When the time dependent Hartree-Fock-Bogoliubov intrinsic equations of motion are solved in the case of seniority one nuclear systems, the unpaired nucleon remains on the same orbital. The blocking effect hinders the possibility to skip from one orbital to another. This unpleasant feature is by-passed with a new set of pairing time dependent equations that allows the possibility that the unpaired nucleon changes its single-particle level. These equations generalize the time dependent Hartree-Fock-Bogoliubov equations of motion by including the Landau-Zener effect. The derivation of these new equations is presented in details. These equations are applied in the case of a supersymmetric fission process, that is, in order to explain the fine structure the $^{14}$C emission from $^{233}$Ra. A new version of the Woods-Saxon model extended for two-center potentials is used in this context.

PACS numbers: 24.10.Eq Coupled channel and distorted wave model; 23.70.+j Heavy-particle decay

I. INTRODUCTION

In the Hartree-Fock (HF) approximation, the self-consistent potential for heavy nuclei is quite smooth, since it includes the convolution of many instantaneous densities. If the potential varies in time, each single-particle wave function moves independently in the smoothly varying field. The Pauli principle is fulfilled being mediated permanently through the mean field potential. The two-body collisions are incorporated in the equations of motion only to the extend to which they contribute to the mean field. In principle, the time dependent HF approach treats the residual interactions exactly only if the mean field is allowed to break all symmetries. A such approach leads to a very big computational problem. In order to avoid such numerical difficulties, usually, the HF mean field is constrained to be at least axial symmetric. In this case, levels characterized by the same good quantum numbers cannot intersect. When the equations of motion are solved in such circumstances, an individual single particle wave function will belong to only one orbital characterized by some good quantum numbers and the mechanism of level slippage is not allowed. In general, two dynamical approaches are used. On one hand, the generator coordinate method assumes that the internal structure of the decaying system is equilibrated at each step of the collective motion. On the second hand, the exchange between collective and internal degrees of freedom is neglected, so that adiabaticity is assumed. Usually, these behavior leads to the unpleasant feature that a system, even moving infinitely slowly, could not end up in its ground state. Some attempts were done to extend the time-dependent HF method in order to includes collision terms. This problem was partially solved by introducing a residual pairing interaction in the Time Dependent Hartree-Fock-Bogoliubov (TDHF) approach. This last method provides the possibility of level slippage for pairs and allows a description of the nuclear dynamics. For example, in the case of an even-even system, this method represents a tool to estimate the dissipation during disintegration processes. A connection with the Landau-Zener effect is included in the TDHF equations. Pairs undergo Landau-Zener transitions on virtual levels with coupling strengths given by the gap $\Delta$. Unfortunately, for seniority one nuclear systems the pairing residual interaction does not affect a single nucleon and, during the deformation of the nucleus from its initial state up to scission, the unpaired particle remains located on the same orbital. The level slippage is again forbidden for the blocked level.

In the case of independent single-particles, neglecting residual interactions, the problem of the unpaired nucleon is solved in terms of the Landau-Zener effect. The Landau-Zener effect reflects a mechanism that allows the possibility that a single nucleon skips from one single-particle level to another one in some avoided level crossing regions. The probabilities that the unpaired nucleon arrives in different final states can be computed by solving a system of coupled channel equations that characterizes the microscopic motion. In the following, a way to introduce a similar mechanism for the unpaired nucleon in superfluid systems is investigated. The TDHF equations will be generalized in order to include the Landau-Zener effect. The classical TDHF equations and the equations that govern the Landau-Zener effect will be obtained as particular cases of the new time dependent pairing equations.

II. LANDAU-ZENER EFFECT

The single-particle levels are function of the deformation parameters that characterize the shape of a nucleus. Levels characterized by the same quantum numbers associated to some symmetry of the system cannot cross and exhibit avoided level crossings. The transition probability of a nucleon from one adiabatic level to another one is strongly enhanced in an avoided crossing region. This promotion mechanism is known as the Landau-Zener effect. In Fig. (a) an ideal avoided crossing $(j, m)$ between
III. SUPERFLUID SYSTEMS

An effect analogous to the Landau-Zener one can be obtained by generalizing the TDHFB equations for the case of seniority one nuclear systems. The problem will be explored in the simplest possible way: a monopole pairing force, and a sufficiently weak pairing such that the nucleons are not redistributed to change significantly the mean field potential. In order to make the problem tractable, two approaches are investigated. The first one is valid for a low lying levels system with a small number of avoided crossing regions, so that variations of densities $\rho_l$ and pairing moment components $\kappa_i$ due to the blocked level can be neglected. The second one takes into account the blocking effect, that is, the fact that $\rho_l(m)$ and $\kappa_i(m)$ depend on the blocked level $m$.

A. Low lying levels

Using quasiparticle creation and annihilation operators $\alpha_k^+$ and $\alpha_k^-$

$$
\alpha_k = u_k a_k - v_k a_k^+;
\alpha_k^+ = u_k a_k + v_k a_k^*;
\alpha_k^- = u_k a_k^+ - v_k^* a_k;
\alpha_k^+ = u_k a_k + v_k^* a_k^*; \tag{3}
$$

it is possible to construct some interactions that help us to promote the nucleon from one diabatic level to another. The two situations plotted in Fig. 1 can be modeled. In the plot (a), the single particle follows the diabatic level $\epsilon_i$ while in (b) it stays on the adiabatic one $\epsilon_j$. Here $a_k^+$ and $a_k$ denote operators for creating and destroying a particle in the state $k$, respectively. The state characterized 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. The interaction able to promote the unpaired nucleon from one adiabatic level to another must be given by product of operators of the type (3).

