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TITLE NUCLEAR INTERLEVEL TRANSFER DRIVEN BY COLLECTIVE OUTER SHELL ELECTRON OSCILLATIONS

AUTHOR(S) G. A. Rinker, J. C. Solem, L. C. Biedenharn

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Nuclear Interlevel Transfer Driven by
Collective Outer Shell Electron Oscillations

G. A. Rinker and J. C. Solem
Theoretical Division, Los Alamos National Laboratory
Los Alamos, New Mexico 87545

and

L. C. Biedenharn
Duke University
Durham, NC 27706

ABSTRACT

We discuss the general problem of dynamic electron-nucleus coupling, and the possibility of using this mechanism to initiate gamma-ray lasing. Single-particle and collective mechanisms are considered. The problems associated with accurate calculation of these processes are discussed, and some numerical results are given. Work in progress is described.

1. The problem of the transfer process

A commonly-proposed gamma-ray laser scheme¹ is shown in Fig.1. One envisions a long-lived storage state I which can be populated by some laser² or radiochemical³ means, and which can be pumped with a relatively small amount of energy to a lasing state I'. This state can then decay with one or more radiative transitions, including at least one with the desired lasing characteristics. The graph represents a simple, ideal situation. The storage state shown (2^-) requires a magnetic quadrupole transition to reach the ground state (0^+) directly. This is strongly inhibited. The transfer step $2^- \rightarrow 1^+$ is electric dipole, normally the strongest multipole. The lasing transition $1^+ \rightarrow 0^+$ is magnetic dipole.

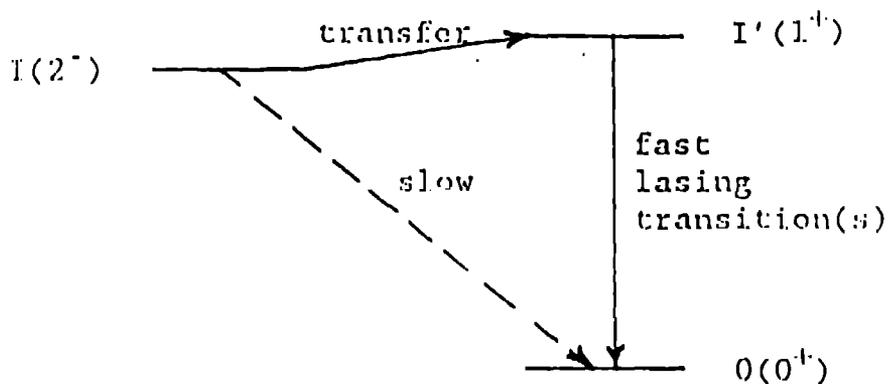


Fig. 1. Proposed gamma-ray laser level scheme.

avored for Borrmann⁴ mode propagation in crystals. The physical problem is to transfer the population I-I' efficiently.

1.1 Direct nuclear photoabsorption

One of the simplest transfer procedures would be to use direct nuclear photoabsorption $\gamma I \rightarrow I'$. This rate can be estimated on dimensional grounds. The photon field is

$$\vec{E} = E_0 \hat{e} \sum_{L=0}^{\infty} i^L \frac{(kr)^L}{(2L-1)!!} P_L(\cos\theta) \quad (1)$$

The interaction matrix element is thus

$$M_L = E_0 (2\pi R_N/\lambda)^L = E_0 (2 \times 10^{-7})^L \quad (2)$$

This expression contains the dimensionless parameter (nuclear radius)/(photon wavelength), which is a small number. ***** Thus the transfer rate is greatly inhibited compared with other competing atomic photoabsorption processes.

