FEB 01 1988

UCRL- 96759, Rev. 1
PREPRINT

CONF-880110--16-Rev 1

NUCLEAR EXCITATION VIA THE MOTION OF ELECTRONS IN A STRONG LASER FIELD

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THIS PAPER WAS PREPARED FOR SUBMITTAL TO
THE O-E LASE 88 CONFERENCE
Los Angeles, CA
10-17 January, 1988

December, 1987

Lawrence Livermore National Laboratory

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UCRL--96759-Rev.1

DE88 005223

Nuclear excitation via the motion of electrons in a strong laser field

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ABSTRACT

A method of switching from a nuclear isomeric state to a lasing state is examined. A semi-classical model of laser-electron-nuclear coupling is developed. In it the electrons are treated as free in the external field of the laser, but with initial conditions corresponding to their atomic orbits. Application is made to testing this model in ^{235}U and to the design criteria of a gamma-ray laser.

1. INTRODUCTION

Much effort has been devoted in recent years to the search of physical mechanisms and techniques that would permit the developement of a gamma-ray laser.[1] In one of the most promising schemes proposed up to now, energy is first stored in a long-lived (isomeric) nuclear state. Then a fast transition to a short-lived level is induced by some external stimulation. It is presently believed that a certain number of conditions have to be met in order to implement this scenario in practice. In particular, it seems at least desirable that the second step - the lasing phase - populate the upper level of a recoiless Mossbauer-type transition. However, whatever the precise sequence of nuclear levels involved, one basic requirement is to be able to produce a significant population of the storage level, and then switch it to the lasing level in a time shorter than its lifetime.

The purpose of this paper is to analyze the possibilities of a method where the isomer population would be switched to the lasing state by the motion of the atomic electrons in the field of a intense laser. In fact, it seems that highly localized electron current densities can be induced in multi-electron atoms by powerful lasers[2] and may couple to nuclear excitations.[3] A direct excitation of low-lying nuclear levels is therefore conceivable and would be a powerful test of this mechanism. On the other hand, we expect the nuclear transition probability per unit time to be very small since the driven electrons move at average distances from the nucleus comparable to atomic dimensions.

We must emphasize that we have no reason to utilize this method to excite low-lying nuclear states. That is solely a test of the method. If this laser-electron-nucleus coupling has any value, it will be in switching from an isomeric to a lasing state. We will discuss the conditions in which that could take place later in the paper.

In the present analysis of this problem, we do not attempt to set up a

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precise description of the laser + electrons + nucleus system. An exact solution is beyond our imagination and a sophisticated approximation (TDHF) is both in progress and very difficult.[4] Our aim is rather to determine if a laser-electron-nucleus coupling mechanism has any chance at all to be useful for a gamma-ray laser. In this spirit, we develop a very simple model for the electron motion which always errs in overestimating the electron-nucleus coupling, but nevertheless has the structure of a more realistic theory. Namely, we assume that the atomic electrons are set into motion by the electric field of the laser as if they were free. They are hypothesized to move on classical trajectories starting at their location in the atom. The nuclear transition probability derived with these assumptions is therefore an upper limit that is indicative of the possibilities of the method and affords an opportunity to examine these effects as a function of the pertinent parameters. We are free to include or not include individual shells and thus can examine both the effects of atomic binding (relative to the laser energy) and possible coherent electron motion.

As a test case, we choose the $\Delta I=3$ transition from the ground state to the 75 eV, 26 mn isomeric level of ^{235}U which has also been discussed by SOLEM et al.[5] in a more simplified model which treats the electron cloud as a rigid body. The frequency of this transition being only ~ 15 times that of presently available U-V lasers,[6] it has been suggested as a very good candidate for testing this mechanism. However, our study extends to more general cases since we determine the dependence of the transition probability on the frequency and multipolarity of the nuclear transition. We also analyze the influence of the characteristics (frequency and intensity) of the laser.

In the second section we develop our model. The results obtained for the nuclear transition, probability with different sets of parameters, are presented in the third section. In the fourth one, we apply them to the case of ^{235}U. The fifth section details design considerations for a realistic gamma-ray laser with the provision that our calculations underestimate the power and frequency requirements of the mechanism. Conclusions are drawn in the sixth section.

