## Communications

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# Laser-induced resonant absorption of $\gamma$ radiation

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This paper reports an analytic estimate developed from a multiphoton model of the cross section for the absorption of a $\gamma$-ray photon when the nuclear recoil is compensated by the simultaneous absorption of an optical photon from the radiation field of a high power laser. Tabulated data show 22 experimentally accessible nuclear transitions which might be expected to demonstrate this effect.

[NUCLEAR REACTIONS Calculated $B(E1)$ for multiphoton absorption; $\gamma$+ optical photon from laser.]

In general, the recoil of an elementary body of mass $M$ interacting with electromagnetic radiation requires an absorbed photon to have an energy higher than the transition energy of the system by an amount $E^2/2Mc^2$, where $E$ is the energy of the quantum. Since it is often the case that the best source of probe photons is the same transition excited in a similar system, for absorption to be subsequently possible it is necessary that the level width of the excited state be large compared to the recoil energy lost by the photon both during emission and absorption. For the emission of optical photons from atoms and molecules this condition is readily fulfilled and for the emission of $\gamma$ rays from nuclei it is usually not.[1]

A recent letter[2] has suggested that the recoil energy of a nucleus imparted by the absorption of a $\gamma$ photon might be effectively compensated by the simultaneous absorption of a second photon of appropriate energy. For $\gamma$ rays of energies of the order of 0.1 to 1 MeV, the sum of the recoil energies lost in emission and absorption generally corresponds to the energy of an optical photon. Thus, it has been proposed[2] that a nucleus immersed in the radiation field of a high-power laser might absorb a $\gamma$ photon emitted from a similar nucleus in another sample through a two-photon process in which the sum of the energy of the $\gamma$ photon, less the recoils, and the energy of the optical photon from the field would equal the transition energy. Absorption cross sections as large as 1 $\text{fm}^2$ were predicted for optical power densities of the order of $10^{10}$ $\text{W/cm}^2$ with not unreasonable assumptions being placed on the structure of model nuclei considered.

In the previous work it was assumed that the resonant multiphoton absorption would proceed through a nearly resonant intermediate state coupled to both initial and final states by $E1$ electric dipole transitions. The resulting rate coefficient obtained was found to depend upon $(\Delta E)^{-2}$, the inverse square of the energy defect by which the intermediate state nearest in energy to either the initial or final nuclear state missed being resonant with the optical photon. For an appreciable effect it was necessary that $\Delta E$ be of a magnitude essentially irresolvable by conventional techniques of nuclear spectroscopy, thus rendering difficult the identification of potential candidates for a demonstration.

This paper presents the estimation of the cross section for the multiphoton absorption of $\gamma$ and optical photons through intermediate states which are the magnetic sublevels of the initial and final states. Transition probabilities have been obtained which are proportional between cutoff limits to $(E_2)^{-2}$, the inverse square of the energy of the optical photon. Cross sections ranging from 0.01 to 1.0 $\text{fm}^2$ are readily obtained for odd stable nuclei in the mass regions corresponding to deformed spheres.

Conventional, second-order perturbation theory has been shown to yield, for the cross section for the absorption of $\gamma$ photons in a field of optical photons,[2]

$$\sigma = 2\Pi^3 K\alpha^2 R^4 \frac{\omega_1\omega_2}{\omega_{nn_0}{}^2} |Q_{nn_0}|^2 \frac{N_2}{\max(\Gamma_1, \Gamma_2)}, \qquad (1)$$

where $\alpha$ is the fine structure constant, $\omega_1$ and $\omega_2$ are the frequencies of the $\gamma$ and optical photons, respectively, $\omega_{nn_0}$ is the transition frequency, $N_2$ is the optical photon flux, $\max(\Gamma_1, \Gamma_2)$ is the larger of the frequency bandwidths of the two photons, $R$ is the nuclear radius ($r_0 A^{1/3}$), $K$ is the integral over line shape functions and is of the order of unity, and $Q_{nn_0}$ is a dimensionless matrix element given by