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

$$
\delta L = \delta < \varphi | H - \hbar \frac{\partial}{\partial t} + H' - \lambda N | \varphi >, \tag{4}
$$

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

$$
\varphi(t) = \sum_m c_m(t) a_m^+ \prod_{l \neq m} (u_l(t) + v_l(t) a_l a_l^+). \tag{5}
$$

FIG. 1: Ideal avoided crossing region. Possible transition states in an avoided crossing region for an unpaired nucleon in the superfluid model.
The Lagrangian contains several terms. The first one is the many body Hamiltonian with pairing residual interactions

\[ H(t) = \sum_{k>0} \epsilon_k(t)(a_k^\dagger a_k + a_k^\dagger a_k) - G \sum_{k,l>0} a_k^\dagger a_k^\dagger a_l a_l. \]  

The residual interactions between diabatic levels characterized by the same quantum numbers that are responsible for the Landau-Zener effect are assumed on the form:

\[ H'(t) = \sum_{i,j \neq \neq} h_{ij}(t)\alpha_i^+ \alpha_j. \]  

The sum runs over diabatic levels \( i \) and \( j \). The number of particle operator is:

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

After some calculations, as detailed in Appendix A, the next time dependent coupled channel equations are obtained [18]:

\[ i\hbar \dot{\rho}_l = \sum_{m} p_m \{ \kappa_l \Delta_m^* - \kappa_l^* \Delta_m \}, \]  

\[ i\hbar \dot{\kappa}_l = \sum_{m} p_m \{ (2\rho_l - 1) \Delta_m + 2\kappa_l (\epsilon_l - \lambda) \}, \]  

\[ i\hbar \dot{h}_m = \sum_{j \neq m} h_{mj} (S_{mj} - S_{jm}), \]  

The next notations are used:

\[ \Delta_m = G \sum_{k \neq m} \kappa_k; \]  

\[ \Delta_m^* = G \sum_{k \neq m} \kappa_k^*; \]  

\[ \kappa_k = u_k v_k; \]  

\[ \rho_k = |v_k|^2; \]  

\[ p_m = |c_m|^2; \]  

\[ S_{jm} = e_j^\dagger c_m. \]  

The reality is more complicated. When a diabatic wave function \( i \) is "reflected" in an avoided crossing region \( (i,j) \), this wave function is transformed in a component of the diabatic wave function \( j \). The reality is more complicated. When a diabatic wave function \( i \) is "reflected" in an avoided crossing region, this wave function must be split into two parts: a transmitted diabatic wave function \( i \) and a reflected adiabatic wave function \( j \). That means, the number of wave functions must be doubled after the passage of each avoided crossing region. Therefore, treating the more realistic situations, the system of coupled channel equations becomes much more complicated. For simplicity, in our approximations we considered only a superposition of \( n \) diabatic wave functions, that means, the diabatic wave function \( i \) is forced to contribute to the amplitude of the diabatic wave function \( j \) (which is not always equivalent to \( j' \)). After some calculations, as detailed in Appendix B a new set of TDHFB equations results:

\[ i\hbar \dot{\rho}_{l(m)} = \kappa_{l(m)} \Delta_m^* - \kappa_{l(m)}^* \Delta_m. \]  

B. Blocking effect

If the blocking effect is taken into consideration, each seniority one BCS wave function is characterized by its own set of \( \rho \) and \( \kappa \) values, and the trial wave function is:

\[ \varphi(t) = \sum_{m} c_m(t) a_m^+ \prod_{i \neq m} (u_i(t) + v_i(t) a_i^+ + v_i(t) a_i^+). \]

The Landau-Zener interaction is postulated as follows:

\[ H'(t) = \sum_{i,j \neq \neq} h_{ij}(t) \alpha_i^+ \alpha_j \]  

\[ \times \alpha_{ij}^+ \alpha_{ji} \prod_{k \neq i,j} \alpha_{k(i)} a_k^+ a_k^+ = \sum_{i,j \neq \neq} h_{ij}(t) (u_i a_j^+ - v_i a_j^+)(u_j a_i^+ + v_j a_i^+) \]  

\[ \times \prod_{k \neq i,j} \alpha_{k(i)} a_k^+ a_k^+, \]  

where the quasiparticle creation and annihilation operators

\[ \alpha_{k(j)} = u_k a_k^+ - v_k a_k^+; \]  

\[ \alpha_{k(j)}^i = u_k a_k^+ + v_k a_k^+; \]  

\[ \alpha_{k(j)}^i = u_k a_k^+ - v_k a_k^+; \]

\[ \alpha_{k(j)}^i = u_k a_k^+ + v_k a_k^+. \]  

The operation is realized on the \( n \) possible diabatic states of the unpaired nucleon. In this paper, the sum over pairs energy generally runs within the index \( k \). When the single-particle sum over \( k \) is realized only for one partner of each reversed pair the result is multiplied with a factor 2. The index \( k \) runs over a workspace that allows the pairing force to operate only over a finite number of active levels around the Fermi energy.
\( i\hbar \dot{\kappa}_l(m) = (2\rho_{l(m)} - 1) \Delta_m + 2\kappa_l(m) (\epsilon_l - \lambda_m), \) 
\( \) (18)

\[ i\hbar \hat{p}_m = \sum_{j \neq m}^n h_{mj} (S_{mj} - S_{jm}), \] 
\( \) (19)

\[ i\hbar S_{jm} = S_{jm} \left\{ -\frac{i}{\hbar} \left[ (\Delta_m |^2 - |\Delta_j|^2 \right) 
+ (\epsilon_m(t) - \epsilon_j(t) - \lambda_m + \lambda_j) \right] 
- \frac{1}{2} \sum_{k \neq m} \left( \Delta_m \kappa^*_k(m) + \Delta^*_m \kappa_k(m) \right) \left( \frac{\rho^2_{c(k(m))}}{|\kappa^*_k(m)|^2} - 1 \right) \right. 
+ \left. \frac{1}{2} \sum_{k \neq j} \left( \Delta_j \kappa^*_k(j) + \Delta^*_j \kappa_k(j) \right) \left( \frac{\rho^2_{c(k(j))}}{|\kappa^*_k(j)|^2} - 1 \right) \right\} 
+ \sum_{l \neq j, m}^n \left[ h_{ml}(t) S_{lj} - h_{jl}(t) S_{lm} \right] + h_{mj}(t) (p_j - p_m). \] 
\( \) (20)

The same notations as in the previous approach are used.