2. MODEL

As mentioned in the introduction, we assume that the effect of the laser is to set into motion all the electrons of the atom as if they were free. In other words, the influence of the Coulomb field of the nucleus as well as the shielding due to the external atomic shells are ignored. In the same spirit, each portion of the electron cloud is supposed to move along a classical trajectory. That is, we choose simplistic initial conditions, namely that each electron is a classical gas whose initial position is governed by its quantum spatial distribution. Neglecting relativistic retardation and re-radiation effects, the motion of the electrons is due entirely to the electric field E of the laser. We assume the laser linearly polarized (E parallel to z) and the wavelength much larger than the atom size (which is true for ultraviolet and soft x-ray lasers). The z-component of the electric field can therefore be written $E_z = E_0 \sin \omega_0 t$ with ω_0 the laser

frequency. With these assumptions, each portion of the electron cloud initially located at (x_o, y_o, z_o) acquires a harmonic motion on a linear trajectory:

$$X_e(t) = X_o$$
$$Y_e(t) = Y_o \tag{1}$$
$$Z_e(t) = Z_o - \gamma \sin \omega_o t$$

with the same amplitude

$$\gamma = \frac{e}{m} \frac{E_o}{\omega_o^2} \tag{2}$$

(e and m are the electron charge and mass respectively). Denoting by $e\rho_o^{(i)}$ charge distribution of the i-th electron before excitation and by $e\rho_o^{(i)}$ its charge distribution when it moves according to (1), we have:

$$\rho_t^{(i)} (\vec{r}_e (t)) = \rho_o^{(i)} (\vec{r}_o)$$

with

$$\vec{r}_e(t) = (x_o, Y_o, Z_o - \gamma \sin \omega_o t) \quad {}_o \quad \vec{r}_o = (x_o, Y_o, Z_o)$$

The Coulomb field created at point $\vec{r}$ by the i-th moving electron is therefore:

$$d^3 r_e(t) \frac{\rho_t^{(i)} (\vec{r}_e(t))}{|\vec{r} - \vec{r}_e(t)|} \equiv d^3 r_o \frac{\rho_o^{(i)} (\vec{r}_o)}{|\vec{r} - \vec{r}_e(t)|}$$

and the time-dependent electron-nucleus interaction is

$$\bar{V}(t) = \sum_{i=i}^{Z} \int d^3 r_o \, \rho_o^{(i)} (\vec{r}_o) \times$$
$$\times \sum_{\sigma} d^3 r \, \Psi^+_{prot} (\vec{r}, \sigma) \frac{e^2}{|\vec{r} - \vec{r}_e(t)|} \Psi_{prot} (\vec{r}, \sigma) \tag{3}$$

where $\Psi^+_{prot} (r, \sigma)$(resp. $_{prot} \Psi (r, \sigma)$ is the creation (resp. annihilation) operator of a proton of spin σ at point r. The Hamiltonian of

the nucleus in the external field created by the Z moving electrons now is

$$\hat{H} = \hat{H}_n + \hat{V}(t)$$

where H_N is the nuclear Hamiltonian, and the time-dependent state $|\phi(t)\rangle$ of the nucleus can be expanded on the eigenstates $|x_\nu\rangle$ of H_N:

$$|\phi(t)\rangle = \sum_n b_n(t)\, e^{-\frac{i}{\hbar} E_n t} |x_n\rangle$$

$$\hat{H}_N |x_n\rangle = E_n |x_n\rangle$$

We assume the nucleus initially in state $|x_0\rangle$ (with angular momentum I_0). The probability of finding the nucleus in state $|x_1\rangle$ (with angular momentum I_1) at time T_1 is

$$P(T) = \frac{1}{2I_0 + 1} \sum_{M_0 = -I_0}^{I_0} \sum_{M_1 = -I_1}^{I_1} |\, b_1(T)\, |^2 \tag{4}$$

with

$$b_n(T) = \delta_{n,0} + \frac{1}{i\hbar} \int_0^T dt \sum_m \langle x_n|\hat{V}(t)|x_m\rangle\, e^{\frac{i}{\hbar}(E_n - E_m)t}\, b_m(t) \tag{5}$$

