Light nuclei production in fusion of heavy ions

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A possible mechanism of the production of light nuclei in fusion reactions is considered. It is shown that the decay of the dinuclear system during its evolution to a compound nucleus yields a substantial rate for the production of light nuclei. The cross section of this process is calculated for the reaction ^{58}Ni+^{58}Ni. The coupling of other modes of motion causes an increase of the asymmetric decay of the dinuclear system.

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I. INTRODUCTION

Two mechanisms can contribute to the production of light nuclei with $Z > 2$ in fusion reactions of heavy ions at low energy. One of them is a cluster evaporation from the compound nucleus [1] and the other is the decay of a dinuclear system (DNS) at a preequilibrium stage of reaction [2,3]. The last mechanism plays a significant role in the process of compound nucleus formation if it proceeds through an increase of the mass asymmetry of the DNS. In this case the fusion is interpreted in the following way: the dinuclear system formed after a capture stage evolves to a compound nucleus via the nucleon transfer from a light nucleus to a heavy one.

It is possible of course that after a capture stage the neck between the nuclei grows quite quickly and a deformed united system is formed. The further processes are determined by the evolution of this nuclear system. Its shape will approach the equilibrium one if the initial distance between the nuclei is less than the value corresponding to the saddle point of the compound nucleus. Otherwise the formation of the compound nucleus is not possible and the system goes to the quasifission channel. Both mechanisms are combined in reality and their relative role can change from reaction to reaction.

In the reactions with heavy nuclei the deformed compound nucleus at the saddle point is more compact than in the case of a dinuclear system. So, the channel connected with the increase of the neck radius contributes mainly to the quasifission cross section. The channel connected with the increase of the mass asymmetry becomes important for the description of the fusion cross section. For example, this was confirmed by the calculation of fusion cross sections in the reactions ^{100}Mo $+^{100}$ Mo and ^{110}Pd $+^{110}$ Pd [4].

In reactions with light nuclei the channel connected with the increase of the neck radius contributes mainly to the compound nucleus formation. However, the competition of two channels is possible. The decay of very asymmetric configurations of the DNS can enhance the yield of light nuclei. Therefore, the investigation of the production of light nuclei gives us additional, although indirect, information on the fusion reaction mechanism.

In this paper we investigate the preequilibrium light particle production in the reaction ^{58}Ni $+^{58}$Ni. This reaction is interesting due to its intermediate place between the reactions with heavy and light nuclei. The compound nucleus formation goes here mainly through the change of the form of the united system. Nevertheless, the channel of the DNS evolution contributes to the production of light nuclei with $Z > 2$. Our model for the production of light nuclei in fusion reactions is given in Sec. II. The results of our calculations for the reaction ^{58}Ni+^{58}Ni and a conclusion are given, respectively, in Secs. III and IV.

II. MODEL

A. Charge distribution

Nuclear fusion proceeding through an increase of the mass asymmetry of the DNS appears to be consistent with the following qualitative picture [4]: (i) after the total dissipation of the initial kinetic energy the rotating DNS is formed. (ii) A diffusion process leads to the exchange of nucleons between the two touching fragments, thus generating a time-dependent distribution in the charge (mass) asymmetry of the DNS. The DNS evolves to a compound nucleus by the transfer of nucleons from a light nucleus to a heavy one. (iii) There is a certain decay probability of the DNS during its evolution to a compound nucleus. Just the asymmetric DNS decay contributes to the production of light nuclei with $Z > 2$ in fusion of heavy ions at low energies. In comparison with deep inelastic transfer reactions, where the DNS decays inevitably, the decay probability of the DNS in fusion reactions is smaller than 1.

As a result of the above stated, the cross section of the production of light nuclei depends on the formation probability P_Z of the DNS configuration with the charge number Z of the light fragment and on the decay probability Λ_Z. Thus, our task is to calculate P_Z and then to determine Λ_Z. In order to calculate P_Z, a diffusion equation can be used [5]. The values of Λ_Z can be calculated by a classical treatment of the DNS using two macroscopic degrees of freedom (the distance between the nuclear centers R and the mass asymmetry of the DNS η) [3].