Two main differences arise between Eqs. (17)-(18) and (17)-(20) that are implicitly determined by the hypothesis assumed in their derivation. Firstly, in Eqs. (17)-(18) the values of \( \dot{\rho} \) and \( \dot{\kappa} \) are obtained through a weighted sum that runs over unpaired states while in Eqs. (17)-(18) these quantities belong to only one diabatic wave function. Secondly, in Rel. (20), \( S_{jm} \) depends on all the densities \( \rho \) and pairing moment components \( \kappa \) of the implied two diabatic wave functions \( j \) and \( m \). As a consequence of these differences, the number of differential equations increases \( n \) in Eqs. (17)-(18). Another consequence is that \( \sum_{k \neq m} \rho_k(m) = N - 1 \) for each diabatic wave function \( m \) in Eqs. (17)-(18) while \( 2 \sum_k \rho_k = N + 2 \rho_F - 1 \) in Eqs. (17)-(20) where \( \rho_F \) denotes the single-particle density of the Fermi level in the initial ground-state configuration. Finally, the chemical potential \( \lambda \) has values associated to the diabatic state under consideration in Eqs. (17)-(18).

\[ \] IV. ENERGY

In this section, only the equations associated to the blocking level approach are displayed. For the low lying level approach, the index \( (m) \) must be dropped. The ground state energy \( E_0 \) of any deformation is obtained in the framework of the BCS formalism by considering the Fermi level \( \epsilon_F \) populated with the unpaired nucleon:

\[ E_0 = 2 \sum_{k \neq F} \rho_{c(k)} \epsilon_k + \epsilon_F - G \sum_{k \neq F} |\kappa_{c(k)}|^2 - G \sum_{k \neq F} \rho^2_{c(k)} , \] 
\( \) (21)

in the static, lower energy state. For the same deformation, the energy of an adiabatic state \( m \) is obtained by considering the unpaired nucleon located on the diabatic state under consideration:

\[ E_m = 2 \sum_{k \neq m} \rho_{c(k)} \epsilon_k + \epsilon_m - G \sum_{k \neq m} |\kappa_{c(k)}|^2 - G \sum_{k \neq m} \rho^2_{c(k)} , \] 
\( \) (22)

where the solutions of the TDHFB equations are used. In the frame of our model, the difference

\[ \Delta E_m = E_m - E_0 , \] 
\( \) (23)

behaves as a specialization energy. So, as inferred in Ref. 19 the quantity \( \Delta E_m \) must increase the potential barrier tunneled by the nuclear system. Different barriers are obtained for each diabatic state under consideration. These appear as dynamic excitations during the decay process. Combining excitations with occupation probabilities of diabatic states, we obtain

\[ E = \sum_n \rho_m E_m , \] 
\( \) (24)

for the average energy and

\[ \Delta \overline{E} = \sum_n \sum_m \rho_m \Delta E_m , \] 
\( \) (25)

for the averaged dissipated energy during the decay. As mentioned in Ref. 20, the collective kinetic energy is temporarily stored as a conservative potential. This energy subsequently decays partially to the dissipation.

The equations (17)-(20) and (17)-(20) involves only single-particle energies. They conserve the average number of particles because \( 2 \sum_{k \neq m} \rho_{c(k)} = N - 1 \) for any \( m \) (or \( 2 \sum_k \rho_k = N + 2 \rho_F - 1 \)) and \( \sum_m \rho_m = 1 \). The average energy can evolve in time as follows:

\[ \dot{E} = \sum_n \rho_m \left\{ 2 \sum_{k \neq m} \rho_{c(k)} \dot{\epsilon}_k + \dot{\epsilon}_m 
- \dot{G} \sum_{k \neq m} |\kappa_{c(k)}|^2 - \dot{G} \sum_{k \neq m} \rho^2_{c(k)} \right\} . \] 
\( \) (26)

For a stationary system, for which \( \dot{\epsilon}=0 \) and \( \dot{G}=0 \), the total energy is conserved, even if individual values of \( p,\rho \) and \( \kappa \) may still be varying with time. In our treatment, the chemical potential has the values \( \lambda_m \) obtained from BCS equations for each energy levels workspace associated to the diabatic wave function \( m \).

\[ \] V. GENERALIZATION

If the blocked levels are eliminated, the system (19)-(12) reduces to:

\[ i\hbar \dot{\rho}_l = \kappa_l \Delta^* - \kappa^*_l \Delta , \] 
\( \) (27)

\[ i\hbar \dot{\kappa}_l = (2\rho_l - 1) \Delta - 2\kappa_l (\epsilon_l(t) - \lambda(t)) , \] 
\( \) (27)

the well known TDHFB equations [4, 5]. On another hand if the pairing is neglected, the third equation of the system (11) can be written

\[ i\hbar (\dot{c}_m c^*_m + \dot{c}^*_m c_m) = \sum_{j \neq m} h_{mj} \left( c_j c^*_m + c^*_m c_j \right) . \] 
\( \) (28)
Introducing explicitly the time dependence of the amplitudes $c_m$

$$c_m(t) = c_{0m}(t) \exp \left( -\frac{i}{\hbar} \int_0^t \epsilon_m(\tau) d\tau \right),$$

the next relation is obtained:

$$i\hbar (\hat{c}_m c_m^* + \hat{c}_m^* c_m) = \epsilon_m c_{0m}(\epsilon_m - \epsilon_j) + \sum_{j \neq m}^{n} h_{jm} \left[ c_{0j} c_{0m}^* \exp \left( -\frac{i}{\hbar} \int_0^t (\epsilon_j - \epsilon_m) d\tau \right) \right]$$

$$- c_{0j}^* c_{0m} \exp \left( \frac{i}{\hbar} \int_0^t (\epsilon_j - \epsilon_m) d\tau \right).$$

(30)