We first evaluate the nuclear transition probability after one laser cycle of duration $T_0 = 2\pi/\omega_0$. Since T_0 is of the order of 10^{-14} - 10^{-15} s and the electron-nucleus interaction is expected to be small, (5) can be evaluated using first order perturbation theory. Then:

$$b_1(T_0) \simeq \frac{1}{i\hbar} \int_0^{T_0} dt\, \langle x_1|\hat{V}(t)\, |x_0\rangle\, e^{i\omega_N t} \tag{6}$$

with

$$\omega_N = (E_1 - E_0)/\hbar$$

Inserting the expression (3) of V(t) and expanding the Coulomb potential into spherical harmonics gives

$$b_1(T_o) = \sum_{i=1}^{Z} \int d^3r_o\, \rho_o^{(i)}(\vec{r}_o) \sum_{\lambda=0}^{\infty} \sum_{\mu=-\lambda}^{\lambda} \frac{4\pi e^2}{2\lambda+1} \frac{1}{i\hbar} \int_0^{T_o} dt\, e^{i\omega_N t}\, Y_\lambda^\mu(\hat{r}_e(t))$$

$$\times \sum_\sigma \int d^2\hat{r}\, Y_\lambda^{\mu *}(\hat{r}) \left(\int_0^{r_e(t)} r^2 dr \frac{r^\lambda}{[r_e(t)]^{\lambda+1}} \langle x_1 | \Psi^+_{prot}(\vec{r},\sigma)\, \Psi_{prot}(\vec{r},\sigma) | x_o \rangle \right.$$

$$\left. + \int_{r_e(t)}^{\infty} r^2 dr \frac{[r_e(t)]^\lambda}{r^{\lambda+1}} \langle x_1 | \Psi^t_{prot}(\vec{r},\sigma)\, \Psi_{prot}(\vec{r},\sigma)\, x_o \rangle \right. \qquad (7)$$

We now assume that the nuclear transition matrix element

$$\langle x \mid \Psi^t_{prot}(\vec{r},\sigma)\, \Psi_{prot}(\vec{r},\sigma) \mid x_o \rangle$$

is peaked around the nuclear surface. This is certainly valid for transitions between low-lying nuclear states since they involve essentially single-particle or surface collective excitations. With this assumption, the first integral in the bracket is non-zero only if $r_e(t)$ lies outside the nucleus and the second one is always very small in comparison to the first: it contains the factor $(r_e(t)/r)$, with $r_e(t)$ smaller than the nuclear radius (we exclude in this study the case of $\lambda = 0$ transitions). Keeping then only the first integral, its upper limit can be extended to infinity. This leads to

$$b_1(T_o) \frac{4\pi e}{i\hbar} \sum_{i=1}^{Z} \int d^3r_o \rho_o^{(i)}(\vec{r}_o) \sum_{\lambda\mu} \frac{S_{\lambda\mu}(T_o \vec{r}_o)}{2\lambda+1} \langle x_1 | \hat{Q}_{\lambda\mu} | x_o \rangle$$

where

$$Q_{\lambda\mu} = \sum_\sigma \int d^3r\, \Psi^+_{prot}(\vec{r},\sigma)\, er^\lambda\, \Psi_{prot}(\vec{r},\sigma)\, Y_\lambda^{\mu *}(\hat{r})$$

is the electric multipole operator of order λ and

$$S_{\lambda\mu}(T_o \vec{r}_o) = \int_0^{T_o} dt\, \Theta(r_e(t) - R_n)\, e^{i\omega_N t} \frac{Y_\lambda^\mu(\vec{r}_e(t))}{[r_e(t)]^{\lambda+1}}$$

Here Θ is the usual step function ($\Theta(x) = 1$ if $x < 0$ and $\Theta(x) = 0$ if $x < 0$) and we have assumed a spherical nucleus of radius R_N.

Using (8) and (4), the transition probability between times 0 and T_c is

$$P(T_0) = \left(\frac{4\pi e}{\hbar}\right)^2 \sum_{i,i'=1}^{Z} \int d^3r_0 \int d^3r_0' \, \rho_0^{(i)}(\vec{r}_0) \, \rho_0^{(i)'}(\vec{r}_0')$$

$$\times \sum_{\lambda\mu} \sum_{\lambda'\mu'} \frac{S^*_{\lambda\mu}(T_c,\vec{r}_0)}{2\lambda+1} \; \frac{S_{\lambda'\mu'}(T_0,\vec{r}_0')}{2\lambda'+1} \times A$$