We determine the cross section of the production of light nuclei by the expression similar to that used in [5]

$$\frac{d\sigma}{dZ} = \frac{\pi}{E_p m A_p} \int_0^\infty dt \int_0^\infty dJ J \Phi(J) P_Z(J,t) G(t) \Lambda_Z(J) , \quad (1)$$

where A_p and E_p are the mass number and the energy of the projectile, respectively. The factor $G(t)$ represents the probability that the interaction time is t. $\Phi(J)$ is the probability of a defined reaction class (deep inelastic transfer, fusion, and so on) with the angular momentum J. Choosing $\Phi(J)$ in (1) we can obtain the contribution of the defined reaction class to the observed charge distribution. The time integration in (1) gives the cross section averaged over possible interaction times. The factor $\pi(E_p m A_p)^{-1} J dJ$ defines the element of the geometric cross section.

Putting $\Lambda_z(J) = 1$ in (1) we get the known expression used for the calculations of the charge distributions in the deep inelastic transfer reactions [5]. Since these reactions occur mainly for J near the critical angular momentum $J_{\rm crit}$, we may use the following parametrizations

$$\Phi(J) = \exp\left(-\frac{|J_{\rm crit} - J|}{\Delta J} \right) , \quad (2)$$

$$G(t) = \frac{1}{\tau_0} \exp\left(-\frac{t}{\tau_0} \right) , \quad (3)$$

where τ_0 is a mean lifetime of the DNS. The expressions (1)–(3) allow one to obtain the contribution of the deep inelastic transfers to the production of light nuclei.

If Z is far from the projectile charge the formation probability P_Z is small for trajectories with $J > J_{\rm crit}$. For J near $J_{\rm crit}$ it increases because of the increase of the interaction time. For $J < J_{\rm crit}$ the interaction time becomes larger and we can simplify (1) as follows:

$$\left(\frac{d\sigma}{dZ} \right)_{J<J_{\rm crit}} = \frac{\pi}{E_p m A_p} \int_0^{J_{\rm crit}} J P_Z(J, \tau_{\rm int}) \Lambda_Z(J) dJ . \quad (4)$$

At this point we suppose that all trajectories with $J < J_{\rm crit}$ have the same interaction time $\tau_{\rm int}$. The value of $J_{\rm crit}$ for the reaction ^{58}Ni+^{58}Ni practically coincides with the value of momentum $J_{B_f=0}$ which corresponds to the situation when the fission barrier of the compound nucleus ^{116}Ba is vanishing. Therefore, at $J < J_{\rm crit}$ the fusion reaction takes place and we can set $\Phi(J) = 1$. By use of (4) the cross section of the production of light nuclei in fusion of heavy ions can be calculated.

It has been assumed above that factors P_Z and Λ_Z can be considered separately. This is possible because the characteristic time for nucleon transition from one nucleus to the other one is less than the decay time of the DNS. In a forthcoming publication we shall consider the nucleon transfer and the DNS decay simultaneously.

B. Formation probability P_Z

To calculate P_Z we use the master equation

$$\frac{d}{dt} P_Z(J,t) = \Delta^{(-)}_{Z+1}(J) P_{Z+1}(J,t) + \Delta^{(+)}_{Z-1}(J) P_{Z-1}(J,t) - [\Delta^{(+)}_Z(J) + \Delta^{(-)}_Z(J)] P_Z(J,t) , \quad (5)$$

which is indeed a suitable tool to describe the evolution of the DNS. In (5) $\Delta^{(\pm)}_Z(J)$ are the transport coefficients which can be calculated microscopically [6] or can be parametrized [5]. Since the transport coefficients obtained in both these approaches are similar [7], we use the parametrization [5]

$$\Delta^{(+)}_Z(J) = kf \exp\left(\frac{U(Z,J) - U(Z+1,J)}{2T} \right) ,$$
$$\Delta^{(-)}_Z(J) = kf \exp\left(\frac{U(Z,J) - U(Z-1,J)}{2T} \right) . \quad (6)$$

