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Laser Enhancement of Nuclear β Decay

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(Received 20 July 1981)

Nuclear $\beta^{\mp}$ decay in the presence of a very intense laser field is considered. If the energy which is available for the decay in the absence of the field is small, the impact of the laser field is very significant, resulting in a large enhancement of the decay and a large upshift of the $\beta^{\mp}$ energies.

PACS numbers: 23.40.-s, 13.40.Ks, 42.60.-v

Applications of lasers to nuclear physics have thus far been restricted to measurements of hyperfine or isotope shifts of atomic levels which provide information about nuclear quantities.[1] Directly influencing nuclear phenomena by a laser field seems hopeless at first glance. However, laser fields (albeit well focused) are of practically infinite extent compared to nuclear dimensions. Hence, although nuclear matrix elements are hardly affected by the laser, the quantum states of charged particles which evolve from nuclear processes are. For example, the state of a free electron in the presence of a laser field is given by the Volkov solution [Eqs. (2) and (3)].

The basic dimensionless parameter which turns out to govern all these effects is (we use units such that $\hbar = c = 1$)

$$\nu^2 = (ea/M)^2 = 7.5\times10^{-11}\lambda^2[\mathrm{cm}]I[\mathrm{W/cm^2}](m/M)^2, \qquad (1)$$

with a, λ, and I the field strength of the vector potential, the wavelength, and the intensity of the laser field; M is the mass of the charged particle and m the electron mass. Since in the presence of a recoiling agent a charged particle can gain energy from the laser field (or lose energy to it) the phase space available for nuclear reactions can be increased. In view of Eq. (1) this applies mainly to reactions with an electron or positron taking part. For $m=M$, $\nu^2 \gtrsim 1$ can be comparatively easily achieved with present-day high-energy pulsed lasers.[2] Hence nuclear β decay seems to be a most promising application of these ideas. What has been said applies to electron scattering off nuclei or Coulomb excitation as well. In this case, however, there is ample phase space anyway, so that the influence of the laser field is minor.

Decays of elementary particles in laser fields have been the subject of extensive investigations before.[3-7] In this Letter we provide basic estimates for nuclear β decay in an intense laser field, leaving details for a future publication.

Charged particles in a plane-wave field of infinite extent are described by the Volkov solution of the Dirac equation.[8] For a circularly polarized monochromatic field $\vec{A}=a(\vec{e}_1 \cos kx + \sigma \vec{e}_2 \sin kx)$, this is

$$\psi(p, s\,|\,x) = \sum_{n=-\infty}^{\infty} \exp[-i(\tilde{p}-nk)x]\exp(-in\psi\sigma)V_n(p)u(p,s), \tag{2}$$

$$V_n(p) = J_n(z) + (ea/4pk)\not{k}[(\not{e}_1 - i\not{e}_2\sigma)J_{n-1}(z)\exp(i\psi\sigma) + (\not{e}_1 + i\not{e}_2\sigma)\exp(-i\psi\sigma)J_{n+1}(z)]. \tag{3}$$

Here $k=(\omega, 0, 0, \omega)$ is the wave vector of the laser field, p the momentum of the charged particle outside of the field ($p^2=M^2$), and

$$\tilde{p} = p + [(ea)^2/2pk]k \tag{4}$$

its effective momentum inside the field $[\tilde{p}^2 = M_*^2 = M^2 + (ea)^2]$. $u(p, s)$ satisfies the free Dirac equation, $(\not{p}-M)u(p,s)=0$, and $\exp(2i\psi)=(p^1+ip^2)/(p^1-ip^2)$, $z=eap_\perp/pk$, $p_\perp=(p_1^2+p_2^2)^{1/2}$, with p^1 and p^2 the components of the momentum transverse to the laser field. $\sigma=\pm1$ denotes the direction of polarization. Since the Bessel functions become maximal for $n \sim z$, the most probable energy transfer with the laser is of the order of $z\omega \sim (ea/m)(m/\omega)\omega \sim 10^6\omega$. Hence we can expect an increase in the available decay energy of about 1 MeV for $ea/m \sim 1$.