The last relation is an equivalent form of the Landau-Zener Eq. obtained in the frame of the single-particle model. Furthermore, if the pairing interaction is neglected, $\rho$ and $\kappa$ can be either zero or one, and the fourth equation of the system reduces to

$$i\hbar \left( \hat{c}_m c_m^* + \hat{c}_m^* c_m \right) = \epsilon_j c_{0m}(\epsilon_m - \epsilon_j) + \sum_{j \neq m}^{n} h_{mj} \left[ c_{0j} c_{0m}^* \exp \left( -\frac{i}{\hbar} \int_0^t (\epsilon_j - \epsilon_m) d\tau \right) \right] - h_{jm} \left( c_{0j}^* c_{0m} - c_{0m}^* c_{0m} \right).$$

(31)

After introducing the exponential dependence, the next relation emerges:

$$i\hbar \left( \hat{c}_m c_m^* + \hat{c}_m^* c_m \right) = \epsilon_m(\epsilon_m - \epsilon_j) + \sum_{j \neq m}^{n} h_{mj} \left[ c_{0j} c_{0m}^* \exp \left( -\frac{i}{\hbar} \int_0^t (\epsilon_j - \epsilon_m) d\tau \right) \right]$$

$$- h_{jm} \left( c_{0j}^* c_{0m} - c_{0m}^* c_{0m} \right).$$

(32)

That is another form of the Landau-Zener relation. So, the Landau-Zener equation for single-particle systems (without residual interactions) and the TDHFB equations for quasiparticles are two particular cases of the coupled channel equations. So, this system represents a generalization of the TDHFB equations in the case of seniority one nuclear systems. Similar arguments are valid also for the system.

### VI. RESULTS

To solve the TDHFB equations, only the variations of the single-particle energies $\epsilon_k$ are needed. The simplest way to obtain the evolutions of single-particle energies is to consider a time-dependent single particle potential in which the nucleons move independently. As evidenced in Ref. such a description is within the spirit of the more rigorous Hartree-Fock approximation, which defines the potential self-consistently.

The $^{14}$C emission from $^{223}$Ra will be treated. The fragments issued in this reaction are spherical while the parent is little deformed, allowing a description in terms of a nuclear shape parametrization given by two spheres smoothly joined within a third surface.

A fine structure in the $^{14}$ radioactivity of the $^{223}$Ra was observed in 1989. In the first experiment, the results indicates that $15\pm3\%$ of $^{14}$C decays are transitions on the ground state of the daughter, while $81\pm6\%$ are transitions on the first excited state. In Ref. using the M3Y potential, it was evidenced that the preformation probability must be more favorable for the excited state than for the ground state with a factor of 180. Such a value cannot be accounted from theoretical models without taking into account dynamical ingredients. This is the main reason that the fine structure phenomenon was selected to validate our equations.

The deformation energy of the nuclear system is the sum between the liquid drop energy and the shell effects, including pairing corrections. The macroscopic energy is obtained in the framework of the Yukawa-plus-exponential model extended for binary systems with different charge densities. The Strutinsky prescriptions were computed on the basis of a new version of the supersymmetric two-center shell model. This version solves a Woods-Saxon potential in terms of the two-center prescriptions as detailed in Appendix.

Because the pairing equations diverges for an infinite number of active levels, a limited number of levels are used in the calculations: 31 levels above and 31 levels under the the unpaired Fermi level in the initial ground state configuration that is, $N - 1 = 62$. These levels are selected in terms of the spin projection $\Omega$ on the symmetry axis and kept as a single particle energies workspace. A constant value of the pairing parameter $G = 0.13$ MeV is used.

The least action trajectory was obtained by generalizing in a three-dimensional space the method initiated in Ref. and used extensively to describe the fission processes. The inertia is computed within the Werner-Wheeler method. The trajectory of the decaying system is obtained simultaneously as function of three generalized coordinates, that is, the elongation $R$ (the distance between the centers of the nascent fragments), the necking parameter $C = S/R_3$ (the curvature of the intermediate surface) and $R_1/R_2$ (the ratio between the radii of the heavy fragment $R_1$ and that of the light one $R_2$). These parameters are explained in Appendix. In Fig. the deformation energy $V$ of the nucleus is plotted as function of the elongation $R$. Three excitations of the nuclear systems that correspond to three adiabatic wave functions are also plotted with dotted lines. These excitations are added to the deformation energies obtained in the framework of the macroscopic-microscopic model in order to calculate the penetrabilities as show below. In Fig. (b) and (c), the variations of the necking and mass-asymmetry generalized coordinates are displayed. At $R \approx 10$ fm, a system formed by two spherical tangent nuclei is obtained. The Woods-Saxon potential is presented in Fig. for a sequence of nuclear shapes along the least action path.

The single-particle level schemes for neutrons and protons must be computed along the minimal action trajectory in order to solve the time dependent pairing equations. It is known that $^{223}$Ra has the spin $\frac{3}{2}$ emerging from $1i_{13/2}$. Adiabatically, the unpaired neutron reaches the $2g_{9/2}$ level of the daughter $^{209}$Pb. As also evidenced
FIG. 2: (a) Deformation energy $V$ as function of the distance between the centers of the nascent fragments $R$. The three excitations due to the diabatic levels $\epsilon_i$, $i = 1, 3$ are also plotted with dotted lines. (b) Variation of the curvature of the median surface and (c) of the mass-asymmetry parameter as function of $R$.

FIG. 3: Mean field Woods-Saxon potential $V_0$ as function of the cylindrical coordinates $\rho$ and $z$ for different values of the elongation along the minimal action trajectory. (a) Elongation $R=2$ fm; (b) $R=5$ fm; (c) $R=10$ fm; and (d) $R=15$ fm.

FIG. 4: Neutron energy diagram along the minimal action path as function of the distance between the centers of the fragments $R$. The levels with spin projection $\Omega$ of interest are plotted with thick lines. The levels are labeled within the spectroscopic factors. At the right the first column corresponds to the daughter nucleus while the second one is related to the $^{14}$C.