$$Q_{nn_0} = \frac{\omega_{nn_0}}{R^2} \sum_{n'} \left[ \frac{(\vec{e}_1 \cdot \vec{V}^{(1)}_{nn'})(\vec{e}_2 \cdot \vec{V}^{(2)}_{n'n_0})}{\omega_{n'n_0} - \omega_2} + \frac{(\vec{e}_2 \cdot \vec{V}^{(2)}_{nn'})(\vec{e}_1 \cdot \vec{V}^{(1)}_{n'n_0})}{\omega_{n'n_0} - \omega_1} \right], \qquad (2)$$

where $\vec{e}_1$ and $\vec{e}_2$ are the dimensionless polarization vectors of the appropriate components of the electromagnetic fields of the $\gamma$ and optical photons, respectively. The matrix elements between initial $n_0$, intermediate $n'$, and final $n$, states of the interaction of the nucleus with the $i$th radiation field $V^{(i)}_{jk}$, have been normalized by the product of field strength and electric charge to consolidate elementary constants appearing in Eq. (1) and to render $Q_{nn_0}$ dimensionless. Then, for single particle states, an $E1$ electric dipole transition becomes

$$V^{(e)} = R, \qquad (3a)$$

while for an $M1$, magnetic dipole transition,

$$V^{(m)} = (\hbar/2Mc)(\vec{L} + g_n\vec{S}), \qquad (3b)$$

where the $\vec{L}$ and $\vec{S}$ are the orbital and spin angular momentum operators, respectively, and $g_n$ is the nuclear gyromagnetic ratio of the particle undergoing the transition. To the order of approximation of the Weisskopf estimates[3] for which $(\vec{e} \cdot \vec{V}^{(e)}_{jk}) \simeq R$, it has been shown[4] that for an $M1$ transition of an "average" nucleon $(\vec{e} \cdot \vec{V}^{(m)}_{jk}) \simeq \sqrt{10}(\hbar/McR)(\vec{e} \cdot \vec{V}^{(e)}_{jk})$, and thus $(\vec{e} \cdot \vec{V}^{(m)}_{jk}) \simeq \sqrt{10}(\hbar/Mc)$. Barring the fortuitous occurrence of the type of accidental degeneracy assumed in previous work,[2] the individual terms of Eq. (2) will have appreciable magnitude only when the intermediate state is either degenerate with the initial state for which $\omega_{n'n_0} \sim 0$ or with the final state, giving $\omega_{n'n_0} \sim \omega_{nn_0}$. Since only resonant two-photon excitation is considered $\omega_1 + \omega_2 = \omega_{nn_0}$ and in either case the frequency dependence of the dominant terms becomes approximately[5] equal to $(1/\omega_2)$. However, the quantum numbers expressed explicitly in Eq. (2) are sets, $nLSJm$, and in the first case the products of matrix elements associated with the dominant terms have the form

$$V^{(1)}_{nLSJm, n_0L_0S_0J_0m'}\, V^{(2)}_{n_0L_0S_0J_0m', n_0L_0S_0J_0m_0}.$$

In the second case for which the intermediate state is degenerate with the final state, the dominant terms contain the products