To calculate the transition rate of nuclear β decay in the presence of a laser field, we simply replace within the usual matrix element (cf. Shalit and Feshbach[9]) the wave function $\exp(-ipx)u(p,s)$ of the free electron by the Volkov solution Eqs. (2) and (3). For simplicity, we shall consider the allowed β decay of unpolarized nuclei. We then find by standard methods, neglecting the recoil of the nucleus,

$$dw_{fi} = \sum_{n=-\infty}^{\infty} \frac{g^2}{8\pi^4}\frac{1}{2J_i+1}\sum_M (C_V|^2|M_F^2| + |C_A|^2|\vec{M}_{GT}|^2)$$
$$\times m^3(\epsilon^2-1)^{1/2}\theta\left(\epsilon_0 - \epsilon - \frac{\nu^2}{2\Delta\epsilon} + n\frac{\omega}{m}\right)L_n(\epsilon_0, \epsilon, \theta)\,d\epsilon\, d\Omega_{\vec{p}_e}, \tag{5a}$$

$$L_n(\epsilon_0, \epsilon, \theta) = \left(\epsilon_0 - \epsilon - \frac{\nu^2}{2\Delta\epsilon} + n\frac{\omega}{m}\right)^2\Bigg[\left(\epsilon - n\frac{\omega}{m}\right. J_n^2 + \frac{\nu^2}{4\Delta\epsilon}(J_{n+1}^2 + J_{n-1}^2)$$
$$+\sigma\nu\left(\frac{n}{(\epsilon^2-1)^{1/2}}\frac{\omega}{m\sin\theta} - \frac{(\epsilon^2-1)^{1/2}\sin\theta}{\Delta\epsilon}\right)J_n J_n'\Bigg], \tag{5b}$$

$$\Delta\epsilon = \epsilon - (\epsilon^2-1)^{1/2}\cos\theta, \quad \theta = \sphericalangle(\vec{p}_e, \vec{k}). \tag{5c}$$

Here the argument of the Bessel functions is $z = eap_{\perp e}/p_e k$. The normalized energy (including the rest mass) of the outgoing electron is $\epsilon = p_e^0/m$, and $\epsilon_0 = [M(Z,N) - M(Z\pm1, N\mp1)]/m$ is the difference of the *nuclear* masses, for $\beta^\mp$ decay, respectively. Equations (5a) and (5b) apply for both β^- and β^+ decay, provided the appropriate Fermi and Gamow-Teller matrix elements, $M_F = \langle f|M_F^{(\pm)}|i\rangle$ and $\vec{M}_{GT} = |f|M_{GT}^{(\pm)}|i\rangle$, respectively, are inserted. The sum over M_i is over the spin component of the target nucleus with total spin J_i. Equations (5a) and (5b) show that in the case of allowed unpolarized β decay the laser field does not discriminate between Fermi and Gamow-Teller transitions.

The sum over n in Eq. (5a) is a remainder of the Volkov solution Eq. (2). Because of the θ function in Eq. (5a) it has a lower limit n_0:

$$n \geq n_0 = m(\epsilon + \nu^2/2\Delta\epsilon - \epsilon_0)/\omega. \tag{6}$$

$n_0 > 0$ indicates that a minimum number n_0 of photons have to be absorbed from the laser field in order that the decay (for specified ϵ and $\Delta\epsilon$) can take place. For $n_0 < 0$ a maximum number $|n_0|$ can be emitted into the field. In each case the energy balance is restored by the neutrino which has been integrated over.

If $n_0 \ll 0$, there is ample phase space available for the decay, and we do not expect (and did not find) significant enhancements due to the laser field. Let us then consider the case $n_0 \approx 0$, where phase space is tight without the laser. We then, to a fair approximation, sum from zero to infinity in Eq. (5a).