1.2 Dual process

An alternate process is depicted in Fig.2. Here, photoabsorption occurs on the atomic electrons, and the residual electron-nucleus interaction is used to transfer this excitation to the nucleus. We call

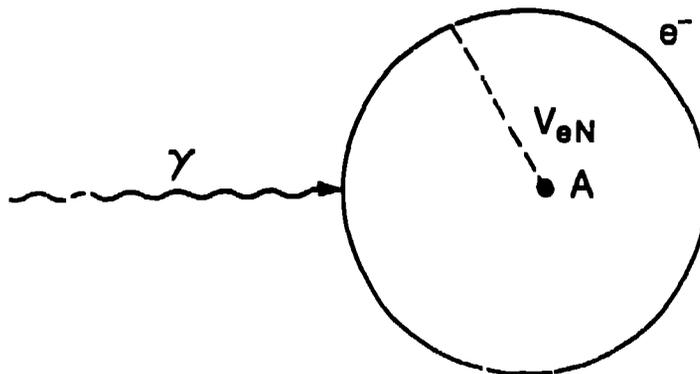


Fig. 2. Dual process for nuclear excitation.

this a dual process because these two interactions take place simultaneously. It has been called a two-step process, but this is not entirely appropriate because there is no well-defined, time-resolved intermediate state.

The photoabsorption matrix element has the same form as Eq.(2), but the nuclear radius is replaced by a length of atomic electron dimensions:

$$M_L = E_0 (2\pi r_e / \lambda)^L \approx E_0 (3 \times 10^{-3})^L \quad (3)$$

This is larger than before by a factor $\approx (1.5 \times 10^4)^L$. However, we must also include the electron-nucleus interaction

$$V_{eN} \approx e^2 \sum_{L=0}^{\infty} \frac{r_{<}^L}{r_{>}^{L+1}} P_L(\cos\theta) + e^2 \sum_{L=0}^{\infty} R_N^L r_e^{-L-1} P_L(\cos\theta) \quad (4)$$

The total effective matrix element is thus

$$M_L \rightarrow E_0 (2\pi r_e / \lambda)^L \cdot \frac{1}{\Delta E} \cdot e^2 R_N^L \langle j | r_e^{-L-1} | j' \rangle \quad (5)$$

There are three important factors. The first is just the electron photoabsorption matrix element. The last is the electron-nucleus interaction taken between electron states j and j' , which may be single-particle or many-body states. This is not necessarily small, since the interaction diverges for small r . Angular momentum selection rules prevent the occurrence of divergent matrix elements, but the fact that this interaction is more singular than the plane wave is the reason that internal conversion dominates radiative decay for a wide class of nuclear transitions. The middle term in Eq.(5) is the energy denominator (propagator) which occurs in any second-order process. Note that a simple dimensional argument gives:

$$\begin{aligned} M_L &\sim E_0 (2\pi r_e / \lambda)^L (e^2 / \Delta E) (R_N^L / r_e^{L+1}) \\ &\sim E_0 (2\pi R_N / \lambda)^L \cdot \frac{1}{\Delta E} \cdot \frac{e^2}{r_e} \quad (6) \end{aligned}$$

The last factor is approximately one atomic energy unit. If this is comparable with ΔE , the expression reduces to precisely what we had before for direct nuclear photoabsorption. Thus we can obtain enhancement only by making $\Delta E \rightarrow 0$, or by exploiting the fact that in general,

$$\langle j | r_e^{-L-1} | j' \rangle \approx \langle j | r_e | j' \rangle^{-L-1} \quad (7)$$

Thus any amplification by the atomic electrons depends upon the exact electron-nucleus matrix elements, which may involve collective electron effects in addition to detailed behavior of the electron wave functions near the nucleus. This is known as the dynamic hyperfine effect in muonic atoms, where it has been studied for many years.⁵ It has more recently been applied to electronic atoms by Morita.⁶

2. Quantitative description of electron-nucleus coupling

Figure 3 shows a detailed level diagram for the dual process. The uncoupled nucleus is on the left, with the lasing decay represented by the rate $\Gamma_{I'0}$. Relevant electron states are in the middle, with the