Here $U(Z,J)$ is the potential energy of the DNS (driving potential) with the charge number Z of the light fragment at angular momentum J. We have used $U(Z,J)$ instead of the ground state energy of the DNS in accordance with the results [7]. In the calculation of $U(Z,J)$, the distance R for each Z corresponds to a position of the potential pocket minimum (see Sec. III). The local thermodynamic temperature T is calculated by means of the expression

$$T = \sqrt{[U(Z_0,J) + E_0^*(J) - U(Z,J)]/a} ,$$

where $a = A/8$ MeV^{-1}. $U(Z_0,J)$ is the potential energy of the initial DNS and A the nucleon number of the DNS. The excitation energy of the initial DNS $E_0^*(J)$ is the difference between the energy $E_{\rm c.m.}$ in the center of mass system and the value of the nucleus-nucleus potential for R corresponding to the bottom of the pocket. Note that values of $E_{\rm c.m.}$ above the Coulomb barrier are considered in the present paper. In (6) f is the geometric factor,

$$f = 2\pi \frac{R_1 R_2}{R_1 + R_2} d \quad (d = 1.0 \text{ fm}) , \quad (7)$$

where R_1, R_2 are the radii of the interacting nuclei. The value k in (6) defines the time scale ($k = 0.5 \times 10^{20}$ s^{-1} fm^{-2}).

Solving the equations (5) we obtain the dependence of charge distribution on time. To obtain the measurable charge distribution we should multiply P_Z by the decay probability.

C. Decay probability of the DNS

As was mentioned above the decay probability Λ_Z for collisions with $J < J_{\rm crit}$ is smaller than 1 and has to be calculated. In the notation of [3], we take the collective Hamiltonian of the DNS in the form

$$H_{\rm coll} = \frac{1}{2}\mu\dot{R}^2 + \frac{1}{2}B_{\eta\eta}\dot{\eta}^2 + B_{R\eta}\dot{R}\dot{\eta} + U(R,\eta,J) \ , \quad (8)$$

where $\mu = mA_1A_2/(A_1+A_2)$ is the reduced mass. The mass asymmetry is defined by $\eta = (A_1 - A_2)/(A_1 + A_2)$, where A_1 and A_2 are the fragment mass numbers. $U(R,\eta,J)$ is the potential energy of the DNS depending on R, η, and J. One-to-one correspondence between η and Z is assumed. The mass coefficients have the form

$$B_{\eta\eta} = \mu\xi^2 + \tilde{B}_\eta \ , \quad B_{R\eta} = -\xi\mu \ , \quad (9)$$

where

$$\tilde{B}_\eta = \frac{mA^2}{12}\sum_{i=1,2}\frac{R_i^2}{A_i}\sum_{\lambda=2}^{\infty}\frac{2\lambda+1}{\lambda}\left(\frac{I_{\lambda+1/2}(2R_i^2/r_w^2)}{I_{1/2}(2R_i^2/r_w^2)}\right)^2 ,$$

$$\xi = -\frac{A}{2}\left[\frac{R_1}{A_1}\left(1-\frac{r_w^2}{2R_1^2}\right) - \frac{R_2}{A_2}\left(1-\frac{r_w^2}{2R_2^2}\right)\right] . \quad (10)$$

Here, $A = A_1 + A_2, r_w$ is the radius of the window between the nuclei, and $I_{\lambda+1/2}(x)$ is the modified Bessel function. Expressions (10) have been obtained in [3].