In principle, the emitted electrons can have arbitrarily high energies by absorbing sufficiently many photons out of the laser field. If, however, the minimum number n_0 [Eq. (6)] which is required for that exceeds the argument z of the Bessel functions significantly, these contributions will be practically zero. Hence we will assume that an upper limit is given by the condition $n_0 \leqslant z$, or explicitly

$$\Delta\epsilon(\epsilon-\epsilon_0)+\tfrac{1}{2}\nu^2 \leqslant \nu(\epsilon^2-1)^{1/2}\sin\theta. \qquad (7)$$

If the left-hand side of Eq. (7) is positive, this yields, for specified ϵ, limits on the possible angles: $0<\theta_1 \leqslant \theta \leqslant \theta_2 < \pi$. Taking the equality sign in Eq. (7) we can solve for $\cos\theta_1$ and $\cos\theta_2$. Real solutions are only obtained if ν is in the interval $\nu_{\min} \leqslant \nu \leqslant \nu_{\max}$, where

$$\nu_{\max,\,\min} = 2\{\epsilon\epsilon_0 - 1 \pm [(\epsilon^2-1)(\epsilon_0^{\,2}-1)]^{1/2}\}. \qquad (8)$$

Equation (8) shows that if a nucleus is β stable without the laser field ($\epsilon_0 < 1$) it practically stays so even in a strong field (apart from contributions with $n_0 > z$).

We now give two examples in which the energies which are available for β decay are low so that we can expect most pronounced effects due a laser field. For ^3H the maximum β^- energy is 18 keV, and hence $\epsilon_0 - 1 = 0.036$. Figure 1, which is a graph of Eq. (8), exhibits that for $\nu > 0.27$ there is also a minimum kinetic energy for the emitted electrons which, for $\nu = 1$, is almost ten times the maximum energy without the laser field. Figure 2 gives the energy distributions of the emitted electrons obtained by integrating Eq. (5a) within limits $\theta_1 \leqslant \theta \leqslant \theta_2$ according to Eq. (7).

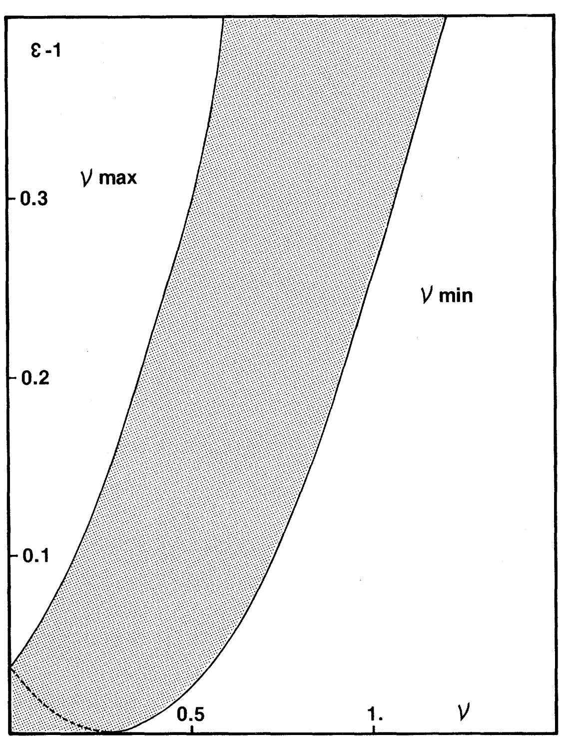

FIG. 1. Allowed final energies of the electron in ^3H decay lie in the shaded part between the curves $\nu_{\max}$ and $\nu_{\min}$ given by Eq. (8). The dashed part of the curve $\nu_{\min}$ does not have physical significance.