$$V^{(2)}_{nLSJm, nLSJm'}\, V^{(1)}_{nLSJm', n_0L_0S_0J_0m_0}.$$

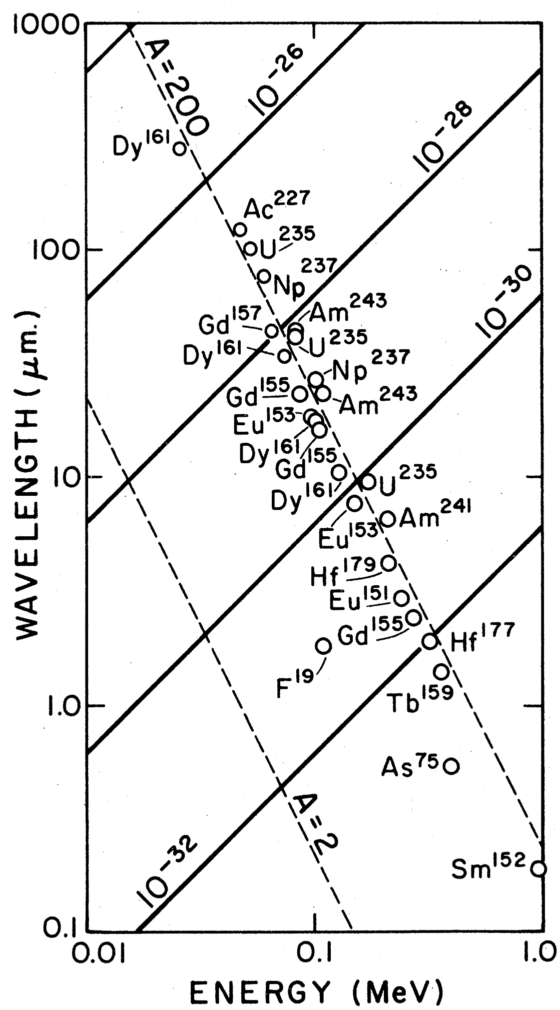


FIG. 1. Domains of the energies of $\gamma$ radiations and of the wavelengths of laser photons compensating the recoils, associated with the nuclear emission and absorption of the $\gamma$ photons induced by intense optical fields. Heavy lines locate parameter values at which the cross sections shown might be achieved in optical fields of $10^{10}$ W/cm$^2$. Dashed lines locate the domains of parameters corresponding to the induced absorption by nuclei of the mass numbers given. 0: tabulated transitions satisfying the necessary selection rules for the nucleus indicated.

The angular parts of the matrix elements contain the Clebsch-Gordan coefficients which impose selection rules on the types of transitions giving nonzero results. Since $\Delta J = 0$ for the optical transition, the largest matrix element will be obtained for a magnetic dipole, $M1$, transition for which $V^{(2)} = V^{(m)}$ and

$$m' - m_o = 0, \pm 1 \quad \text{or} \; m - m' = 0, \pm 1,$$

depending upon the polarization $\vec{e}_2$ for the above two cases, respectively. The largest matrix element for the corresponding $\gamma$ transition is obtained for an $E1$, electric dipole type for which $V^{(1)} = V^{(e)}$, assuming $J - J_0 = \pm 1, 0$. Using the substitutions suggested above gives

$$Q_{nn_0} \sim \sqrt{10}\, \frac{\omega_{nn_0}}{\omega_2} (\hbar / Mc) R^2 Q, \tag{4}$$

where $Q$ is the sum of the products of the Clebsch-Gordan coefficients over the intermediate states $n'$, consistent with the selection rules and polarizations. Since the $\gamma$ transitions is of the $E1$ type and $\omega_1 \sim \omega_{nn_0}$, the lifetime of the final state will determine the maximum width appearing in Eq. (1). The full spectral width corresponding to the value given by the Weisskopf estimates[6] is

$$\Gamma_1 = 3.3 \times 10^{13} A^{2/3} E_1{}^3 .$$

Assuming $Q \sim 1$, $\max(\Gamma_1, \Gamma_2) = \Gamma_1$ and substituting these together with Eq. (4) into Eq. (1) gives after converting units

$$\sigma = 3.5 \times 10^{-50} A^2 \Phi_2 E_1{}^{-6}, \tag{5}$$

where $\sigma$ is in $\text{cm}^2$, $E_1$ is the $\gamma$ energy in MeV, and $\Phi_2$ is the optical flux in $\text{W/cm}^2$. There is, of course, a low energy cutoff which occurs when it is no longer possible to ensure $\Gamma_2 < \Gamma_1$ while maintaining a constant laser power of $\Phi_2$.

Figure 1 plots the domains of the energy and wavelength, respectively, of the $\gamma$ and optical photons over which absorption cross sections of the values shown might be achieved in optical fields of $10^{10}$ W/cm$^2$. Specific points corresponding to tabulated transitions[7] satisfying the selection rules for the multiphoton process have been shown in the figure for nuclei of long-lived or stable species. Empirically, these can be seen to cluster around the line representing the parametric representation of the expression $A = 200$. Values of cross section ranging to 3.2 fm$^2$ for the 25.6 keV transition of $^{161}$Dy can be seen in Fig. 1.