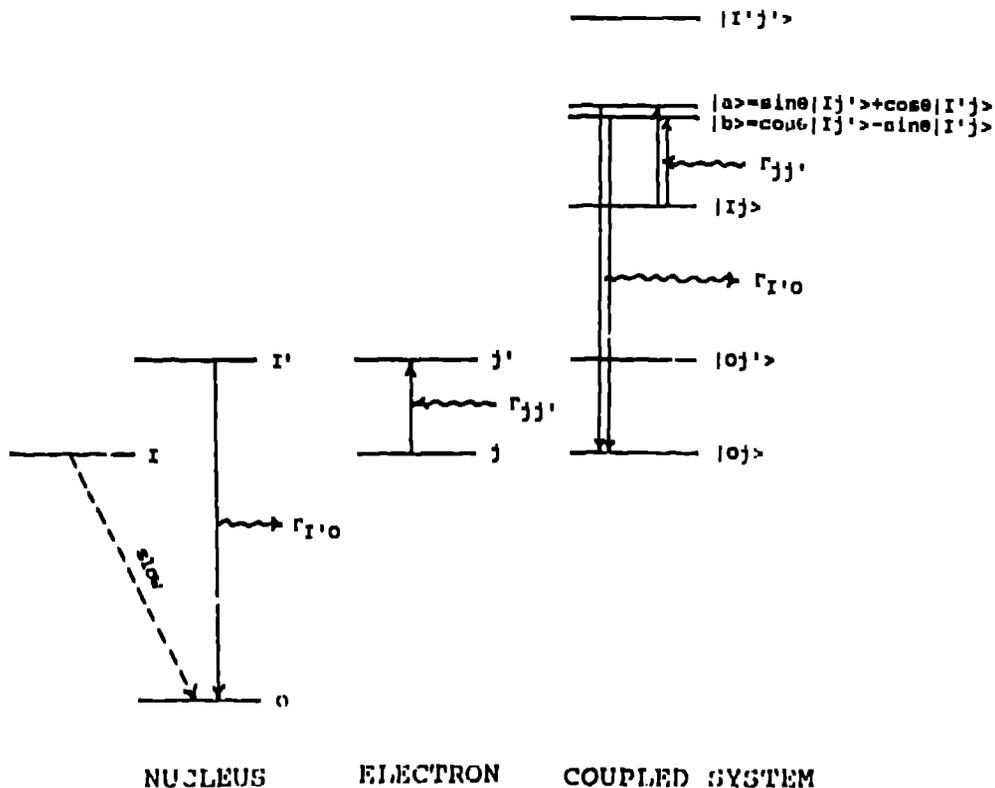


Fig. 3. Level diagram for the dual process.

photoabsorption rate represented by Γ_{jj} . The states of the coupled system are shown on the right. It is assumed that the initial state of the coupled system is $|Ij\rangle$. This state is pumped via the electron components with the rate Γ_{jj} , to the mixed states $|a\rangle$ and $|b\rangle$, which then decay via their nuclear components to the ground state $|0j\rangle$. The hamiltonian eigenstates $|a\rangle$ and $|b\rangle$ are described in terms of their mixing angle θ in the nearly-degenerate subspace $|2\rangle, |3\rangle$:

$$\begin{aligned} |a\rangle &= \sin\theta|2\rangle + \cos\theta|3\rangle & |2\rangle &= |Ij'\rangle \\ |b\rangle &= \cos\theta|2\rangle - \sin\theta|3\rangle & |3\rangle &= |I'j\rangle \end{aligned} \quad (8)$$

The hamiltonian in this subspace is

$$H = \begin{pmatrix} E_2 & V_{23} \\ V_{23} & E_3 \end{pmatrix} \quad \begin{aligned} E_2 &= E_I + E_j, \\ E_3 &= E_{I'} + E_j \\ V_{23} &= \langle 2|V_{eN}|3\rangle \end{aligned} \quad (9)$$

