A generalisation of the time-dependent Hartree-Fock-Bogoliubov microscopic equations of motion is realised by including the Landau-Zener effect for unpaired nucleons. The method is described in full detail. Apart the pairing residual interaction, a new term is included in the total Hamiltonian, that simulates the repulsion between levels characterised by the same good quantum numbers. The system for the microscopic equations of motions are obtained following the variational principle. The classical equations for the Landau-Zener effect and those for the TDHFB emerge as particular cases. The formalism is applied to describe the fine structure phenomenon in the cluster decay.

INTRODUCTION

Recently, a new method was proposed to approach the fine structure phenomenon of cluster- and alpha-decays in odd-nuclei. These decay modes are considered as supersymmetric fission processes. Assuming the existence of few collective variables, associated to some coordinates describing the nuclear shape parametrization, this approach allows to handle approximately the behaviour of many other intrinsic variables. For any values of the shape generalised coordinates, the single-particle level scheme and the potential barrier are determined. During the whole process and after the disintegration, the single particle occupation probabilities of the some orbitals of interest were computed using the only Landau-Zener promotion mechanism. Parameters with the meaning of spectroscopic factors were determined. This method was successfully used to describe the $^{14}$C- and alpha-decays [1–3]. Despite the overlook of residual interactions, the agreement obtained with data gives a strong experimental support to the formalism. Unfortunately, the approach based on the Landau-Zener effect alone is not able to describe the fine structure in the case of the disintegration of even-even nuclei. That motivated us to develop the model in order to introduce the pairing. A new set of coupled channel equations is obtained [4]. The way to obtain this system is described in detail.
The decaying system provides a time dependent single-particle potential in which the nucleons move independently. This description is consistent within the essence of the Time-Dependent-Hartree-Fock approximation where the many body wave function is constrained to be Slater determinants and with the Time-Dependent-Hartree-Fock-Bogoliubov (TDHFB) problem where the function is constrained to be of BCS form all time. In the HFB model, it is believed that the level slippage is treated automatically due to the simultaneous presence of many Slater determinants in the wave function. Such calculations are very cumbersome, especially for an odd number of neutrons or protons. Some simplifications are often introduced. For example, only pairing residual interactions are retained. Usually, the TDHFB equations are solved for even-even nuclei. A way to bypass this problem is to use the superfluid model constructed over a single particle potential and introducing the Landau-Zener effect to replace the effect given by a part of the residual interactions.

In the next section, the equations describing the microscopic dynamics of a many nucleon system with pairing residual interaction, including the Landau-Zener effect applied to the unpaired nucleon, are deduced. This effect allows the transition of an unpaired nucleon between levels with the same quantum numbers associated to some symmetries of the system. Solving the equations, the response of the system is estimated when the nuclear shape parametrization is changed. In the following, a nuclear shape parametrization given only by two intersected spheres, the radius of the sphere associated to the light fragment kept as constant, will be used. The single degree of freedom remains the elongation characterised by $R$, the distance between the centers of the nascent fragments.

**THE LANDAU-ZENER EFFECT**

The single-particle levels are function of the deformation parameters. Levels characterised by the same quantum numbers associated to some symmetry of the system cannot cross and exhibit avoided level crossings. The transition probability of one nucleon from one adiabatic level to another adiabatic level is strongly enhanced in an avoided crossing region. This promotion mechanism is known as the Landau-Zener effect [5–7]. In Fig. 1 the avoided crossing between two adiabatic levels $e_k$ and $e_l$ is displayed. The diabatic levels are $e_m$ and $e_j$. If the variation of the generalised coordinate is produced slowly and the nucleon is initially located on the level $e_l$, after the passage of the avoided crossing region, the nucleon will practically remain on the same level (panel (b) on the Fig. 1). In this case, the motion is adiabatic and the nucleon follows the adiabatic state $e_j$. If the variation of the generalised coordinate is produced with a high velocity, then the nucleon will skip with a great probability on the level $e_k$. The motion is diabatic, and the nucleon follows
the diabatic state $\epsilon_j$ (panel (a) on the Fig. 1). The two situations can be described by considering that the states are in general diabatic and an interaction between the diabatic states manage the probability that the nucleon remains on the adiabatic level.