It is well seen that $\xi = 0$ at $A_1 = A_2$ and that $B_{R\eta}$ is negligible around the symmetric configuration. However, for configurations with large mass asymmetries $B_{R\eta}$ is not small and should be taken into account. Due to this nondiagonal coupling the energy contained in the mass asymmetry mode is transferred to the radial mode of motion. Therefore, the DNS approaches the radial potential barrier when moving from the Businaro-Gallone ($\eta_{\rm BG}$) maximum of the potential $U(R,\eta,J)$ to more asymmetric configurations. This leads to an increase of the decay probability.

To obtain $\Lambda_Z(J)$, we have to determine the distribution function $f(R,\eta,p_R,p_\eta,t)$ of collective coordinates and conjugate momenta. Neglecting the dependence of the inertia tensor on R the Fokker-Planck equation for f corresponding to (8) can be obtained in the form

$$\begin{aligned}\frac{\partial f}{\partial t} =& -(\mu_{RR}p_R + \mu_{R\eta}p_\eta)\frac{\partial f}{\partial R} \\ & -(\mu_{\eta\eta}p_\eta + \mu_{R\eta}p_R)\frac{\partial f}{\partial \eta} + \frac{\partial U}{\partial R}\frac{\partial f}{\partial p_R} \\ & + \left[\frac{\partial U}{\partial \eta} + \frac{1}{2}\frac{\partial \mu_{RR}}{\partial \eta}p_R^2 + \frac{\partial \mu_{R\eta}}{\partial \eta}p_Rp_\eta + \frac{1}{2}\frac{\partial \mu_{\eta\eta}}{\partial \eta}p_\eta^2\right]\frac{\partial f}{\partial p_\eta} \\ & + \gamma\mu_{RR}\frac{\partial}{\partial p_R}(p_R f) + \gamma\mu_{R\eta}\frac{\partial}{\partial p_R}(p_\eta f) + D\frac{\partial^2 f}{\partial p_R^2} \ .\end{aligned} \quad (11)$$

Here, γ is the radial friction coefficient and D the diffusion coefficient connected with γ by the Einstein fluctuation-dissipation relation

$$D = \gamma T^* \ , \quad (12)$$

where T^* is an effective temperature,

$$T^* = \frac{\hbar\omega_R}{2}\coth\left(\frac{\hbar\omega_R}{2T}\right) . \quad (13)$$

In (13) $\hbar\omega_R/2$ is the zero vibration energy and T is defined above. The tensor μ_{ij} is inverse to the tensor of inertia of (9):

$$\mu_{RR} = \xi^2/\tilde{B}_\eta + 1/\mu \ , \quad \mu_{R\eta} = \xi/\tilde{B}_\eta \ , \quad \mu_{\eta\eta} = 1/\tilde{B}_\eta \ . \quad (14)$$

In order to simplify the solution of (11) we take the average value of $\partial U/\partial\eta$ as it has been done in [3],

$$\left\langle\frac{\partial U}{\partial \eta}\right\rangle = -\frac{A}{2}\frac{U^f - U^i}{A_2^f - A_2^i} \ , \quad (15)$$

where $U^i(U^f)$ is the value of the potential for the initial (final) configuration with $A_2 = A_2^i$ (final $A_2 = A_2^f$) nucleons.

The coupling of modes of motion is important for large $|\eta|$. The decay of the DNS configurations with small η is mainly due to thermal and quantum fluctuations. We assume for simplification that at $|\eta| < \eta_{\rm BG}$ the relative distance between the nuclei is defined by the position of the minimum of the nucleus-nucleus potential. At $|\eta| > \eta_{\rm BG}$ due to the coupling of the modes of motion the DNS approaches the radial barrier and the decay probability increases. For the consideration of this process we can use the approximation (15).