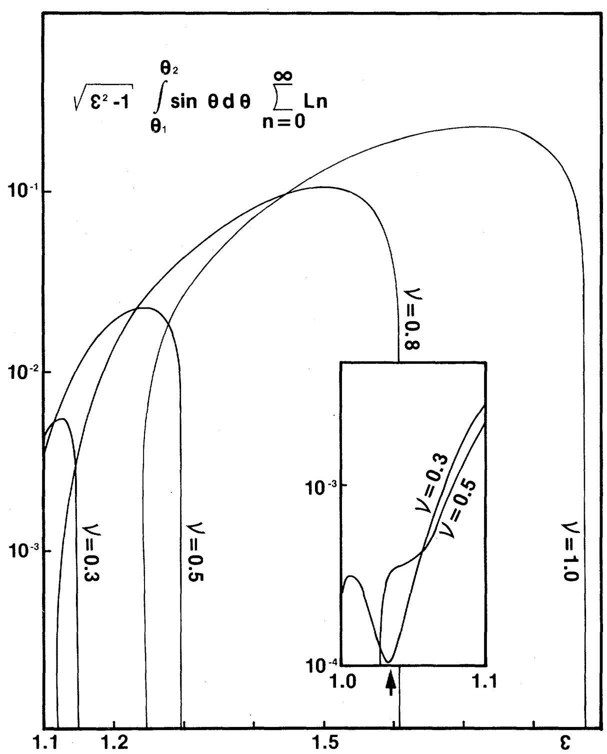

FIG. 2. Logarithmic plot of the electron spectrum of ^3H decay for $\nu = 0.3$, 0.5, 0.8, 1.0. The inset at the lower right shows the values for $1 \leqslant \epsilon \leqslant 1.1$ with doubled scale of the abscissa.

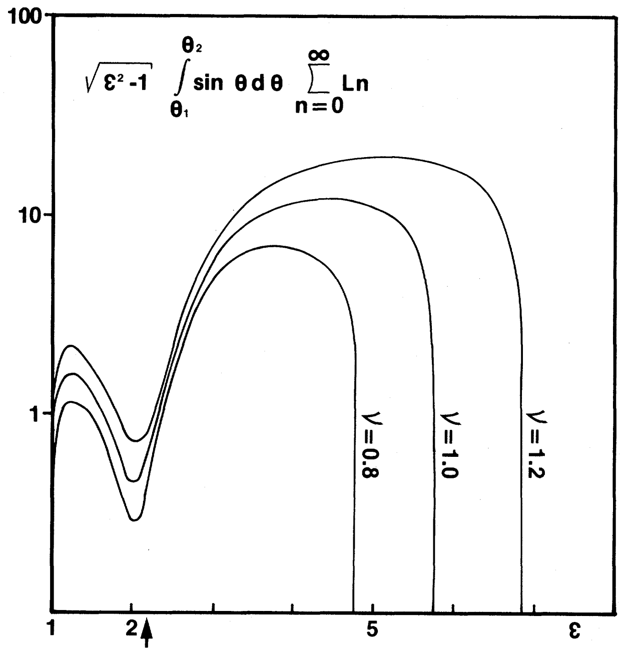

FIG. 3. Logarithmic plot of the electron spectrum of ^{18}F β^+ decay for $\nu=0.8$, $\nu=1$, and $\nu=1.2$. The arrow indicates the maximum energy in the absence of the laser field.

The integration has been done analytically. The most prominent feature is the shift of the final electron energies to much higher values. The curve for $\nu=0.3$ shows a double-peaked distribution[10] of which for $\nu=0.5$ only a small shoulder survives.

The enhancement of the total decay rate, i.e., the ratio of the total decay rates with and without the laser field, for ^{3}H decay is 10^2, 8×10^2, 7×10^3, and 2×10^4 for $\nu=0.3$, 0.5, 0.8, and 1.0, respectively.[11] Some heavy nuclei, e.g., $^{210}_{82}$Pb, $^{228}_{88}$Ra, and $^{241}_{94}$Pu, have lifetimes and β^--decay energies comparable to $^3_1$H, so that we can expect enhancements of the same order of magnitude although the question of Coulomb corrections deserves a more careful analysis in these cases.

For β^+ decay the situation is less favorable, because, if phase space is tight, i.e., the final positron energies are low, electronic capture instead of β^+ decay is dominant. An example in which the final energies are still comparatively low, but electronic capture can be neglected, is $^{18}_9$F with $\epsilon_0=2.29$. In Fig. 3 the positron energy distribution derived again by integrating Eq. (5a) is plotted for $\nu=0.8$, 1.0, and 1.2; the corresponding enhancements are $R=7$, 18, and 40, respectively. In this case the double-peaked spectrum which we have already encountered for $^3_1$H decay, is very obvious, the influence of the laser on the low-energy part being only minor.