FIG. 5: (a) Selected neutron energy levels that can be occupied by a single neutron as function of the internuclear distance. Thick lines are the diabatic levels $\epsilon_i$, $i=1-3$, while thin lines are used for the adiabatic ones. (b) Interactions energies $h_{13}$ in the avoided crossing regions.
ability of the unpaired neutron from the adiabatic level $\Omega = 3$ can be understood by an enhanced transition probability in Refs. [12, 18], the fine structure in the level scheme.

For this purpose, Eqs. (13) - (20) are used. Some features concerning the less rigorous low lying levels approach [9] - [10] can be found in Ref. [18].

In Fig. 6 the three selected diabatic levels $\epsilon_m (m = 1, 2, 3)$ are plotted together with the interaction energies $\hbar \omega$ determined by using spline interpolations around level crossings. Diabatically, the unpaired neutron, initially located on the level $\epsilon_1$ that starts from the spherical orbital $1i_{11/2}$ will arrive on the final state $1i_{11/2}$, that is the first single particle excited state of the daughter, after the passage of three avoided level crossing regions.

The initial conditions are determined by solving the BCS equations for the 3 possible seniority one wave functions at $R \approx 1.5$ fm, where the first minimum of the deformation energy is located. The time dependent pairing equations are integrated numerically using the Runge-Kutta method. The occupation probabilities $p_m$ and the dissipated energies given by formula (23) are determined along the minimal action path for an internuclear velocity $\partial R = 1.4 \times 10^5$ m/s. This value can be translated in a time required to penetrate the barrier of about $1.4 \times 10^{-20}$ s.

In Fig. 6 the probability of occupations $p_m$ with an unpaired neutron of the three diabatic levels are presented as function of the internuclear distance. In the bottom panel, the three dissipated energies $\Delta E_m$ and the energy dissipated in the proton level scheme are also displayed.

The branching ratio $r$ between the partial half-life for transitions to the ground state of the daughter and the partial half-life to the first excited state is given by

$$r = \frac{p_1 \exp(-K_1)}{p_2 \exp(-K_2)},$$

(33)

where the index corresponds to the diabatic level $\epsilon_1$ or $\epsilon_2$, and

$$K_m = \frac{2}{\hbar} \int_{R_g}^{R_m} \sqrt{\mu V(R) + D_m(R) - V(R_g)} dR$$

(34)

are the WKB integrals. Here $R_g$ is the ground state elongation, $R_m$ is the exit point from the barrier for the channel $m$, $V(R)$ is the macroscopic-microscopic energy, $D_m(R) = \Delta E_m(R) + \Delta E_p(R)$ is the dissipated energy, $\Delta E_m$ being the specialization energy given by Rel. [23], $\Delta E_p$ denoting the dissipated energy of the proton subsystem, and $\mu$ is the reduced mass. $\Delta E_p$ is calculated with Eqs. [24]. The barriers obtained for $V(R) + D_m(R)$ are plotted in Fig. 2 (a) with dotted lines. The experimental values of $r$ range between 5.4 and 5.9. Our theoretical value is $r=5$, which is in an excellent agreement with experimental data.

In conclusion, two approaches that generalize the Landau-Zener equations for seniority one superfluid systems are presented. The new formalism is valid for any kind of mean field approximations that include a monopole pairing field. The equations that describe

![FIG. 6: (a) Occupation probability $p_1$ for the diabatic level $\epsilon_1$ as function of the internuclear distance $R$. (b) Occupation probability $p_2$ of $\epsilon_2$. (c) Occupation probability $p_3$ of $\epsilon_3$. (d) The solid line corresponds to the dissipated energy $\Delta E_1$ for $\epsilon_1$, the dashed line is the dissipated energy $\Delta E_2$ for $\epsilon_2$, the dotted-dashed line is the dissipation energy $\Delta E_3$ for $\epsilon_3$, and the dotted line is the dissipation energy $\Delta E_p$ for the proton level scheme.](image)
our approaches offer information about the spectroscopic amplitudes and the dissipated energies in different final channels. The new equations were used to reproduce the qualitative and quantitative features of the fine structure phenomenon in cluster decay. Up to now, this phenomenon was not described adequately in the frame of the center shell model based on a Woods-Saxon potential was obtained. A new version of the superasymmetric two center shell model based on a Woods-Saxon potential was developed and used in this context.

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APPENDIX A

The TDHFB equations when the blocking level is neglected will be derived in this Appendix. Following the same prescriptions as in Ref. [9], using the time dependent Hamiltonian [5] within corrections [7], and the trial wave functions [5], the expected value of the Lagrange function is obtained:

$$\frac{1}{2} \sum_{i,j} \kappa_{ij} \Delta^* \Delta + i \hbar \sum_{i} \left( \sum_{k \neq m} \kappa_{ik} \Delta_k \right)$$

where $\left| \phi_k \right> = a_k^+ \left| 0 > \right.$, due to the normalization. The high order term $\left| v_k \right| ^4$ of formula (A3) is neglected. The equality $\hat{u}_k = - (\tilde{v}_k^* v_k + \tilde{v}_k v_k^*)/2u_k$ is also used.

To minimize the functional, the expression (A1) is derived with respect to the independent variables $v_i^*$ and $v_i$. Two equations follow:

$$\sum_n \left| c_n \right| ^2 \left\{ 2v_i^* (\epsilon_i - \lambda) - G \sum_{k \neq m} \kappa_{k} \right\} = 0,$$

(A7)

$$\sum_n \left| c_n \right| ^2 \left\{ 2v_i (\epsilon_i - \lambda) - G \sum_{k \neq m} \kappa_{k} \right\} - \frac{\partial}{\partial v_i} \left( \frac{\tilde{v}_k^* v_k + \tilde{v}_k v_k^*}{2u_k} \right) = \left( \frac{\kappa_{i} \Delta_{m} - \kappa_{i}^* \Delta_{m}^*}{2m} \right),$$

(A8)

where the notations for densities $\rho = | v | ^2$ and pairing moment components $\kappa = u v$ are introduced.