Using quasiparticle creation and annihilation operators $\alpha^+_k$ and $\alpha_k$

$$\alpha^+_k = u_k a_k - v_k a^+_k, \quad \alpha_k = u_k a_k + v_k a^+_k$$

$$\alpha^*_k = u^*_k a^*_k - v^*_k a^*_k, \quad \alpha^*_k = u^*_k a^*_k + v^*_k a^*_k$$

it is possible to construct the operators that help us to promote the nucleon from one diabatic level to another. Here $a^+_k$ and $a_k$ denote operators for creating and destroying a particle in the state $k$, respectively. The state characterised by a bar signifies the time-reversed partner of a pair. The parameters $v_k$ and $u_k$ are the occupation and vacancy amplitudes, respectively. Because only the relative phase between the parameters $u_k$ and $v_k$ matters, in the following $u_k$ is considered as a real quantity and $v_k$ a complex one.

$$u_k = u^*_k, \quad v_k = -v^*_k.$$  

For diabatic states, the second case (b) can be obtained within an operator of the type:

$$O_{mj} = \alpha^+_m a_j + \alpha^+_j a_m.$$  

That means, the probability to remains on the same adiabatic level is managed by the operator $O_{mj}$ and, as in the classical Landau-Zener effect, by the interaction energy between the diabatic levels. The maximum value of the interaction energy is in the point of nearest approach between the adiabatic levels as plotted in Fig. 1(b).
THE VARIATIONAL METHOD

In order to obtain the equations of motion, we shall start from the variational principle taking the Lagrangian as:

\[ L = \left\langle \phi \left| H - i\hbar \frac{\partial}{\partial t} + \mathcal{H}' - \lambda N \right| \phi \right\rangle \]  \hspace{1cm} (1)

and assuming the many-body state formally expanded as a superposition of time dependent BCS seniority one diabatic wave functions

\[ |\phi\rangle = \sum_i c_i(t) a_i^+ \prod_{j \neq i} (u_j(t) + v_j(t) a_j^+ a_j) |0\rangle \]  \hspace{1cm} (2)

The Lagrangian contains several terms. The first one is the many body Hamiltonian with pairing residual interactions

\[ H = \sum_{k>0} c_k (a_k^+ a_k + a_k^2 a_k^2) - G \sum_{k,l>0} a_k^2 a_k^2 a_l a_l \]

The residual interactions between diabatic levels characterized by the same quantum numbers is

\[ H' = \frac{1}{2} \sum_{m,j} h_{mn}^n O_{mj} = \]

\[ = \frac{1}{2} \sum_{m,j} h_{mn}^n \left[ (u_m a_m - v_m a_m^2) (u_j a_j^+ - v_j a_j^2) + (u_j a_j - v_j a_j^2) (u_m a_m^2 - v_m a_m) \right] \]

The sum runs over \( n \), the total number of avoided crossings between adiabatic levels \( m \) and \( j \), for any pair \((mj)\). The number of particle operators is:

\[ N = \sum_{k>0} (a_k^+ a_k + a_k^2 a_k^2) \]

After some calculations we obtain:

\[ \left\langle \phi \left| H - i\hbar \frac{\partial}{\partial t} + \mathcal{H}' - \lambda N \right| \phi \right\rangle = \]

\[ = \sum_m |c_m|^2 \left[ \sum_{k>0} 2 v_k \left( \epsilon_k - \lambda \right) + \left( \epsilon_m - \lambda \right) - G \sum_{k>0} u_k v_k \left| c_m \right|^2 \right] - \]

\[ - i \hbar \sum_m |c_m|^2 \left[ \sum_{k>m, k>0} \frac{1}{2} (v_k^* \dot{v}_k - \ddot{v}_k v_k) \right] - \)

\[ - i \hbar \sum_{m>0} |c_m|^2 \dot{c}_m + \frac{1}{2} \sum_{m,n} h_{mn}^n \left( c_m^* c_j + c_j^* c_m \right) \]  \hspace{1cm} (3)
Where the terms like:

$$\sum_{m} |c_m|^2 \sum_{k} v_k^d$$

Are neglected. The independent parameters are $v_k$, $v_k^*$, $c_m$ and $c_m^*$ subjected to conditions:

$$u_k^2 + |v_k|^2 = 1$$

$$\sum_{m>0} |c_m|^2 = 1$$

In the following, the calculations are made in a way similar to that given in Ref. [8]. First, the equation (3) is derived with respect $v_l$ and $v_l^*$ and set to zero

