dominated by the S wave, due to the additional centrifugal barrier for the D wave, which is important in this region. At higher energies, according to MRG [9] and cluster model calculations [10], the D wave is dominant. There is a region where the contributions of the S and D waves are comparable. This transition from the dominance of the S to the D wave leads to changes in the angular distributions of the final particles in the case of excitation by unpolarized photons, or to oscillations in the Σ -asymmetry for polarized photons [10].

The aim of the present work is the accurate measurement of angular distributions of tritons emitted in the ${}^7\text{Li}(\gamma, t)\alpha$ reaction, at two characteristic energy bins: 6.4–6.7 MeV, where the S and D waves contributions are comparable, and 8.5–9 MeV, where the contribution of the D wave is dominant, in order to: (i) clarify the experimental situation of

triton angular distributions at low energies, using monochromatic photons with energies below the thresholds of contaminating reactions, and (ii) to test the prediction of the potential cluster model with attractive local potential and forbidden states [11].

The experiment has been carried out at the IPEN/CNEN-SP IEA-R1 research reactor, using monochromatic photon beams from neutron capture reactions [12]. The neutron targets were located at the midpoint of the BH4-12 tangential beam hole of the reactor (Fig. 1). The system allows (using different neutron target materials) the use of up to 30 discrete photon lines with energies from 5 to 11 MeV and energy resolution of a few eV.

The thermal neutron flux at the target position was $\sim 6.2 \times 10^{11}$ n/cm² s, monitored by a self-powered neutron detector (SPND) installed close to the neutron target. The photon flux is collimated by a long (~ 228 cm) lead collimator. In order to reduce neutron contamination, the beam is filtered after collimation by 36.5 cm of polyethylene and 9.5 cm of borated paraffin. The target consists of a 45×70 mm², self-supporting metallic natural lithium foil (92.5% ${}^7\text{Li}$ and 7.5% ${}^6\text{Li}$), 100 μm or 50 μm thick (depending on the photon

FIG. 1. Experimental setup: (1) neutron target; (2) lead collimator; (3) concrete wall; (4) polyethylene; (5) borated paraffin; (6) Li target; (7) Si detectors; (8) Ge(Li) detector.

FIG. 2. Typical triton spectrum obtained at $E_\gamma \sim 8.80$ MeV.

energy). The target was placed inside a vacuum chamber, which also housed the silicon strip detectors and a calibrated alpha source for energy calibration. The beam spot on the target was 3×6 cm², for a target tilt of 30°.

The charged products of the reaction were detected by an arrangement of seven silicon strip detectors [13], located at triton emission angles, θ_t , equal to 30°, 50°, 70°, 90°, 110°, 130° and 150°. The distance from the detector to the center of the target was 17 cm, with full detector (all 16 strips) solid angle, $\Delta\Omega$, of 0.09 sr. The detector angular acceptance (for a point target) was $\Delta\theta_p = \pm 8.4^\circ$ and $\Delta\phi = \pm 8.4^\circ$. The actual beam dimensions increased the detector angular acceptance to $\Delta\theta_p = \pm 12^\circ$ and $\Delta\phi = \pm 11^\circ$, according to a Monte Carlo simulation [14].

Two neutron targets were used in the experiment, Ti and Ni. The Ti target photon spectrum presents two strong lines: 6.418 and 6.759 MeV [15] (weighted average 6.56 MeV), which sum 80% of the intensity responsible for the emission of tritons with energies above the detection threshold. The rest of the intensity (20%) is distributed among several low intensity lines. For the Ni target, the photon spectrum presents two strong lines at 8.998 and 8.533 MeV [15] (weighted average 8.80 MeV), which concentrate 60% of the total intensity and a line at 6.837 MeV with 15% of the total intensity. The rest of the intensity (25%) is distributed among several low intensity lines.

Within the conditions of our experiment ($E_\gamma < 10$ MeV) only the reactions

and

may contribute to the measured spectra. The second one is caused by neutrons from both beam contamination and background in the experimental area, which interact with the ${}^6\text{Li}$ present in the target. Other reactions either have thresholds above the photon energies used, or negligible cross sections.

FIG. 3. Typical triton spectrum obtained at $E_\gamma \sim 6.56$ MeV.

Both reactions produce the same particles in the final states, but with different energy spectra. Due to the reaction kinematics, α -particles carry less energy than tritons. Moreover, the energy loss in the target is larger for alphas than for tritons. For those reasons, the number of α particles detected is much smaller than that of tritons, and their counts are concentrated in the low-energy part of the spectrum.

The background from reaction (2) was evaluated using a dummy neutron target (with very low intensity photons). The dominant contribution comes from thermal neutrons in the experimental area, the neutron contamination in the beam being very small. The thermal neutron background was reduced with a Cd shield around the scattering chamber.

Figure 2 shows a typical spectrum, measured using the Ni neutron target and a 150 μm thick Li target. One can clearly see the background at the lower channels and two bumps. The high energy bump, peaked around channel 150 (~ 3.5 MeV), corresponds to tritons produced by reaction (1), induced by photons from the 8.80 MeV lines; the other one, peaked around channel 90 (~ 2.3 MeV), corresponds to tritons produced both by reactions (2) and (1), the last induced by the secondary photon line at 6.837 MeV.

Figure 3 shows an analogous spectrum obtained using the Ti neutron target and a 50 μm thick Li target. This spectrum presents a single bump, formed by tritons produced by reaction (1), induced by the principal photon line (6.56 MeV) with a small contribution ($\sim 5\%$) from reaction (2). This contribution was taken into account in the analysis.

The differential cross sections obtained experimentally in this work were normalized to unity at $\theta = 90^\circ$, and are shown in Figs. 4 and 5 by the open circles. Those figures also show results from measurements using real [16] (band 1) and virtual [17] (band 2) photons. The bands were obtained by reconstructing the differential cross sections from the published Legendre coefficients. The absolute values of the cross sections presented in Refs. [16] and [17] are $\sim 20\%$ different. Both results were normalized to unity at 90° , in order to allow a better visualization. The dashed curve represents the results obtained in the framework of the binary cluster potential model [10], also normalized to unity at 90° . According to this model, the $E1$ component of the differential cross section is dominated around 9 MeV by the D scattering wave, resulting in an angular dependence with a maximum at 90° . The behavior is in good agreement with our measure-

FIG. 4. Differential cross section for $E_\gamma \sim 8.80$ MeV. This work (data points); measurements with real photons from Ref. [16] (band 1); measurements with virtual photons from Ref. [17] (band 2); binary cluster potential model calculation from Ref. [10] (dashed curve). All results are normalized to unity at 90° .

ment and that of Ref. [16], and not with Ref. [17]. According to this model [11], the contributions from S and D scattering waves are comparable around 6.5 MeV, resulting in an isotropic angular distribution. Again, there is good agreement

FIG. 5. Differential cross section for $E_\gamma \sim 6.56$ MeV (see caption of Fig. 4).

between our results and those from Ref. [16].

The agreement between measurements and the predictions of the binary cluster potential model indicates that the characteristic change in the angular distributions of the emitted tritons, which is a signature of the overflow of S and D scattering waves of $E1$ transitions, is correctly described by the model. This conclusion is corroborated by the results of a Σ -asymmetry measurement for the reaction ${}^7\text{Li}(\gamma, t)\alpha$ performed in this energy region [10].

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