The solution for θ is exactly

$$\tan\theta = \frac{E_2 - E_3}{2V_{23}} \left\{ 1 \pm \left[1 + \frac{4V_{23}^2}{(E_2 - E_3)^2} \right]^{1/2} \right\} \quad (10)$$

Note that

$$\tan\theta \rightarrow \pm 1 \quad \text{as } E_2 \rightarrow E_3, \text{ and}$$

$$\tan\theta \rightarrow \theta \approx \frac{V_{23}}{E_2 - E_3} \quad \text{as } V_{23} \rightarrow 0 \quad (11)$$

Finally, the steady-state population N' of the lasing state I' is

$$N' = \frac{2 \tan^2 \theta}{1 + \tan^4 \theta} \frac{\Gamma_{jj'}}{\Gamma_{I'O}} \rightarrow \frac{\Gamma_{jj'}}{\Gamma_{I'O}}, \quad E_2 \rightarrow E_3$$

$$\rightarrow 2 \frac{V_{23}^2}{(E_2 - E_3)^2} \frac{\Gamma_{jj'}}{\Gamma_{I'O}}, \quad V_{23} \rightarrow 0$$

= mixing coefficient $\cdot \frac{\text{electron photoabsorption}}{\text{nuclear decay}}$. (12)

Note that in the limit of complete mixing ($\tan^2 \theta \rightarrow 1$), the steady-state population is precisely what one would obtain in balance between photoabsorption and decay, except that the photoabsorption occurs with the electron rate instead of the nuclear rate. For less complete mixing, the population is reduced by a factor $2\theta^2$, where $\theta = -V_{23}/(E_2 - E_3)$. We define the amplification factor

$$K = \frac{2 \tan^2 \theta}{1 + \tan^4 \theta} \frac{\Gamma_{jj'}}{\Gamma_{II'}} \quad (13)$$

This is the factor by which the population inversion is increased over what it would be by direct nuclear photoabsorption. This factor varies between 0, for small mixing, and a maximum value of order $(r_e/R_N)^{2L}$. The naive discussion of Eq.(6) given previously leads to an expected value of order unity.

2.1 Single-particle electron response

In order to calculate V_{23} for single-particle transitions, we have made use of electron wave functions from atomic Dirac-Hartree-Fock-Slater calculations. We have made a general canvass of L=1, 2, and 3 transitions for atomic numbers Z=12, 40, 68, and 92, using nuclear transition strengths of a few single-particle units. Results⁷ for L=1 and 2 are shown in Fig.4, where the matrix element V_{23} is plotted against electron transition energy ΔE . For L=1, the matrix elements all lie between 10^{-3} and 1 electron volt. For L=2, they lie between 10^{-6} and 10^{-3} eV. We found no case for L=3 where the matrix element was greater than 0.5×10^{-6} eV.