The quantity

$$A = \frac{1}{2I_0+1} \sum_{M_0} \sum_{M_1} \langle x_1|\hat{Q}_{\lambda\mu}|x_0\rangle \; \langle x_1|\hat{Q}_{\lambda'\mu'}|x_0\rangle$$

is readily expressed with the Wigner-Eckart Theorem and the definition

$$B(E\lambda,\, I_0 \to I_1) = \sum_{\mu=-\lambda}^{\lambda} \sum_{M_1=-I_1}^{I_1} |\langle x_1|\hat{Q}_{\lambda\mu}|x_0\rangle|^2 \tag{10}$$

of $B(E\lambda;\, I_0 \to I_1)$ [7]. One obtains:

$$A = \delta_{\lambda\lambda'}\delta_{\mu\mu'} \; B(E\lambda;\, I_0 \to I_1) \,/\, (2\lambda+1)$$

Taking the lowest non-zero value of λ as the most probable multipolarity of the transition, the transition probability finally is:

$$P(T_0) = \left(\frac{4\pi e}{\hbar}\right)^2 \sum_{\mu=-\lambda}^{\lambda} \int \sum_{i=1}^{Z} d^3r_0 \, \rho_0^{(i)}(\vec{r}_0) \, S_{\lambda\mu}(T_0\,\vec{r}_0)\,|^2 \frac{B(E\lambda)}{(2\lambda+1)^3} \tag{11}$$

with

$$\lambda = |I_1 - I_0| \neq 0$$

In the following we shall first base our discussion on the partial transition probability:

$$P(T_0,\vec{r}_0) = \left(\frac{4\pi e}{h}\right)^2 \sum_{\mu=-\lambda}^{\lambda} |\; S_{\lambda\mu}(T_0,\vec{r}_0)\;|^2 \frac{B(E\lambda)}{(2\lambda+1)^3} \tag{12}$$

and the one-electron transition probabilities:

$$p^{(i)}(T_0) = \int d^3r_0 \, \rho_0^{(i)}(\vec{r}_0) \; P(T_0,\vec{r}_0) \tag{13}$$

How the complete transition probability (11) can be related to (12) and (13), that is the effect of the coherence of the motion of the electrons, will be discussed in the next Section.

Finally, the transition probabilities (12) and (13), derived for one laser cycle T_0, will be readily extended to one laser pulse $T_p = N \cdot T_0$ ($N \approx 300$). In fact, as we shall see in the next Section, the transition probabilities for one cycle are small and first order perturbation theory is still valid during one laser pulse. Replacing then T_0 by T_p in (12) and using the periodicity of the integrand of $S_{\lambda\mu}(T_0, \vec{r_0})$, one finds that $P(T_p, r_0)$ is equal to $N.P(T_0, r_0)$ if $\hbar\omega_N/\hbar\omega_0$ is close to an integer and vanishingly small otherwise. Consequently, during one laser pulse, only those nuclear transitions whose energies are an integer number of times the laser photon energy can be excited by the moving electrons.

3. RESULTS.

We first study the behaviour of the probability $P(T_0, \vec{r}_0$ of eq. (12). This quantity depends on:

- the parameters $r_{\perp 0} = (x_0^2 + y_0^2)^{1/2}$ and z entering $\vec{r}_0$ (the detailed expressions developed in the Appendix of Ref. 14 show that (12) depends on x_0 and y_0 only through the combination $r_{\perp 0}$)
- the energy $\hbar\omega_N$ and the multipolarity λ of the nuclear transition.
- the frequency ω_0 and the maximum electric field strength E_0 of the laser (treated as an incident classical external electro-magnetic field).

In most applications, we suppose that the laser works in the ultraviolet frequency range ($\hbar\omega_0 = 5$ eV), and gives an electric field strength E_0 between one and five atomic units e/a_0^2 ($a_0 = .529\ 10^{-8}$ cm). This corresponds to an electromagnetic intensity between 7.10^{16} and 2.10^{18} W/cm. These characteristics are similar to those envisaged by Boyer and Rhodes [6]. With these parameters, the amplitude of the sinusoidal motion of the electrons is greater than:

$$\gamma = \frac{eE_0}{m\omega_0^2} \sim 30\ a_0 \tag{15}$$

The numerical application we have performed applies to the ^{235}U nucleus ($Z = 92$, $R = .14\ 10^{13} a_0$) and the $B(E\lambda)$ in eq. (12) has been taken to be one Weisskopf unit (we assume the nucleus is spherical with radius P_N):

$$B(E\lambda) = \frac{e^2}{4\pi} \left(\frac{3}{\lambda + 3}\right)^2 (R_N)^{2\lambda} \tag{16}$$