The pocket of nucleus-nucleus potential for the reaction ^{58}Ni $+^{58}$ Ni is deep enough for $\eta = \eta_{\rm BG}$. Therefore, for the solution of (11) we can use the global momentum approach [8] and write the function f in the form of a multidimensional Gaussian with time-dependent parameters. Using (11) we can obtain a system of equations for the first and second momenta of the distribution function [3]. The main assumptions and initial conditions for the solution of this equation are given in [3]. For simplicity we consider the classical motion along η. The distribution function of the distances between the fragment centers is of a particular interest for us:

$$\begin{aligned}P(R,J,t) &= \int f(R,\eta,p_R,p_\eta,t)d\eta dp_R dp_\eta \\ &= [2\pi\chi_{RR}(t)]^{-1/2}\exp\left(-\frac{[R-\overline{R}(t)]^2}{2\chi_{RR}(t)}\right) . \end{aligned} \quad (16)$$

Here $\overline{R}(t)$ and $\chi_{RR}(t)$ are the average value of R and the variance of the radial distribution, respectively. Due to our assumption on the classical motion along η we can write

$$P(R,J,t) = P\left(R,J,\overline{\eta}(t)\right) \ .$$

Using (16) the decay probability can be obtained as follows:

$$\Lambda_Z(J) = \int_{R_b}^{\infty} P(R,J,\overline{\eta})dR \ , \quad (17)$$

where $\overline{\eta}$ corresponds to Z, and R_b defines the barrier position of the nucleus-nucleus potential for given Z. The main advantage of using the Fokker-Planck equation for the dynamic description of the DNS is the possibility to include the penetration through the potential barrier and the influence of thermal and quantum fluctuations.

III. RESULTS

A. Potential energy of the DNS

The calculation of the potential energy of the DNS is needed in order to obtain the transport coefficients (6) and to estimate the value (15). Let us consider here only the interaction of spherical fragments, neglecting a possible deformation. The value of $U(R,Z,J)$ is defined as

$$U(R,Z,J) = B_1 + B_2 + V_{\rm Coul}(R) + V_n(R) + V_{\rm rot}(R,J) - [B_{12} + V'_{\rm rot}(J)] \ , \quad (18)$$

where B_1, B_2, and B_{12} are the binding energies of the fragments and the compound nucleus; V_n, $V_{\rm Coul}$, and $V_{\rm rot}$ are the nuclear, Coulomb, and centrifugal parts of the nucleus-nucleus potential, respectively. The value of U in (18) is normalized to the energy of the rotating compound nucleus by $B_{12} + V'_{\rm rot}$.

The Coulomb potential can be determined as in [9]. The experimental data [10] show that the sticking condition is practically satisfied for the DNS. Therefore, for $V_{\rm rot}$ we have

$$V_{\rm rot}(R,J) = \frac{\hbar^2 J(J+1)}{2(\mu R^2 + j_1 + j_2)} \ , \quad (19)$$

where $j_i = 2mA_iR_i^2/5$ are the moments of inertia of spherical fragments. $V'_{\rm rot}(T)$ is given by

$$V'_{\rm rot}(J) = \frac{\hbar^2 (J+1)}{2j_c} \ , \quad (20)$$

$j_c = 2mAR_c^2/5$, where R_c is the radius of the compound nucleus.

Different versions of the phenomenological potential $V_n(R)$ describing elastic scattering and reaction cross sections in heavy ion collisions can be found in the literature. A detailed analysis [11] of various theoretical schemes allows us to take a folding procedure for the construction of $V_n(R)$. Unfortunately, it is not possible to use the simple proximity potential for the description of the interaction of light nuclei because it overestimates the depth of the potential pocket.