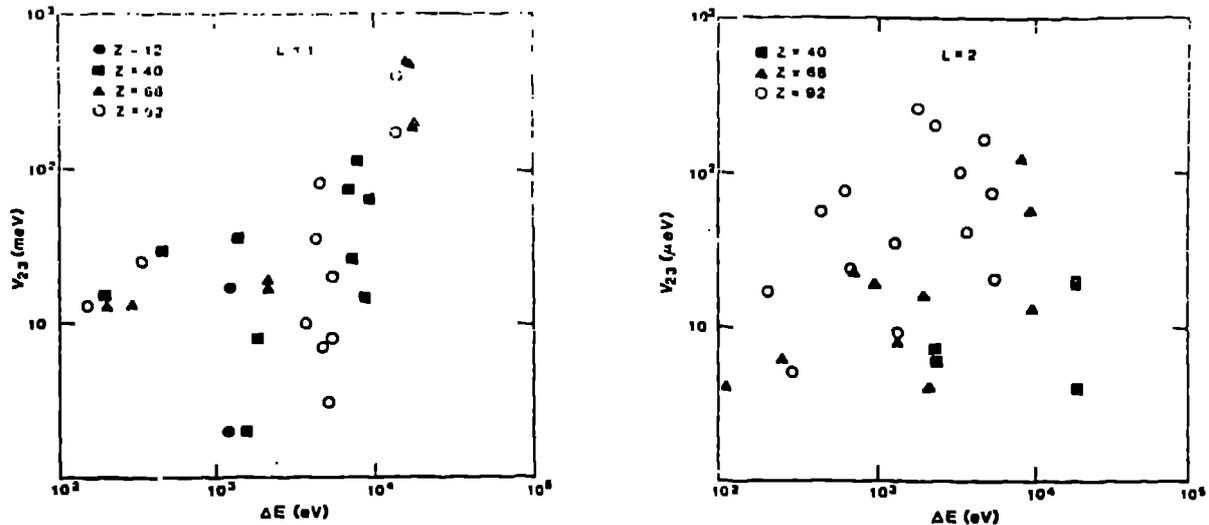


Fig. 4. Dipole and quadrupole matrix elements for selected electron transitions.

3. Collective electron response

3.1 The Articles of Faith

The possibility of collective electron motion modifies the above considerations. The formalism remains the same in principle, but the number of electron degrees of freedom increases enormously. In addition, there is a nonlinear interplay between the Coulomb interaction with the nucleus, the self energy of the electron gas, and the interaction with the external photon field. The nonlinear spatial and temporal interactions have led to three conjectures about an atom's response to the fields generated by a high-intensity laser. We facetiously refer to these conjectures as "The Articles of Faith:"

1. Atomic electrons can amplify the field produced at the nucleus;
2. The electrons can produce harmonics of the driving field, inducing nuclear transitions of energy greater than the quantum energy of the driving field; and
3. The electrons can generate electric fields at the nucleus of higher multipolarity than the driving field.

3.2 Calculation of collective interaction

Quantitative investigation of these questions on a quantum-mechanical basis involves many degrees of freedom. In most attempts made so far, basis truncation errors are severe. A classical approach reorganizes and

averages the degrees of freedom, lessening truncation effects but introducing physical errors. There exist some atomic problems for which a classical description has proved useful in obtaining a quantitative description.⁸ However, we must keep in mind that the electron-nucleus coupling involves discrete states and is inherently quantum-mechanical. For this aspect of the problem, classical approximations can yield only qualitative insight.

Our approach is a stepwise combination of classical and quantum methods. To calculate the interlevel transfer, we

(1) Use a classical model (Vlasov gas) to describe the collective electron response to an applied laser field;⁹

(2) Diagonalize a quantum electron hamiltonian¹⁰ using the applied photon field and the self-consistent electron density response (1); and

(3) Use the mixed quantum electron states to compute electron-nucleus matrix elements as for the single-particle states.

A fourth step, to integrate the time-dependent Dirac equations for the coupled system, is in progress. This step will be required in order to calculate nuclear pumping rates. Other statistical approaches (in particular, the Thomas-Fermi model and a self-interacting Vlasov gas) are being considered for step 1.

Figure 5 shows some typical results of step 1. Plotted are the L-1, 2, and 3 moments of the atomic electron density distribution calculated as a classical Vlasov gas, for an applied dipole laser field of one atomic unit in strength and wavelength $\lambda=198\text{nm}$. The moments are defined as

$$Q_L = \frac{(2L+3)(2L+1)}{4\pi} \frac{Z}{n} \sum_{i=1}^n (r_i/R)^L P_L(\cos\theta), \quad (14)$$

where n is the number of Vlasov test particles (10000), and R is an atomic dimension ($2a_0$) beyond which particles are considered lost. Initial ionization is apparent in the strong peak in the first laser cycle, after which the density settles down to a more regular response. The dipole density closely follows the applied field. The quadrupole and octupole responses, however, show some evidence of higher harmonics and are clearly nonzero; in spite of the fact that the applied field has no components beyond $L-1$. This supports the second and third articles of faith. The noise is somewhat reduced if the magnetic force due to the photon field $\vec{v} \times \vec{B} = \vec{v} \times (\vec{k} \times \vec{E})$ is included, but no significant change in the oscillation amplitudes or harmonics is produced.