The one-electron transition probability (12) has been computed using the method described in the Appendix of Ref. 14. The dependence of $P(T_0, r_0)$ on $r_{\perp 0}$ is displayed in the Figures (1a) to (1c) of Ref. 14 for some sets of the other parameters. *The observed variations appear very well reproduced* by the formula (A-22) and (A-25) of Ref. 14 derived in the part 2 of the Appendix of Ref. 14. They can be rewritten here:

$$P(T_0, r_{\perp 0}, Z_0) \leq \begin{cases} P(T_0, r_{\perp 0} = R_N, Z_0) & r_{\perp 0} \leq R_N \\ P(T_0, r_{\perp 0} = R_N, Z_0) \times \left(\frac{R_N}{r_{\perp 0}}\right)^{2\lambda} & R_N \leq r_{\perp 0} \leq R_{max} \end{cases} \tag{17}$$

where $R_{max} \sim 5\ a_0$ is much larger than the atom radius. In all the calculations we have done, the prescription of (17) is very accurately followed.

Since $P(T_0, r_{\perp 0}, z_0)$ is *maximum* for $r_{\perp 0} = R_N$, *we now* concentrate on the behaviour of $P(T_0, r_{\perp 0} = R_N, z_0)$ as a function of z_0 and of the other parameters.

First, the dependence of this quantity on λ and on the laser flux ϕ is, in all the calculations, very closely reproduced by the eq. (A-28) derived in the part 2-b/ of the Appendix of Ref. 14. Namely:

$$P(T_0, r_{\perp 0} = R\ , Z_0) \quad \alpha \quad \frac{1}{\phi} \frac{1}{(2\lambda+1)^4\ (\lambda+3)^2} \tag{18}$$

The decrease of the transition probability when ϕ increases may seem surprising. This behaviour, however, is easily understood from the relation (15) between ϕ and the amplitude γ of the electron motion: when ϕ increases and the frequency remains fixed, the electrons move farther from the nucleus and the average Coulomb field experienced by the nucleus decreases. The dependence (18) on λ shows that the transition probability for $\lambda = 1$ is 12 times the one for $\lambda = 2$ and 67 times the one for $\lambda = 3$.

The figures (2a) to (2c) of Ref. 14 display the variation of $P(T_0, r_{\perp 0} = R_N, z_0)$ as a function of the initial position z_0 of the electron for different values of $\hbar\omega_0$ and $\hbar\omega_N$. One observes that the probability is quasi-periodic and becomes very small for some values of z_0. The period of the oscillations tends to be larger as $\hbar\omega_0$ and $\hbar\omega_N$ decrease. This behaviour can be understood by looking at the

formula (A-28) of the Appendix of Ref. 14: the probability is proportional to $(\cos(\alpha_0\omega_N/\omega_0))^2$ with $z_0 - \gamma\cos\sigma_0 = 0$. When $z_0 \ll \gamma$, $\alpha_0 \sim \pi/2 - z_0/\gamma$ and the period with respect to z_0 of this term approximately is

$$\pi\gamma \frac{\omega_0}{\omega_N} = \pi \frac{e}{m} \frac{E_0}{\omega_0\omega_N}$$

We now examine the one-electron transition probability (13). In the part 3 of the Appendix of Ref. 14, we show that an upper limit for this probability is given by (see (A-45)) of Ref 14:

$$p^{(i)}(T_0) = p^{max}(T_0, r_{\perp 0} = R_N) \times (D^{n\ell j} + E^{n\ell j}) \qquad (19)$$

with

$$p^{max}(T_0, r_{\perp 0} = R_N) = \max_{|z_0| \leq R_{max}} P(T_0, r_{\perp 0} = R_N, z_0) \qquad (20)$$

Here the electron i is supposed to belong to the (nlj) shell. The multipolarity-dependent coefficients k_λ are given by (A-43): $k_1 = 3.5$, $k_2 = 1.3$ and $k_\lambda = 1.0$ for $\lambda \geq 3$. $D^{n\ell j}$ and $E^{n\ell j}$ can be found for ^{235}U in the Table I.