$$\sum_{m>0} |c_m|^2 \left \{ 2v_l (\epsilon_l - \lambda) - G \left( \sum_{k \neq m}^{k>0} \sum_{k \neq m}^{k>0} \kappa_k \left( \frac{v_l v_l^*}{2u_l} \right) + \left( u_l - \frac{\rho_l}{2u_l} \right) \sum_{k \neq m}^{k>0} \kappa_k^* \right) + i\hbar v_l^* \right \} = 0$$

$$\sum_{m>0} |c_m|^2 \left \{ 2v_l (\epsilon_l - \lambda) - G \left( \sum_{k \neq m}^{k>0} \sum_{k \neq m}^{k>0} \kappa_k \left( \frac{v_l v_l^*}{2u_l} \right) + \left( u_l - \frac{\rho_l}{2u_l} \right) \sum_{k \neq m}^{k>0} \kappa_k^* \right) - i\hbar v_l \right \} = 0$$

(4)

Where:

$$\kappa_k = u_k v_k, \quad \rho_k = |v_k|^2.$$

Multiplying the first equation of the system (4) with $v_l$ and the second with $v_l^*$, and subtracting them, the next relation is obtained:

$$i\hbar \dot{v}_l = \sum_m |c_m|^2 \left \{ \kappa_l \Delta_m^* - \kappa_l^* \Delta_m \right \}.$$

(5)

Where:

$$\Delta_m = G \sum_{k \neq m}^{k>0} \kappa_k; \quad \Delta_m^* = G \sum_{k \neq m}^{k>0} \kappa_k^*$$

are the pairing strength constants. To deduce the previous equation we made use of the condition of number of particle conservation:

$$\sum_{k \neq m}^{k>0} v_k^* v_k = N - 1$$

Another equation can be obtained:
\[
\frac{\partial \kappa_l}{\partial t} = -\frac{v_l}{2u_l} \rho_l + u_l \dot{v}_l = \\
-\frac{v_l}{2u_l} \sum_m c_m^l \left\{ \kappa_i \Delta_m^* - \kappa_i \Delta_m \right\} \\
+ \frac{\sum_m c_m^l \left[ 2v_l (\varepsilon_l - \lambda) - \Delta_m^* \left( -\frac{v_l v_l}{2u_l} \right) \right] - \left( u_l - \frac{\rho_l}{2u_l} \right) \Delta_m}{ih} + u_l
\]

So that:

\[
ih \dot{\kappa}_j = \sum_m c_m^j \left\{ (2\rho_l - 1) \Delta_m + 2\kappa_i (\varepsilon_l - \lambda) \right\}
\] (6)

The equation (3) is now derived with respect to \(c_m\) and \(c_m^*\) and set to zero. The next relations are obtained:

\[
-i h \dot{c}_m^* = c_m^* \left[ 2 \sum_{k \neq m} \frac{1}{2} \left( \varepsilon_k - \lambda \right) + \left( \varepsilon_m - \lambda \right) - G \left( \sum_{k \neq m} u_k v_k \right) \right] - \\
-i h c_m^* \sum_{k \neq m} \frac{1}{2} \left( \varepsilon_k - \lambda \right) \right) + \frac{1}{2} \sum_{m} h_{mn}^p c_j^*
\]

\[
ni h \dot{c}_m = c_m \left[ 2 \sum_{k \neq m} \left( \varepsilon_k - \lambda \right) + \left( \varepsilon_m - \lambda \right) - G \left( \sum_{k \neq m} u_k v_k \right) \right] - \\
-i h c_m \sum_{k \neq m} \frac{1}{2} \left( \varepsilon_k - \lambda \right) \right) + \frac{1}{2} \sum_{m} h_{mn}^p c_j^*
\]

The first equation of the system (7) is multiplied with \(-c_m\) and the second with \(c_m^*\), and the two equations are added. The next equation is obtained:

\[
i (\dot{c}_m c_m^* + \dot{c}_m^* c_m) = \frac{1}{2} \sum_{mn} h_{mn}^p \left( c_j c_m^* - c_j^* c_m \right)
\] (8)

For a single passage in an avoided crossing region, only two parameters, for example \(c_m\) and \(c_j\) can change. On the other hand, from the condition of conservation can be obtained:

\[
\dot{c}_m c_m^* + c_m \dot{c}_m^* = - \left( \dot{c}_j c_j^* + c_j \dot{c}_j^* \right)
\]
This condition is fulfilled by the equation (8) above, so the equation conserves the norm.