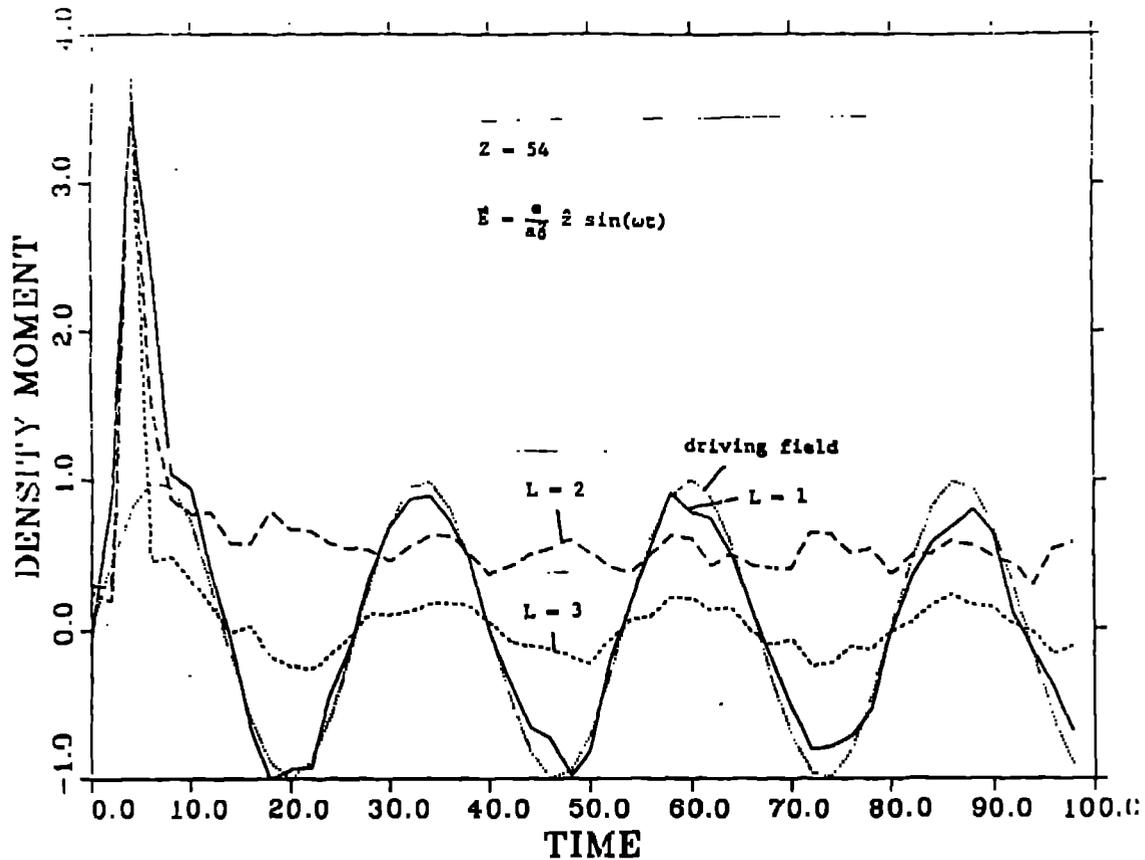


Fig. 5. Density moments for the Vlasov gas.