The values of the p^{Max} of eq. (20) are listed in the second column of the Table II for different sets of the parameters $\hbar\omega_N$, $\hbar\omega_0$ and λ. We have assumed $\phi = 7.10^{16}$W/cm^2 for the laser intensity since the probability corresponding to other values of ϕ can be obtained exactly from eq. (18). One can verify that the behaviour of (20) with λ and $\hbar\omega_0$ follows - with an accuracy of about 8-9 % - the formula (A-31) of the Appendix of Ref. 14. Using dimensionless quantities and allowing $B(E\lambda)$ to be different from one Weisskopf unit, this formula can be rewritten here:

$$p^{Max}(T_0, r_{\perp 0} = R_N) = \frac{(\hbar\omega_0/5eV)^2}{\phi/\phi_0} \frac{B(E\lambda)}{(2\lambda+1)^4 (\lambda+3)^2} C(\hbar\omega_N) \qquad (21)$$

where $\phi_0 = 7.10^{16}$ W/cm^2 and $B(E\lambda)$ is expressed in Weisskopf units. Let us recall that this expression can be used only if the amplitude $\gamma = e E_0/m\omega^2$ is much larger than the atom radius. Using Table II of Ref. 14 it appears that $C(\hbar\omega_N)$ is a very slowly changing function of $\hbar\omega_N$. It takes values between 40 and 60 for 20 eV < $\hbar\omega_N$ < 500 keV. This result may appear surprising in view of the form (9) of $S(T_0, r_0)$: one expects that the exponential factor averages the integral to zero when ω_N is large. In fact, the expressions developed in the

Appendix reveal that only the values of t belonging to a very small interval Δt give a significant contribution to the integral. The discussion leading to the formula (A29) shows that

$$\Delta t \sim R_N/(\gamma\omega_0) = \frac{R_N}{2\pi\gamma} T_0 \approx 10^{-6} T_0 .$$

It follows that, unless $\hbar\omega_N$ larger than at least 1 MeV, the exponential in (9) does not oscillate in the interval Δt and consequently the integral is not expected to go to zero. From a physical point of view, this result means that the electrons are actually effective in exciting the nucleus only during the very short time they spend in the vicinity of the nucleus, that is, when the Coulomb field experienced by the nucleus is maximum. As a matter of fact, Δt is approximately the time required by a electron moving according to (1) with $r_{10} = 0$ to pass through the nucleus. However, this independence of the laser frequency from the nuclear transition is misleading. The number of cycles in a laser pulse, assumed to be 300 cycles in our case, restricts the nuclear transition to an integer multiple of the laser frequency. Practical considerations will restrict this multiple to an integer low enough to manage this energy matching.

The values listed in Table I of Ref. 14 show that the one-electron probability (19) is maximum for electrons belonging to the 1s and 2s shells, whatever the value of λ. The probability induced by one electron of the 3s1/2, 2p1/2 and 4s1/2 shells are 4%, 2% and 1% respectively the one induced by a 1s electron. Electrons on the other shells contribute only less than 1%. This remark has important consequences concerning the effect of possible coherence of the motion of the electron cloud, that is the relation between the total transition probability (11) and the one-electron transition probabilities (13). In fact, the summation over i in (11) can be restricted to the 4 electrons belonging to the 1s and 2s shells. Consequently:

$$P(T_0) = 4 \sum_i P^{(i)}(T_0) \tag{22}$$

with $P(T_0)$ given by (13) and (19). The effect of the coherence of the electron motion is therefore to multiply the incoherent transition probability by a factor of 4. This is much less than the Z factor that would have appeared by directly trying to majorate the square of the sum over i in (11).

Finally, the third column of Table II of Ref. 14 displays what we obtain as an upper limit for the incoherent transition probability induced by all the atomic electrons. It has been obtained by summing (19) over all the atomic shells. According to (22), the total coherent transition probability is about 4 times larger.