In a realistic experiment the precise wavelength needed to compensate the recoil might not be available so that, as suggested previously,[2] the final tuning of the two-photon resonance might be done by Doppler shifting the $\gamma$ source. If $\delta E$ represents the energy defect between the available laser photon and the compensating energy needed, then to within terms of first order in $(v/c)$ the necessary velocity is approximately[2] $v \sim 2(\delta E)/Mc$, which is the velocity needed to compensate the entire recoil by the Doppler shift alone reduced by $(E_1/2\delta E)$, a large number. Thus a reasonable coincidence between the wavelength of an available laser line and the wavelength needed for precise compensation would generally result in a requirement for a velocity of practicable magnitude.

Although high peak powers have been assumed in the presentation of the results shown in Fig. 1, in an actual model experiment only the average power is important. If a geometry were arranged in which a density of $10^{22}$ cm$^{-3}$ nuclei were enclosed in a cavity resonant for the optical photons with a confinement time of the order of 1 $\mu$sec, $3 \times 10^{26}$ absorbers per cm$^2$ of laser radiation incident upon the cavity could be illuminated. Since the induced cross section is linear in the laser power, for example, at $10^9$ W/cm$^2$, the absorption probability for the $^{161}$Dy transition discussed above would reach unity per cm$^2$ of incident laser radiation, although it would be physically spread over the larger transverse area of the medium in the cavity. Then the average rate of occurrence of absorption events would depend upon the product of the gamma flux from the source and the duty cycle of the laser radiation. In the example of $^{161}$Dy for a gamma flux of $10^9$ cm$^{-2}$ sec$^{-1}$ an average absorption rate of 1 sec$^{-1}$ W$^{-1}$ would be achieved.

It can be reasonably concluded that recoil compensation through the absorption of a photon from a laser field represents a potentially viable means for implementing gamma ray absorption and fluorescence spectroscopy over the range of transition energies from 0.01 to 1.0 MeV. Tabulated transitions satisfying the selection rules span this range of energies and at the higher end correspond to laser wavelengths at which maximum powers are available from existing devices, thus pragmatically compensating the inverse dependence of the induced cross sections on energy. The fine tuning of the induced absorption resonances with the conventional techniques of Doppler shifting offers a means of obtaining much higher resolution at these energies than has been previously possible.

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[1]K. G. Malmfors and R. Mossbauer, *Alpha, Beta and Gamma-Ray Spectroscopy*, edited by Kai Siegbahn (North-Holland, Amsterdam, 1965), pp. 1281–1311.

[2]C. B. Collins, S. Olariu, M. Petrascu, and Iovitzu Popescu, Phys. Rev. Lett. 42, 1397 (1979).

[3]See for example, P. J. Brussard and P. W. M. Glaudemans, *Shell-Model Applications in Nuclear Spectroscopy* (North-Holland, Amsterdam, 1977), Chaps. 9 and 10.

[4]This approximation, of course, neglects the angular dependence which forbids both $(\vec{e}\cdot\vec{V}^{(m)})$ and $(\vec{e}\cdot\vec{V}^{(e)})$, as defined, from taking simultaneous nonzero values. The approximation taken from Ref. 3 concerns only the radial integrals and assumes the angular parts will be separately treated.

[5]A more detailed development of the transition amplitudes of Eq. (2) would generally remove the explicit divergence of the order of $\omega_2^{-1}$ resulting from transitions for which $n'=n_0$ or $n$. See, for example Eq. (54) of S. Olariu, Iovitzu Popescu, and C. B. Collins, Phys. Rev. D (to be published). If either of the initial or final levels were degenerate, there would be a dependence upon $(\omega_2-\Delta)^{-1}$ resulting from transitions through the degenerate intermediate states, where $\Delta$ is the shift of the levels associated with the removal of the degeneracy.

[6]This is the estimate for a single photon transition for which $(\vec{e}\cdot\vec{V}^{(e)}_{nn_0})\sim R$ and is thus consistent with the order of approximation used in this work.

[7]C. M. Lederer and V. S. Shirley, *Table of Isotopes*, seventh edition (Wiley, New York, 1978).