The condition of conservation of the number of particles

$$2 \sum_k \left| v_k \right| ^2 = N + 2 \rho_F - 1,$$

(A9)

was used so that

$$\sum_k \left( \tilde{v}_k^* v_k + \tilde{v}_k v_k^* \right) = 0,$$

(A10)

$$\frac{\partial}{\partial v_i} \left( \tilde{v}_k^* v_k + \tilde{v}_k v_k^* \right) = 0,$$

(A11)

and

$$\frac{1}{2} \frac{\partial}{\partial v_i} \tilde{v}_k^* v_k = - \tilde{v}_k^*.$$ (A12)

Multiplying Eqs. (A7) and (A8) with $v_i$ and $v_i^*$, respectively, and subtracting, the first TDHFB equation (9) is obtained:

$$i \hbar \partial t = \sum_n \left| c_n \right| ^2 \left( \kappa_{i} \Delta_{m} - \kappa_{i}^* \Delta_{m}^* \right),$$

(A13)

where $\sum_n \left| c_n \right| ^2 = 1$ and $\Delta_m = G \sum_{k \neq m} \kappa_k$.

Another equation can be obtained:

$$\kappa_{i} = - \frac{\partial}{\partial v_i} \rho_l + u \tilde{v}_l = \frac{\rho_l}{2u_l} - \frac{\partial}{\partial v_i} \sum_n \left| c_n \right| ^2 \left( \kappa_{i} \Delta_{m} - \kappa_{i}^* \Delta_{m}^* \right) + \frac{1}{2} \frac{\partial}{\partial v_i} \left( \frac{\tilde{v}_k^* v_k + \tilde{v}_k v_k^*}{2u_k} \right) \Delta_m,$$

(A14)

so that the second TDHFB equation (10) follows:

$$i \hbar \kappa_{i} = \sum_n \left| c_n \right| ^2 \left( 2 \rho_l - 1 \right) \Delta_m + 2 \rho_l (\epsilon_l - \lambda) \right\} \Delta_m,$$

(A15)

Using the property

$$\sum_n \left| c_n \right| ^2 = 1,$$

(A16)
so that
\[ \sum_m^n \hat{c}_m^* \hat{c}_m - \sum_m^n \hat{c}_m^* \hat{c}_m = 0, \] (A17)
the Eq. (A1) is derived with respect to \( c_m \) and \( c_m^* \) and set to zero. The next relations are obtained:
\[ -i\hbar \hat{c}_m = c_m^* [2 \sum_{k \neq m} |v_k|^2 (\epsilon_k - \lambda) + (\epsilon_m - \lambda)] - \frac{G}{\sum_k u_k^* v_k} \sum_{k \neq m} \frac{1}{2} (v_k^* \hat{v}_k - \hat{v}_k^* v_k) + \sum_{j \neq m} h_{mj} c_j^*, \] (A18)
\[ i\hbar \hat{c}_m = c_m [2 \sum_{k \neq m} |v_k|^2 (\epsilon_k - \lambda) + (\epsilon_m - \lambda)] - \frac{G}{\sum_k u_k v_k} \sum_{k \neq m} \frac{1}{2} (v_k \hat{v}_k - \hat{v}_k^* v_k) + \sum_{j \neq m} h_{mj} c_j. \] (A19)

Multiplying the relations (A18) and (A19) with \( c_m \) and \( c_m^* \) and subtracting them the next relation follows:
\[ i\hbar (\hat{c}_m c_m + c_m^* \hat{c}_m) = \sum_{j \neq m} h_{mj} (c_j^* c_j - c_j^* c_m). \] (A20)

It is a form of the third TDHF equation (11). For a passage through only one avoided crossing region \((m, j)\), only two amplitudes, \( c_m \) and \( c_j \), can change. On another hand, from the condition of conservation it can be obtained:
\[ \hat{c}_m c_m^* + c_m^* \hat{c}_m = -(\hat{c}_j c_j^* + c_j^* \hat{c}_j). \] (A21)

This condition is fulfilled by the above Eq. (A20), so the equation conserves the norm. Changing indexes, multiplying with amplitudes and subtracting relations (A18) and (A19) the next equation follows:
\[ i\hbar (\hat{c}_m c_j + c_j^* \hat{c}_m) = c_m c_j^* \left\{ \frac{1}{2} (|\Delta_m|^2 - |\Delta_j|^2) + 2 |v_j|^2 (\epsilon_j - \lambda) + (\epsilon_m - \lambda) - \Delta_m - 2 |v_m|^2 (\epsilon_m - \lambda) - (\epsilon_j - \lambda) - \Delta_j \right\} + c_m c_j^* \frac{\hbar}{2} (v_m^* \hat{v}_j - \hat{v}_j^* v_m - v_j^* \hat{v}_j + \hat{v}_j^* v_j) \] (A22)
\[ \text{From Eqs. (A7) and (A8), the expressions in the last parenthesis that involves \( \hat{v} \) and \( \hat{v}^* \) can be obtained:} \\
\[ \sum_{m \neq j} \frac{\hbar}{2} (v_m^* \hat{v}_j - \hat{v}_j^* v_m) = \rho_j (\epsilon_j - \lambda) + \sum_{m \neq j} \frac{\hbar}{2} (\rho_j^2 - \kappa_j^2), \] (A23)

Using the notations (13) and rearranging the terms the fourth TDHF equation (12) is obtained.

**APPENDIX B**

The TDHF equations when the blocking effect is taken into consideration are derived in this Appendix.