The repulsive core of the folding potential is obtained by using density-dependent nucleon-nucleon forces [12]:

$$V_n(R) = \int \rho_1(\mathbf{r}_1)\rho_2(\mathbf{R} - \mathbf{r}_2)\mathcal{F}(\mathbf{r}_1 - \mathbf{r}_2) d\mathbf{r}_1 d\mathbf{r}_2 \ . \quad (21)$$

The effective nucleon-nucleon interaction has the form

$$\mathcal{F}(\mathbf{r}_1 - \mathbf{r}_2) = C\left[F_{\rm in}\frac{\rho_0(\mathbf{r}_1)}{\rho_{00}} + F_{\rm ex}\left(1 - \frac{\rho_0(\mathbf{r}_1)}{\rho_{00}}\right)\right] \times \delta(\mathbf{r}_1 - \mathbf{r}_2) \ , \quad (22)$$

$$F_{\rm in,ex} = f_{\rm in,ex} + f'_{\rm in,ex}(N_1 - Z_1)/A_1(N_2 - Z_2)/A_2 \ ,$$

where N_i, Z_i are neutron and proton numbers of the fragment "i", $\rho_0(\mathbf{r}) = \rho_1(\mathbf{r}) + \rho_2(\mathbf{r})$, $\rho_{i=1,2}$ is the density of the fragment "i". The following set of parameters of the nucleon-nucleon interaction has been used in our calculations:

$$C = 300 \text{ MeV fm}^3, \quad f_{\rm in} = 0.09, \quad f_{\rm ex} = -2.59,$$

$$f'_{\rm in} = 0.42, \quad f'_{\rm ex} = 0.54 \ .$$

When one considers the interaction of the ion with mass number $A_2 \geq 20$ and a heavy fragment, the nuclear density is taken in the Saxon-Wood form

$$\rho_i(r) = \frac{\rho_{00}}{1 + \exp[(r - R_i)/a_0]} \quad (23)$$

with $\rho_{00} = 0.17 \text{ fm}^{-3}$, $R_i = r_0 A_i^{1/3}$. For smaller A_2 another functional dependence of $\rho_2(r)$ is more realistic:

$$\rho_2(r) = A_2\left(\frac{\kappa^2}{\pi}\right)^{3/2} \exp(-\kappa^2 r^2) \ . \quad (24)$$

The parameters r_0, a_0, and κ have been chosen that they satisfactorily describe the position (R_b) and height (E_b) of the interaction barrier [13]. For this aim small variations of $r_0 = (1.1–1.12)$ fm and $a_0 = (0.45–0.52)$ fm are sufficient. The parameter κ has been allowed for free variation from 0.44 for ^{16}O to 0.671 for ^{4}He.

The dependence of $U(R,Z,J)$ on Z is presented for different angular momenta in Fig. 1. The distance R

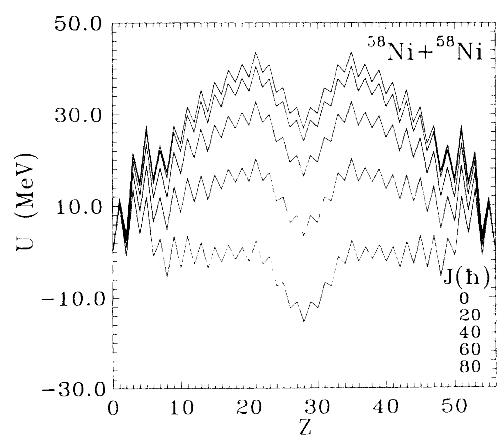

FIG. 1. Driving potential [Eq. (18)] for the system ^{58}Ni $+^{58}$Ni as a function of Z for different values of J. The distance R of each configuration corresponds to the position of the potential pocket minimum. The energy scales are normalized to the total energy of the rotating compound nucleus. The sequence J = 0, 20, 40, 60, 80$\hbar$ is assigned to curves from top to bottom.

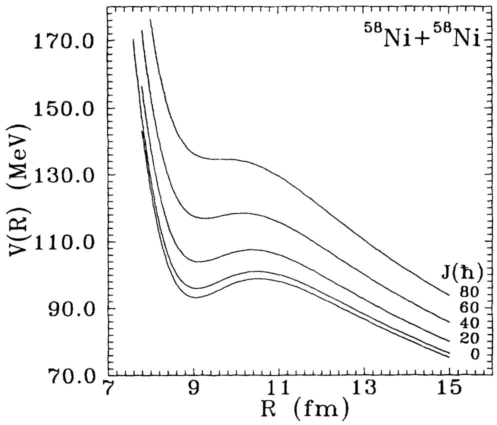

FIG. 2. Radial dependence of nucleus-nucleus potential for the system ^{58}Ni $+^{58}$Ni at different values of J. The sequence $J = 0$, 20, 40, 60, $80\hbar$ is assigned to curves from bottom to top.