Some remarks of caution are in order: (a) In deriving all of our explicit results we took $n_0=0$ as the lower limit of summation so that the summation over n could be done analytically. This tends to underestimate the low-energy part and to overestimate the high-energy part of the spectrum, although it might be partially compensated by the fact that there are also contributions from terms with $n_0>z$ which we neglected too. Consequently, the enhancement ratios given above should be considered as order-of-magnitude estimates. (b) The applicability of the Volkov solution [Eqs. (2) and (3)], which is valid for an infinitely extended laser field, to the case of a focused laser pulse is crucial. Hence the laser pulse should be taken as long and the spot size as large as possible in view of the required extreme intensity. (c) Coulomb corrections have been neglected. Their general tendency should be the same as in the case without a laser field.

We benefited from discussions with C. Chandler. This work has been supported by the U. S. Air Force Office of Scientific Research, the Max-Planck-Gesellschaft zur Förderung der Wissenschaften e.v., and the Alexander von Humboldt Stiftung.

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[10]The double-peaked distribution has been observed in

Ref. 5 for free neutron decay. It is difficult to compare our results with Ref. 5 since we expressed all quantities in terms of the electron energy ϵ outside of the laser field whereas the author of Ref. 5 chose the effective energy inside the laser field. Moreover, Ref. 5 contains several obvious misprints.

[11]E.g., for the Garching iodine laser we have $\nu^2 \sim 1$ in a spot size of 10 μm. Under the assumption that the field penetrates 1 μm into the sample, during the 300-psec pulse duration the number of positrons from a ^{18}F sample would increase from ~ 0.05 when the laser is off to ~ 1 when the laser is on. Thus time coincidence measurements along these lines are feasible even though the laser duty cycle is very low.

Isospin Splitting of Isovector High-Spin "Stretched" Particle-Hole Excitations in Non-Self-Conjugate Nuclei

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A simple Lane model is used to parametrize the energy systematics of the isospin splitting of high-spin magnetic states in non-self-conjugate nuclei. A strength parameter $V_1 = 106 \pm 10$ MeV is found.

PACS numbers: 25.30.Cg, 21.10.Dr, 21.10.Hw

Although the energy systematics of the isospin splitting of the giant electric[1,2] dipole and monopole[3] states have been well studied, very little is known about the isospin splitting of unnatural parity or magnetic states in nuclei. It is of particular interest to focus on high-spin "stretched" magnetic states [e.g., $(g_{9/2}f_{7/2}^{-1})_{8^-}$] where in electron scattering the cross section[4,5] is dominated by a single particle-hole configuration. Since collective effects are unimportant, the isospin splitting energy may be more directly related to a simple symmetry energy formula of the Lane type[6] with a symmetry energy parameter V_1 characteristic of a single-particle value. In this Letter we present measurements on $M8$ transitions in the $T=2$ nucleus ^{60}Ni and use these results together with those from other $T_z \neq 0$ nuclei to determine the strength parameter V_1 and to present evidence that V_1 is approximately multipole independent for magnetic states in non-self-conjugate nuclei.

The electron scattering data were taken at the Bates Linear Accelerator. Electrons were scattered from a 99.79%-enriched ^{60}Ni target of thickness 20.06 mg/cm^2 and dimensions 5×5 cm^2. An energy spectrum of electrons scattered from ^{60}Ni at $E=205$ MeV and $\theta_{\rm lab}=140°$ is shown in Fig. 1. Other incident energies were chosen at angles of $\theta_{\rm lab}=140°$ and $\theta_{\rm lab}=180°$ such that data covered the momentum transfer region 1.5 to 2.7 fm^{-1}. The scattered electrons were momentum analyzed in the Bates Linear Accelerator–Massachusetts Institute of Technology high-resolution magnetic spectrometer[7] and detected in an array of multiwire gas and Čerenkov counters.[7] Energy calibrations were made with peaks of known excitation energies. Overall normalization was obtained by comparing elastic scattering from ^{60}Ni and H