At this point, we would like to stress that the inner shells - whose contribution to the total probability is maximum - are occupied by the most tightly bound electrons. As shown in Table I of Ref. 14, the binding energies of the 1s, 2s and 2p electrons are larger than 15 keV. Since the kinetic energy of a free electron is, with the laser parameters we have adopted, at

most 4 keV, it is very unlikely that the electrons belonging to these shells can be set into motion by the laser. In addition, the electric field experienced by the innermost electrons certainly is considerably screened by the electrons belonging to the other shells. Table I of Ref. 14 also shows that the electrons belonging to the 3s, 3p and 3d shells have binding energies in the range 3-5 keV. The influence of the laser on these electrons may therefore be questionable. However, as we want to keep the possibility of dealing with high laser fluxes and, furthermore, are interested in deriving upper limits to the nuclear transition probabilities, we shall assume that the electrons of these latter shells do contribute to the nuclear excitation. Due to this remark, one should therefore be aware that the probabilities obtained when including all the Uranium electrons are very strongly overestimated. For this reason, we have listed in Table II of Ref. 14 the transition probabilities computed: a/ from the two 3s electrons which are the most bound ones that may contribute (fourth column), b/ when excluding the 1s and 2s electrons (fifth column) and c/ when excluding the 1s, 2s and 2p electrons (sixth column). The latter values certainly are the closest to reality and will be used in the following applications. Inclusion of coherence effects would, as before, result in a factor of 4 since, in this case, the 3s and 4s shells give the major part of the transition probability.

4. APPLICATION TO THE POPULATION OF THE 75 eV ISOMER OF ^{235}U.

Given the nuclear transition probabilities per laser cycle found in the preceding Section, it is interesting to estimate what would be the number of excited nuclei actually produced in a laser-induced excitation experiment. We choose, as before, to excite the 26 minutes, 75 eV isomer of ^{235}U, and assume this can be done by means of the following experimental set-up. A sample of ^{235}U hexafluoride gas is enclosed in a small container sitting in vacuum. The laser beam passes through the container by means of windows and is focused in the center of the container. For reasons related to the technology of highly transparent windows, the gas pressure inside the container is assumed to be only 1 Torr (1/760 atmosphere) at normal temperature [11]. In this situation, all the nuclei contained in the laser focus volume can be supposed to experience the full laser flux and energy [12]. The laser characteristics we choose are those considered previously: photon energy $\hbar\omega_0$ = 5 eV, flux $\phi = 7.10^{16}$ to 2.10^{18} W/cm^2, pulse duration = 300 laser cycles (~ 2.10^{13} s.), pulse repetition rate: one per second, focus volume (10 μm)^3. With the latter figure one finds that the number of ^{235}U nuclei involved in the excitation process is ~ 3. 10^7. We assume the experiment can be carried out during the whole isomer lifetime without change in the number of U nuclei contained in the focus volume.

Table I displays upper limits for the number of isomers that could be produced in such an experiment. The contributions of different electrons shells are displayed. In each case, we have listed the upper limit of the nuclear transition probability per laser cycle (taken from the line $\hbar\omega_N$ = 75 eV, $\hbar\omega_0$ = 5 eV, λ = 3 of Table II of Ref. 14) and the isomer yields corresponding to different exposure times. We first give the number of isomers produced by the motion of the two electrons of the 3s1/2 shell since, as was discussed previously, they are the deepest electrons that are likely to

be set into motion by the laser (their binding energy is comparable to the kinetic energy of a free electron in the electric field of a laser of flux 10^{18} W/cm). As expected, their contribution is fairly large: it represents 60 % that of the 80 less bound electrons. The last two lines include the contribution of the 2p, 2s and 1s shells. Although these shells will not participate in the nuclear excitation, we have displayed the corresponding isomer yields to show what would be the results if they were incorrectly taken into account. The maximum number of isomers that could be produced in the envisaged experiment during the level lifetime does not exceed 50-60,000. One must recall (end of Section II) that the nuclear transition must be an integer multiple of the laser frequency. Whether or not this yield would support a viable experimental measurement, we leave to the expertise of our more practical colleagues.

5. APPLICATION TO THE TRANSFER FROM THE STORAGE LEVEL TO THE LASING STATE.

As a second application, we now study if the excitation mechanism analyzed in this work can be used to induce the lasing phase of a gamma-ray laser, that is, to cause the transfer from the storage level to a short-lived state lying higher in energy. Our discussion will be based on the nuclear transition probability per laser pulse $P(T_p)$ deduced from (19). Taking as before the pulse duration to be 300 cycles, we have:

$$P(T_p) \simeq 300\ P^{Max}(T_o, r_{\perp o} = R_N) \sum_{n\ell j} g^{n\ell j}\ (D^{n\ell j} + K_\lambda E^{n\ell j}) \qquad (23)$$