From (7), by changing \( c_m \) with \( c_j \), the following system can be also written:

\[
\begin{align*}
-\hbar c_j^* c_m &= c_m^* c_j \\
&= \frac{\hbar^*}{2} \sum_{k \neq j} \left( v_k^* (\varepsilon_k - \lambda) + (\varepsilon_j - \lambda) - G \sum_{k > 0} u_k v_k \right) \\
&\quad - \frac{\hbar^*}{2} \sum_{k > 0} \left( v_k^* \dot{v}_k - \dot{v}_k v_k \right) + \frac{1}{2} \sum_{m,j,n} h_{m,j}^* c_m^* c_m \\
\end{align*}
\]

\[
\begin{align*}
\hbar \dot{c}_m c_j^* &= c_m^* c_j \\
&= \frac{\hbar^*}{2} \sum_{k \neq m} \left( v_k^* (\varepsilon_k - \lambda) + (\varepsilon_m - \lambda) - G \sum_{k > 0} u_k v_k \right) \\
&\quad - \frac{\hbar^*}{2} \sum_{k > 0} \left( v_k^* \dot{v}_k - \dot{v}_k v_k \right) + \frac{1}{2} \sum_{m,j,n} h_{m,j}^* c_m^* c_j \\
\end{align*}
\]

So that by subtracting we obtain:

\[
\begin{align*}
\hbar \left( \dot{c}_m c_j^* + \dot{c}_j c_m \right) &= \frac{1}{2} \sum_{m,j,n} h_{m,j}^* \left( |c_j|^2 - |c_m|^2 \right) + \\
&\quad + c_m^* c_j \left( 2 |v_j|^2 (\varepsilon_j - \lambda) + (\varepsilon_m - \lambda) - \Delta_m - 2 |v_m|^2 (\varepsilon_m - \lambda) - (\varepsilon_j - \lambda) - \Delta_j \right) + \\
&\quad + c_m^* c_j \left( \frac{\hbar^*}{2} (v_m^* \dot{v}_m - \dot{v}_m v_m - v_j^* \dot{v}_j + \dot{v}_j v_j) \right)
\end{align*}
\]

Entering in an avoided crossing, the set of parameters \((c_m, c_j)\) can have only the values \((0, 1)\) or \((1, 0)\) because the nucleus cannot have two unpaired nucleon in the same time. So, the last two terms in the right hand of the previous equation are set to zero in regions outside the avoided crossing region. Moreover, if the blocking effect is neglected, and we take into account that the energies of two adjacent levels \((m\) and \(j)\) are close and are practically equal in an avoided crossing region, the last equation becomes:

\[
\begin{align*}
\hbar \left( \dot{c}_m c_j^* + \dot{c}_j c_m \right) &= \frac{1}{2} \sum_{m,j,n} h_{m,j}^* \left( |c_j|^2 - |c_m|^2 \right) \\
&= \frac{1}{2} \sum_{m,j,n} h_{m,j}^* \left( |c_j|^2 - |c_m|^2 \right) \\
&\quad + c_m^* c_j \left( 2 |v_j|^2 (\varepsilon_j - \lambda) + (\varepsilon_m - \lambda) - \Delta_m - 2 |v_m|^2 (\varepsilon_m - \lambda) - (\varepsilon_j - \lambda) - \Delta_j \right) + \\
&\quad + c_m^* c_j \left( \frac{\hbar^*}{2} (v_m^* \dot{v}_m - \dot{v}_m v_m - v_j^* \dot{v}_j + \dot{v}_j v_j) \right)
\end{align*}
\]

Denoting

\[
P_m = |c_m|^2; \quad s_{jm} = c_j^* c_m.
\]

With the property...
We obtain:

\[ i\hbar p_m = \frac{1}{2} \sum_n h^n_{mj} (s_{mj} - s_{jm}) \]

\[ i\hbar s_{jm} = \frac{1}{2} \sum_n h^n_{mj} (p_j - p_m) \]

where \( p_m \) is a real quantity and \( s_{jm} \) an imaginary one.

The relations (5), (6) and (10) forms a coupled channel equations that represents a generalisation of the TDHFB ones. The unpaired nucleon slippage is taken into account in this new system. It will shown that the TDHFB equations and the classical equations that describe the Landau-Zener effect are particular cases of the our system.