Using the corrections (14), and the trial wave functions (14), the expected value of the Lagrange function is:
\[ < \varphi | H - \hbar \frac{\partial}{\partial \varphi} + H' - \lambda N | \varphi > = \sum_m^n | c_m^2 [\varphi] (\epsilon_k - \lambda)^2 (\epsilon_m - \lambda) + (\epsilon_k - \lambda)^2 (\epsilon_m - \lambda) - G | \sum_{k \neq m} u_k^* v_k]^2 \] (B1)

In order to minimize the functional, the expression (B1) is derived with respect to the independent variables \( v_{l(m)}^* \) and \( u_{l(m)}^* \) by taking into account the subsidiary condition (A12). The next relations follows:
\[ \sum_m^n | c_m^2 [2] (2 v_{l(m)}^* (\epsilon_l - \lambda m) - G | \sum_{k \neq m} \kappa_k^* k_{l(m)}^* - \frac{v_{l(m)}^* v_{l(m)}^*}{2u_{l(m)}^*}) + (u_{l(m)}^* - \frac{\rho_{l(m)}^2}{2u_{l(m)}^*}) \sum_{k \neq m} \kappa_k^* k_{l(m)}^* + i \hbar \dot{v}_{l(m)} = 0, \] (B2)
\[ \sum_m^n | c_m^2 [2] (2 v_{l(m)}^* (\epsilon_l - \lambda m) - G | \sum_{k \neq m} \kappa_k^* k_{l(m)}^* - \frac{v_{l(m)}^* v_{l(m)}^*}{2u_{l(m)}^*}) + (u_{l(m)}^* - \frac{\rho_{l(m)}^2}{2u_{l(m)}^*}) \sum_{k \neq m} \kappa_k^* k_{l(m)}^* - i \hbar \dot{u}_{l(m)} = 0. \] (B3)

This system can be solved by considering that the expression in the curly bracket is zero for each value of \( m \). Following a similar way as in Appendix 1, the first two TDHF equations associated to an unpaired nucleon in the state \( m \) emerge:
\[ i \hbar \dot{v}_{l(m)} = \kappa_{l(m)} \Delta_{m} - \kappa_{l(m)} \Delta_{m}, \] (B4)
\[ i \hbar \dot{u}_{l(m)} = (2 \rho_{l(m)} - 1) \Delta_{m} + 2 \kappa_{l(m)} (\epsilon_l - \lambda m). \] (B5)

To obtain the probability that an unpaired nucleon is located on a state \( m \), the expression (B1) must be derived with respect \( c_m \) and \( c_m^* \). Two equations follow:
\[ -i \hbar \dot{c}_m = c_m^* [2 \sum_{k \neq m} |v_k|^2 (\epsilon_k - \lambda m) + (\epsilon_m - \lambda m) - G | \sum_{k \neq m} u_k v_k]^2 \] (B6)
\[ -i \hbar \dot{c}_m^* = c_m [2 \sum_{k \neq m} |v_k|^2 (\epsilon_k - \lambda m) + (\epsilon_m - \lambda m) - G | \sum_{k \neq m} u_k v_k]^2 \] (B7)

Multiplying with complex conjugates and subtracting, the next relation
\[ i \hbar (\hat{c}_m c_m^* + \hat{c}_m^* c_m) = \sum_{j \neq m} h_{mj} (c_j^* c_m - c_j^* c_m), \] (B8)
is obtained. Using the notations [13] the Eq. [19] follows.

From relations [B6] and [B7], another relation can be deduced

\[
\frac{\hbar}{2} \epsilon_m c_j \left[ \frac{1}{\beta} \left( |\Delta m|^2 - |\Delta j|^2 \right) + (\epsilon_m - \epsilon_j - \lambda_m + \lambda_j) + 2 \sum_{k \neq m} \rho_k (\epsilon_k - \lambda_m) - 2 \sum_{k \neq j} \rho_k (\epsilon_k - \lambda_j) \right]
\]

The derivatives \( \dot{v} \) and \( \dot{v}^* \) appear in the previous expression. In order to evaluate quantities where these derivatives intervene, Eqs. [B2] and [B3] are used. The next relation follows.

\[
\frac{\hbar}{2} \left( v_{lm}^* c_j - \dot{v}_{lm} v_{lm} \right) = \frac{\hbar}{2} \left( 2 v_{lm} (\epsilon_j - \lambda_m) \right)
\]

so that relation [B9] becomes:

\[
\dot{v}_{j,m} = \frac{\hbar}{2} \left( \frac{1}{\beta} \left( |\Delta m|^2 - |\Delta j|^2 \right) + (\epsilon_m - \epsilon_j - \lambda_m + \lambda_j) + 2 \sum_{k \neq m} \rho_k (\epsilon_k - \lambda_m) - 2 \sum_{k \neq j} \rho_k (\epsilon_k - \lambda_j) \right)
\]

After some rearrangements of terms, and using notations [13] the Eq. [20] is obtained.

**APPENDIX C**

A two-center shell model with a Woods-Saxon potential was developed recently [31]. An axial symmetric nuclear shape parametrization is used to determine the mean field potential. This nuclear shape parametrization is given by two ellipsoids (of different semi-axis and eccentricities) smoothly joined with a third surface given by the rotation of a circle around the axis of symmetry as displayed in Fig. 7. The parametrization is characterized by 5 degrees of freedom that can be associated, for example, to the elongation \( R = z_2 - z_1 \), to the necking \( C = S/R_3 \), to the mass asymmetry \( \eta = a_1/a_2 \), to the deformations of the two fragments \( b_i/a_i, i = 1, 2 \). Treating the \(^{14}\text{C} \) emission, the deformations of the two fragments can be neglected and the mass asymmetry parameter is considered as \( \eta = R_1/R_2 \). The mean field potential is defined in the frame of the Woods-Saxon model:

\[
V_0(p, z) = -\frac{V_0}{1 + \exp \left[ \frac{\Delta(p, z)}{a} \right]}
\]

where \( \Delta(p, z) \) represents the distance between a point \((p, z)\) and the nuclear surface. This distance is measured only along the normal direction on the surface and it is negative if the point \((p, z)\) is located in the interior of the nucleus. \( V_0 \) is the depth of the potential while \( a \) is the diffuseness parameter. In our work, the depth is \( V_0 = V_0[1 \pm \kappa(N_0 - N_0 + Z_0)] \) with plus sign for protons and minus sign for neutrons, \( V_0 = 51 \text{ MeV}, a = 0.67 \text{ fm}, \kappa = 0.67 \). Here \( A_N, N_0 \) and \( Z_0 \) represent the mass number, the neutron number and the charge number of the parent, respectively. This parametrization, referred as the Blomqvist-Wallborn one in Ref. [31], is adopted because it provides the same radius constant \( a_N \) for the mean field and the pairing field. That ensures a consistency of the shapes of the two fields at hyperdeformations, i.e., two tangent ellipsoids.