Figures 6 and 7 show the averaged electron density along the z-axis as a function of time, defined by

$$\rho(z) = \int_{-\infty}^{\infty} dx \int_{-\infty}^{\infty} dy \rho(x,y,z) \quad (15)$$

4. Coupling to the nucleus

The classical electron density distributions in Figs.5-7 could be used in principle to generate time-varying multipole potentials at the nucleus, in order to calculate excitation rates. Such a procedure would be quantitatively inaccurate, however, because the classical approximations fail at distances small compared with the electron Compton wavelength. Even the semiclassical Thomas-Fermi model gives the wrong density behavior near the nucleus. Our step (2) described above is an attempt to avoid this difficulty. These calculations are not complete, but we have discovered several interesting facts: (1) The intense electric field used for the classical results in Figs.5-7 also produces strong

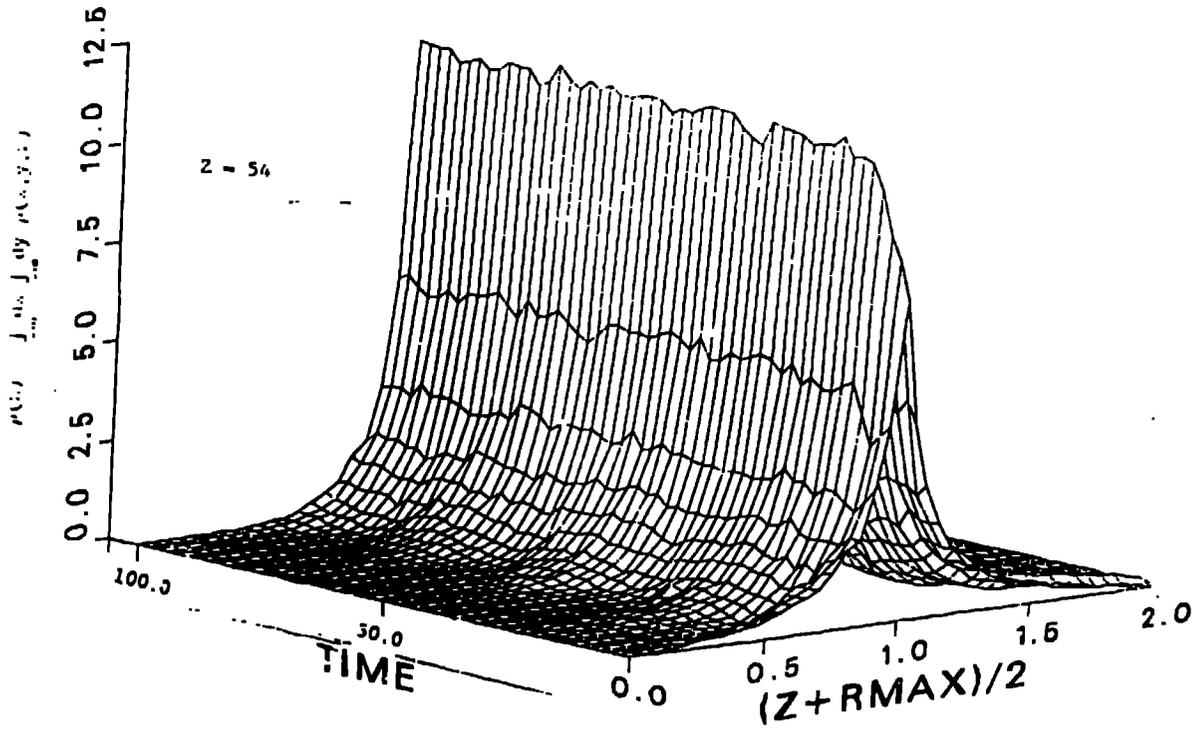


Fig. 6. Electron density along the z-axis as a function of time.