where the summation over (nlj) extends to the electron shells participating to the nuclear excitation. We first assume the laser photon energy is $\hbar\omega_o \simeq 5$ eV, the laser flux is $\phi = \phi_o = 7.\ 10^{16}$ W/cm^2 and all electrons except those of the 1s, 2s and 2ν shells participate to the excitation. The quantity $P^{Max}(T_o,\ r_{\perp o} = R_N)$ is then very well reproduced by eq. (21). The transfer probabilities obtained from (23) in this case are listed in Table II for $B(E\lambda) = 1$ Weisskopf unit and different values of the multipolarity λ of the nuclear transition. Let us recall that the results are almost independent of the energy of the transition between 20 eV and 500 keV although the experimental problem of matching the nuclear transition energy to an integral multiple of the laser frequency will limit the magnitude of this ratio (we have taken $C(\hbar\omega_N = 50$ in (21)). It appears that in order to obtain a transfer probability per pulse larger than 10^{-4} - a value that seems at least desirable - with the above laser characteristics, one has to choose a nuclear transition of multipolarity $\lambda = 1$ and preferentially of collective nature ($B(E\lambda) \geq 1$).

It must be noted that acting on the parameters $\hbar\omega_o$ and ϕ of the laser would not help much in improving the transfer probabilities. At first sight, it seems that increasing the ratio $(\hbar\omega_o)/\phi$ contained in $P^{Max}(T_o, r_{\perp o} R)$ would be useful. However, since the average kinetic energy acquired by one electron from the laser is inversely proportional to this ratio, less electron shells are expected to participate to the nuclear excitation and the sum in (23) tends to decrease. For instance, if the above ratio is multiplied by 10, only the 5s and less bound shells should, in the best case, be included and the probability $P(T_p)$ increases by a factor of at

most 5. Higher values of the ratio would result in no electron motion at all. We would stress, however, that the case of low flux and high $\hbar\omega_0$ is somewhat beyond the limits of the model used in this work which neglects the influence of the central Coulomb potential on the motion of the electrons.

6.. CONCLUSIONS.

A realistic but computationally simple model has been employed to examine nuclear transitions induced by laser-electron coupling. Applications to both test the model and to design criteria for a gamma-ray laser have been made. Nearly independent of laser characteristics we find the percentage of nuclear transitions of multipolarity λ to be

$$\approx 10^{-(3+\lambda)} \quad , \quad \lambda = 1,2,3$$

for 1 Weisskopf unit transition matrix element.

Most electromagnetic transitions between nuclear states have strengths several orders of magnitude less than 1 Weisskopf unit. In particular, one can expect the transition between a spin isomeric state and a proximate lasing state to be severely inhibited. Therefore, these results emphasize the importance of finding isomeric states of a collective nature.[13] They also indicate the value of a quantum mechanical calculation of this motion such as that of ref. (4) which would include self shielding and a more precise treatment of electron currents near the nucleus.

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ACKNOWLEDGMENTS

The authors would like to thank D. Iracane and R. Dei-Cas for many discussions and helpful comments. Two of us (J.-F. B. and D. G.) wish to acknowledge the warm hospitality received in the Nuclear Theory Group at LLNL, where most of this work has been done. Lastly, one of us (MSW) would like to acknowledge helpful discussions with A. Szöke and C. Rhodes.

This work has been performed under the auspices of the Department of Energy by the LLNL contract number W-7405-ENG-48.

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Table I.

^{235}U ELECTRONS TAKEN INTO ACCOUNT	NUCLEAR TRANSITION PROBABILITY PER CYCLE	NUMBER OF ISOMERS PRODUCED		
		during one laser cycle	during one pulse	in isomer lifetime
two 3s 1/2	1.3 10^{-9}	.04	11.7	~ 18,300
All except 1s,2s,2p	3.4 10^{-9}	.10	30.6	~ 47,700
All except 1s,2s	4.2 10^{-9}	.13	37.8	~ 59,000
All	4.0 10^{-8}	1.2	360.	~ 560,000

Table II.

TRANSITION MULTIPOLARITY λ	MAXIMUM TRANSFER PROBABILITY PER LASER PULSE ($B(E\lambda)$ = 1 Weisskopf)	
1	1.7	10^{-4}
2	7.5	10^{-6}
3	.13	10^{-6}