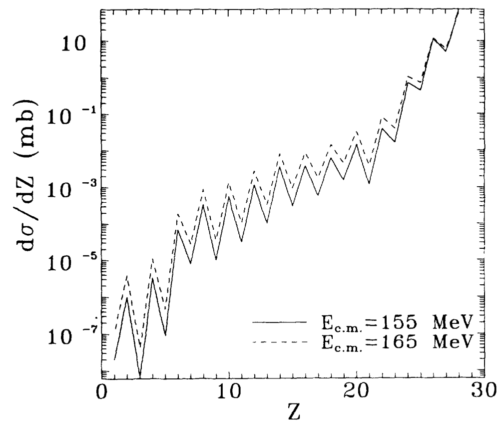

FIG. 3. Calculated cross sections of deep inelastic transfers products in the reaction ^{58}Ni+^{58}Ni at $E_{\mathrm{c.m.}} = 150$ MeV (solid line) and $E_{\mathrm{c.m.}} = 165$ MeV (dashed line) (see text).

for each configuration corresponds to the position of the potential pocket minimum. Binding energies were taken from [14,15]. The mass number of the light nucleus was extracted from the minimization of $U(R, Z, J)$. A large influence of the shell structure of the interacting nuclei on $U(R, Z, J)$ can be seen. This is in agreement with the strong influence of the structure of the light nucleus on the nucleon exchange between the nuclei [6].

At high values of J the energy of the symmetric configuration approaches the energy of the compound nucleus and then becomes less than this energy. At low bombarding energies and high J the DNS cannot overcome the BG maximum and the channel of compound nucleus formation is closed. Instead of compound configurations quasimolecular configurations with a sufficiently long lifetime can occur [16,17]. If we consider neutron deficient nuclei far from the stability line which have relatively small binding energies, then some of their excited states can be thought of as formed by two strongly bound interacting fragments. Constituent fragments are strongly bound because, being lighter, they have such an N/Z ratio that corresponds to the stability line. Due to the balance in binding energies these cluster-type states can appear at relatively low excitation energies. The investigation of the relationship between the DNS configurations and exotic nuclear shapes is a very interesting problem [18].

As it is seen in Fig. 2 the potential pocket disappears for $J > 60\hbar$. The fission barrier of the compound nucleus is equal to zero at $J > 60\hbar$. Therefore, at high angular momenta $[\Lambda_Z(J) = 1]$ only deep inelastic transfers contribute to the production of light nuclei.

B. Deep inelastic transfers contribution

Using (1)–(3) and (5)–(7) we have calculated the charge distribution for deep inelastic transfers. The results for the reaction ^{58}Ni+^{58}Ni at $E_{\mathrm{c.m.}} = 150$ and 165 MeV are shown in Fig. 3. The set of parameters $J_{\mathrm{crit}} = 60\hbar$, $\Delta J = 10\hbar$, $\tau_0 = 2.5 \times 10^{-21}$ s has been

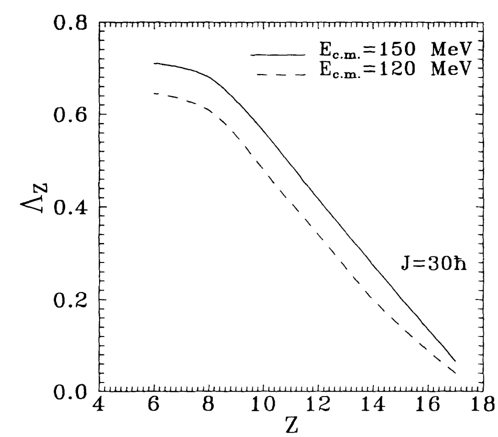

FIG. 4. Decay probability Λ_Z of the dinuclear system as a function of Z in the reaction ^{58}Ni $+^{58}$Ni at $J = 30\hbar$.