In Fig. 3 the mean field potential \( V_0 \) is plotted as function of cylindrical coordinates \( p \) and \( z \) for four nuclear shape configurations obtained along the minimal action path.

The spin-orbit coupling is assumed of the form

\[
V_{ls} = -2\lambda \left( \frac{1}{2mc} \right)^2 (\nabla V_0 \times \vec{p}) \vec{s}
\]

where \( \lambda = 35 \) is a dimensionless coupling constant, \( m \) is the nucleon mass while \( c \) denotes the speed of the light. The spherical components of the operator

\[
L = \nabla V \times p
\]

in cylindrical coordinates are

\[
L^\pm = \mp \hbar e^{\pm i\varphi} \left( \frac{\partial V_0}{\partial p} \frac{\partial}{\partial z} - \frac{\partial V_0}{\partial z} \frac{\partial}{\partial p} \pm i \frac{\partial V_0}{\partial \rho} \frac{\partial}{\partial \varphi} \right)
\]

so that

\[
L_s = \frac{1}{2} \left( L^+ s^- + L^- s^+ \right) + L_z s_z
\]
The next step is to obtain the solutions of the Schrödinger equation
\[
\left[-\frac{\hbar^2}{2m} \Delta + V_0(\rho, z) + V_{ls}(\rho, z) + V_C(\rho, z)\right] \Psi(\rho, z, \varphi) = E \Psi(\rho, z, \varphi)
\]  
(C7)

For protons, a Coulomb term \(V_C\) is added as in Ref. \[30\]. No analytical solutions can be found for such potentials. A suitable eigenvector basis able to diagonalize the Woods-Saxon potential can be obtained with the double center harmonic oscillator model.

A complete analytical eigenvector basis can be only obtained for the semi-symmetric two-center oscillator. This potential corresponds to a shape parametrization given by two ellipsoids that possess the same semi-axis perpendicular on the axis of symmetry. The potential is
\[
V_0(\rho, z) = \begin{cases} 
\frac{1}{2} m \omega_1^2 (z - c_1)^2 + \frac{1}{2} m \omega_2^2 (\rho - a_1)^2, & z < 0, \\
\frac{1}{2} m \omega_1^2 (z - c_2)^2 + \frac{1}{2} m \omega_2^2 (\rho - a_2)^2, & z \geq 0,
\end{cases}
\]  
(C8)

where \(\omega\) denotes the stiffness of the potential along different directions as follows, \(\omega_{z1} = \omega_0 \frac{R_0}{R_{z1}}, \omega_{z2} = \omega_0 \frac{R_0}{R_{z2}}, \omega_\rho = \omega_0 \frac{R_0}{a_1}, \omega_0 = 41A_0^{1/3}, R_0 = r_0A_0^{1/3}\), in order to ensure a constant value of the potential on the surface. The origin on the \(z\)-axis is considered the location of the plane of intersection between the two ellipsoids.

An analytic system of eigenvectors can be obtained for \(V_0\) by solving the Schrödinger equation:
\[
\left[-\frac{\hbar^2}{2m_0} \Delta + V_0(\rho, z)\right] \Psi(\rho, z, \varphi) = E \Psi(\rho, z, \varphi)
\]  
(C9)

The analytic solution of Eq. (C9) is obtained using the ansatz
\[
\Psi(\rho, z, \varphi) = Z(z)R(\rho)\Phi(\varphi)
\]  
(C10)

with
\[
\Phi_m(\varphi) = \frac{1}{\sqrt{2\pi}} \exp(im\varphi)
\]  
(C11)

\[
R_{nm}(\rho) = \sqrt{\frac{2n!}{(n + m)!}} a_\rho \exp \left( -\frac{\alpha^2 \rho^2}{2} \right) \left( \alpha_\rho \rho \right)^m L_n^m(\alpha_\rho^2 \rho^2)
\]  
(C12)

\[
Z_\nu(z) = \begin{cases} 
C_{\nu1} \exp \left( -\frac{\alpha^2 (z - c_1)^2}{2} \right) H_\nu_{z1}[-\alpha z_1(z + c_1)], & z < 0, \\
C_{\nu2} \exp \left( -\frac{\alpha^2 (z - c_2)^2}{2} \right) H_\nu_{z2}[\alpha z_2(z - c_2)], & z \geq 0,
\end{cases}
\]  
(C13)

where \(L_n^m(x)\) is the Laguerre polynomial, \(H_\nu(z)\) is the Hermite function, \(\alpha_i = (m_0 \omega_i/\hbar)^{1/2}\) \((i = 1, 2, \rho)\) are length parameters, and \(C_{\nu1}\) and \(C_{\nu2}\) denote the normalization constants. The quantum numbers \(n\) and \(m\) are integers while the quantum number \(\nu\) along the \(z\)-axis is real and has different values for the intervals \((-\infty, 0]\) and \([0, \infty)\). Imposing conditions for the continuity of the wave function and its derivative, together with those for the stationary energy and orthonormality, the values of \(\nu_1, \nu_2, C_{\nu1}\) and \(C_{\nu2}\) are obtained. Details concerning these solutions and expressions for the normalization constants are found in Refs. \[32, 33\]. For reflection-symmetric shapes, the solutions along the \(z\)-axis are also characterized by the parity as a good quantum number. The basis \(\{C_{\nu1}\}^2\) for the two-center oscillators can be used for a various ranges of models more of less phenomenological \[31, 33\].

On another hand, they are different ways to obtain the single-particle energies for a two-center Woods-Saxon potential. Other recipes are given in Ref. \[36\] where the potentials are expanded in terms of harmonic oscillators functions.
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