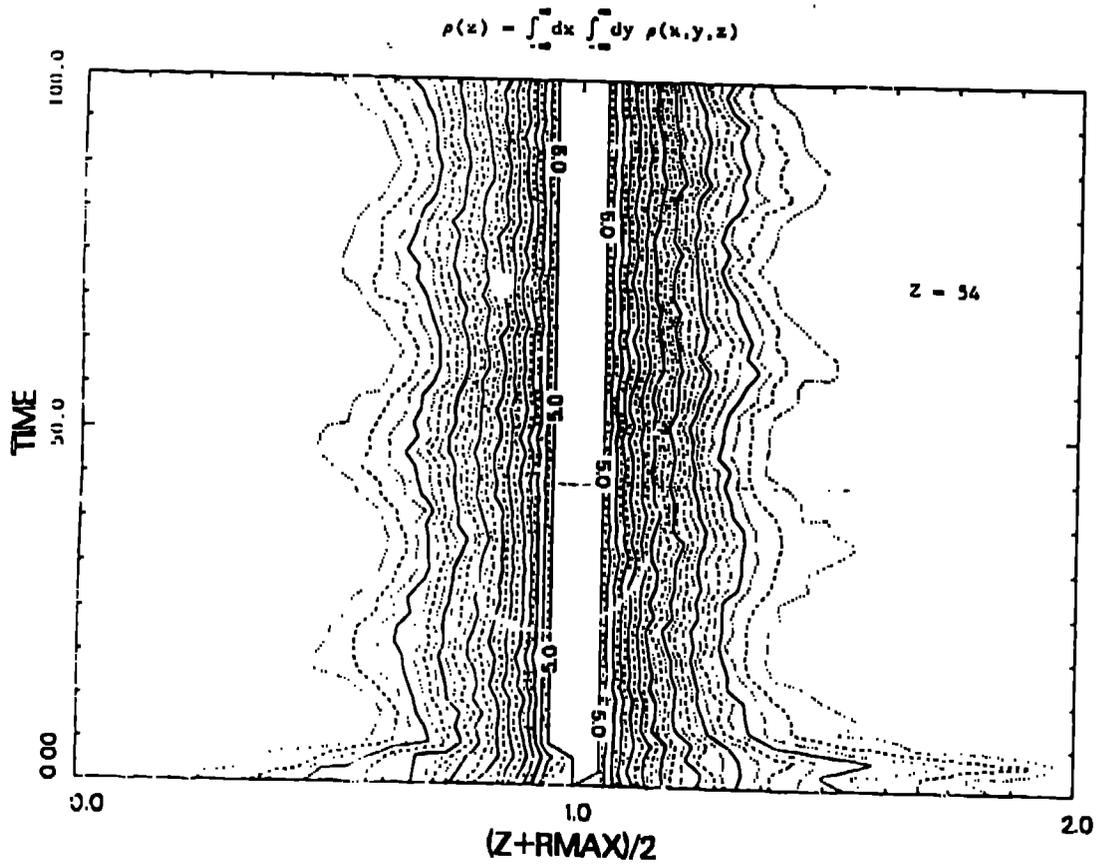


Fig. 7. Electron density along the z-axis as a function of time.

mixing among quantum electron states, including effective $\Delta L=2,3\dots$ couplings arising from higher orders in the interaction potential; (2) The self-consistent dipole interaction from the Vlasov gas is a shielding effect, reducing the effect of the applied field by $\approx 15\%$; (3) The quadrupole and octupole interactions from the Vlasov gas are smaller, but not zero; (4) Magnetic effects are negligibly small, as in the classical part of the calculation.

5. Conclusions and future work

These heuristic studies of nuclear interlevel transfer driven by the induction-field interaction with atomic electrons have led to the following tentative conclusions:

(1) $L=1$ pumping transitions are preferred. We find in general that higher multipole rates are down from this by factors of order $(10^3)^L$.

(2) The dual-process amplification factor is bounded from above by $\Gamma(\gamma j \rightarrow j') / \Gamma(\gamma I \rightarrow I')$ and is strongly dependent upon the electron states and degree of coupled-state degeneracy.

Our own program for future work includes increasing the basis for the quantum electron treatment, solving the time-dependent wave equations, and investigating the effects of self-interaction in the Vlasov gas.

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