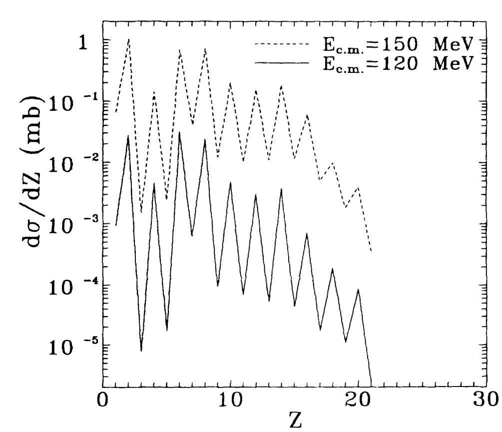

FIG. 5. Calculated charge distribution $(d\sigma/dZ)_{J<J_{\mathrm{crit}}}$ in the reaction ^{58}Ni $+^{58}$Ni at $E_{\mathrm{c.m.}} = 150$ MeV (dashed line) and $E_{\mathrm{c.m.}} = 120$ MeV (solid line).

used. With increasing collision energy the yield of light nuclei increases and a larger number of partial waves contributes to the deep inelastic transfers. Comparing the results in Fig. 3 with the experimental data it is necessary to have in mind that for $Z > 21$ (Businaro-Gallone maximum) the contribution of the highest partial waves is not taken into account (although it exists), because these partial waves do not contribute to the cross section at smaller Z.

C. Contribution of the trajectories with $J < J_{\rm crit}$

If at $J < 60\hbar$ the kinetic energy is large enough, the DNS can overcome the Businaro-Gallone maximum. After this point the system goes to the compound nucleus and its asymmetry increases. Solving (11) and using (17) we obtained the values of $\Lambda_Z(J)$ for each configuration. The thermal and quantum fluctuations rule the value of $\Lambda_Z(J)$ near the Businaro-Gallone maximum. At $\eta > \eta_{\rm BG}$ $\Lambda_Z(J)$ increases with decreasing Z (Fig. 4). The calculated cross sections $(d\sigma/dZ)_{J<J_{\rm crit}}$ (4) are presented in Fig. 5. The dependence of $\Lambda_Z(J)$ on Z is very important and leads to enhanced yields for light nuclei in heavy ion collisions. If we compare our results with the typical ones of the evaporation model, we can see that our calculations give the same cross section for the production of ^{12}C and ^{16}O, although the evaporation cross section for ^{12}C is usually an order of magnitude larger than that for ^{16}O [19]. Thus, our results show that the cross section of the preequilibrium decay of the DNS can be an order of magnitude as large as that given by statistical predictions of light nuclei evaporation from a compound nucleus. Therefore, this mechanism has to be taken into account in the analysis of experimental data.

IV. CONCLUSION

We have calculated the cross sections of the production of light nuclei in the reaction ^{58}Ni+^{58}Ni. Preequilibrium decay of the DNS gives a large contribution to the light nucleus emission. Therefore, careful measurement of the charge distributions in the reaction ^{58}Ni+^{58}Ni at different collision energies will allow us to estimate the relationship between the increase of charge asymmetry and the growing of the neck radius in the DNS. Enhanced experimental yields of light nuclei in comparison with the statistical model predictions will demonstrate the presence of the fusion channel connected with the DNS evolution along charge (mass) asymmetry.

We should mention also that if the energy of the symmetric DNS is smaller than the energy of the compound nucleus, molecular-like states at high angular momenta are formed.

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