For a two-level system the classical Landau-Zener equations can be written as:

\[ i\hbar \frac{d}{dt} \begin{pmatrix} c_1 \\ c_2 \end{pmatrix} = \begin{pmatrix} \varepsilon_1 & H \\ H & \varepsilon_2 \end{pmatrix} \begin{pmatrix} c_1 \\ c_2 \end{pmatrix} \]

After some simple calculations we find:

\[ i\hbar (\dot{c}_1 c_1^* + \dot{c}_2 c_2^*) = H c_1^* c_2 - H c_2^* c_1 \]

\[ i\hbar \dot{c}_1 c_2 = H c_1 c_2^* - H c_2 c_1^* \]

Using the notations for \( p \) and \( s \) we obtain a particular case of the system of equations (10), when the residual interactions are neglected.

When one level is occupied by an unpaired nucleon and no promotion mechanisms are allowed, only one level will be occupied all time by the nucleon. That is, \( c_m = 1 \) for a particular level while \( c_j = 0 \), for any \( j \). In this case, our system is reduced to the well known TDHFB equations described in Refs. [8, 9].

So the TDHFB equations and the classical equations for the Landau-Zener effect are particular cases of our system. Therefore the system represents a generalisation of the TDHFB equations.

**RESULTS FOR CLUSTER EMISSION**

The formalism deduced in the previous section is available for any mean field potential. In the following, the single particle level schemes are obtained with the supersasymmetric two-center shell model [10] improved in Ref. [1]. This model allows an estimation of the adiabatic levels during the disintegration, starting with one nucleus and ending with two separate fragments. This model is based on the asymmetric two center harmonic oscillator. The Hamiltonian
contains also corrections as the spin-orbit and orbital momentum \(L^2\) couplings and takes into account the depths of potentials of the nascent fragments. The single particle neutron levels are displayed in Fig. 2. The \(^{223}\)Ra has the spin 3/2 emerging from \(1i_{11/2}\). Adiabatically, this unpaired neutron reaches the \(2g_{9/2}\) level of the daughter \(^{209}\)Pb. In will be shown that, the fine structure of the \(^{14}\)C radioactivity can be understood by an enhanced transition probability of the unpaired neutron from the adiabatic level with \(\Omega = 3/2\) emerging from \(1i_{11/2}\), to the adiabatic level with the same spin projection emerging from \(1j_{15/2}\). The level scheme presented in Fig. 2 shows that the \(1i_{11/2}\) level reaches adiabatically the \(2g_{9/2}\) state, the \(1j_{15/2}\) reaches the \(1i_{11/2}\) state, and the \(3d_{5/2}\) level reaches \(1j_{15/2}\) state of the daughter \(^{209}\)Pb. In system with cylindrical symmetries, Landau-Zener transitions can be realized only between levels with the same spin projection \(\Omega\).

Four avoided level crossings can be observed in Fig. 2 produced at \((R - R_i)/(R_f - R_i) = 0.75, 1, 1.1\) and 1.2. Without these interactions, diabatically, the levels \(1i_{11/2}, 1j_{15/2}\) and \(3d_{5/2}\) attain the \(1i_{11/2}, 1j_{15/2}\) and \(2g_{9/2}\) daughter orbitals, respectively.

![Fig. 2 – Neutron level scheme for \(^{14}\)C spontaneous emission from \(^{223}\)Ra with respect the normalized elongation \((R - R_i)/(R_f - R_i)\). The levels (with \(\Omega = 3/2\)) emerging from \(1i_{11/2}, 1j_{15/2}\) and \(3d_{5/2}\) are represented with thick lines. \(R_f = R_1 + R_2\) for the configuration of the initial nucleus considered as a sphere and for that of two tangent spheres, respectively.](image-url)
In Refs. [8, 9, 11], a particular form of the equations (5) and (6) (without Landau-Zener included) were used to study the dumping (or the dissipation) process. The dissipation was defined as the flow of energy between collective and intrinsic modes. This energy flow must be irreversible. Unfortunately, according to the discussion of Ref. [9], neglecting the Landau-Zener effect, the calculation is microscopically reversible in the sense that if all the coordinates are time-reversed the system retraces its path. An unknown fraction of the energy that is identified as dissipated energy is in fact collective kinetic energy. Through the presence of the Landau-Zener term in the equations of motion the requirement of irreversibility is satisfied. Therefore, the Landau-Zener probabilities \( p_m \) describe a true dissipation process, while the variations of \( \rho_j \) and \( \kappa_k \) values take partially into account the collective kinetic energy. In these circumstances, it can be assumed that, after the disintegration, the spectroscopic factor can be approximated with only the leading term:

\[
\Phi_m = |\langle \Phi | \Phi_{0m} \rangle|^2 \propto p_m
\]

where \( \Phi \) is the BCS wave function obtained from the equations of motion after the scission, \( \Phi_{0m} = a_m \prod_j (u_{0j} + v_{0j}) a_j^\dagger  \ket{0} \), with \( m = 2g9/2, 1i11/2, 1j15/2 \), are the orthogonal set of seniority one wave functions constructed on the asymptotic overlapped energy levels of the daughter and the emitted nucleus with the ground state \( u_{0j} \) and \( v_{0j} \) values.

The partial half-life for the state \( i \) is obtained with the equation in the WKB approximation:

\[
T_{1/2}^m = \frac{\hbar \ln 2}{2E_i p_m} \exp \left( \frac{\hbar^2}{\hbar} \int A_1 A_2 E'' dR \right) = \frac{\hbar \ln 2}{2E_i p_m} \exp(K_m)
\]

where \( p_m \) is the final occupation probability and has the meaning of a spectroscopic factor. \( E'' \) is the deformation energy computed in the framework of the Yukawa-plus-exponential model extended for binary systems with different charge densities [12]. This deformation energy is corrected within shell effect computed within Strutinsky’s method as in Ref. [11]. The deformation energies for the excited states are obtained by adding the difference between the single particle energy of the excited state and the single particle energy of the nucleon in the ground state. The values of \( K_i \) are 2.10 \times 10^{27}, 1.27 \times 10^{29} and 1.87 \times 10^{31} for transitions to the daughter ground state, first excited state and second excited state, respectively. These deformation energies are plotted with thick line in Fig. 3. \( R_{g.s.} \) is the ground state elongation used as starting point of the
Fig. 3 – The potential barrier $E_d$ measured from the ground state of the parent as function of the normalized elongation coordinate. The barriers of the first and second excited states obtained by adding the single particle excitation energy are also plotted.

decay, $R_i$ is the channel exit point from the barrier and $E_v$ is the zero point vibration energy.

The effort is mainly focused to reproduce the experimental ratio $r$ of the partial half-lives for transitions to the first excited state and to ground state,

$$r = \frac{T^{11/2}_{g9/2}}{T^{11/2}_{10/2}} = \frac{p_{g9/2}}{p_{11/2}} \exp(K_{11/2})$$

which has approximately the value 0.218. A critical parameter in the model is the quantity which characterizes the variations in time of the generalized coordinates. In the present work, it is matter of the velocity $v$ of the inter-nuclear distance (or the elongation). Some models are able to determine this quantity [13], but in the present work the velocity of the elongation is considered as a fitting parameter and its optimal value will be briefly discussed in comparison with previous results. In Fig. 4, the final single-particle occupation probabilities as function of the internuclear distance velocity are plotted. As expected, for large velocities the system becomes mainly diabatic and the occupation probabilities of the first daughter excited state tends to 1 while that of the ground state reaches 0. For lower velocities, the g.s. transitions are favored. In the bottom of Fig. 4, the ratio $r$ is plotted. When $v \approx 9 \times 10^4 - 3 \times 10^5$ fm/fs, $r$ reaches the vicinity of the experimental value. The value of $v$ determined in this way is approximately one order of magnitude lower than that determined from the tunnelling times of Ref. [13] using macroscopic models, that means, approximating the effective mass with a value close to the reduced mass. However, in the region of avoided crossing levels, the more realistic cranking approximation
predicts that the inertia increase with orders of magnitude. If the effective mass
increases, the deformation velocity must decrease in order to conserve the
available energy. Therefore, the lower values of $v$ are qualitatively consistent
with the use of the cranking model for the inertia.

The previous treatment provides a different way to attack the fine structure
of supersymmetric fission for odd and even nuclei. To our knowledge, it is the
unique model which can explain the favoured transitions to the excited state in
the case of cluster-decay. Also, the inclusion of the Landau-Zener effect in the
equations of motion offers a more complete description of dumping phenomena
during the disintegration.

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