Barrier modification in sub-barrier fusion reaction $^{64}$Ni+$^{100}$Mo using Wong formula with Skyrme forces in semiclassical formalism

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Abstract. We obtain the nuclear proximity potential by using semiclassical extended Thomas Fermi (ETF) approach in Skyrme energy density formalism (SEDF), and use it in the extended $\ell$-summed Wong formula under frozen density approximation. This method has the advantage of allowing the use of different Skyrme forces, giving different barriers. Thus, for a given reaction, we could choose a Skyrme force with proper barrier characteristics, not-requiring extra “barrier lowering” or “barrier narrowing” for a best fit to data. For the $^{64}$Ni+$^{100}$Mo reaction, the $\ell$-summed Wong formula, with effects of deformations and orientations of nuclei included, fits the fusion-evaporation cross section data exactly for the force GSkI, requiring additional barrier modifications for forces SIII and SV. However, the same for other similar reactions, like $^{58,64}$Ni+$^{58,64}$Ni, fit the data best for SIII force. Hence, the barrier modification effects in $\ell$-summed Wong expression depend on the choice of Skyrme force in semiclassical ETF method.

1. Introduction
The unexpected behavior of some fusion-evaporation cross sections at energies far below the Coulomb barrier, has challenged the theoretical models to explain the, so called, fusion hindrance phenomenon in true coupled-channels calculations (ccc) for reactions such as $^{58}$Ni+$^{58}$Ni, $^{64}$Ni+$^{64}$Ni, and $^{64}$Ni+$^{100}$Mo [1]. The ccc could, however, be sensitive to the so far unobserved, hence not-included, high-lying states. Misicu and Esbensen [2] were the first who succeeded in describing the above said three reactions in terms of a density-dependent M3Y interaction, modified by adding a repulsive core potential [3]. The repulsive core changes the shape of the inner part of the potential in terms of a thicker barrier (reduced curvature $\hbar\omega$) and shallower pocket. Here, the nuclei are considered spherical in ground-state and the dynamical quadrupole and octupole deformations ($\beta_2$, $\beta_3$) are included (only $\beta_2$ in the case of $^{58}$Ni+$^{58}$Ni).

The dynamical cluster-decay model (DCM) of preformed clusters by Gupta and collaborators [4, 5] is found recently [6, 7] to have barrier modification effects as the inbuilt property, where “barrier lowering” at sub-barrier energies arise in a natural way in its fitting of the only parameter of model, the neck-length parameter. The difference of actually-calculated barrier from the actually-used barrier height, corresponding to the neck-length parameter for best-fitted fusion-evaporation cross section, gives the “barrier lowering” in DCM, whose values are found to increase as the incident energy decreases to sub-barrier energies. Calculations are based on $\beta_2$ deformations and orientation $\theta_i$-dependent nuclear proximity potential of Blocki et al. [8].
Very recently, the Wong formula [9] is also extended by Gupta and collaborators [10, 11] to include its angular momentum ℓ-summation explicitly, which is also shown to contain the barrier modification effects due to the ℓ-dependent barriers. However, for the 58,64Ni-based fusion-evaporation cross sections [1], a further modification of barriers is found essential for below-barrier energies, which is implemented empirically either by “lowering the barrier” or “narrowing the barrier curvature” by a fixed amount for all ℓ’s in the potential calculated by using the proximity potential of Blocki et al. [8], with multipole deformations β2 − β4 and θℓ-integrated for co-planer nuclei. Apparently, the depth of the potential pocket plays no role, in both the DCM and ℓ-summed Wong formula (the two models are same for capture reactions).

In this contribution, we use within the ℓ-summed Wong model, the nuclear proximity potential obtained recently [12] for the Skyrme nucleus-nucleus interaction in the semiclassical ETF approach. Using SEDF, the universal function of proximity potential is obtained as a sum of the parameterized spin-orbit-density-independent and the spin-orbit-density-dependent universal functions (UF’s), with different parameters of UF’s obtained for different Skyrme forces [12]. This method has the advantage of introducing the barrier modifications at sub-barrier energies, if needed, by either (i) modifying the Fermi density parameters (the half-density radii and/ or surface thicknesses, for “exact” SEDF calculations [12]), (ii) the constants of the parametrized UF’s [12] or (iii) change the Skyrme force itself since a different Skyrme force would give different barrier characteristics (height and curvature). This later possibility is exploited here in this paper. It is possible that some Skyrme force would fit the data for one reaction, but not for another reaction and hence requiring “barrier modification” or another Skyrme force.

Section 2 gives briefly the semiclassical ETF method using SEDF, including details of approximations used for adding densities. Section 3 discusses the ℓ-summed Wong formula [10]. Our calculations are given in section 4, and a brief summary of results in section 5.

2. The semiclassical extended Thomas Fermi (ETF) model

The SEDF in semiclassical ETF method provides a convenient way for calculating the interaction potential between two nuclei. In the Hamiltonian density, the kinetic energy density τ as well as the spin-orbit density J are functions of the nucleon density ρq, q = n, p. For the composite system, the densities can be added in either adiabatic or sudden approximation, but we are interested in sudden densities since the different terms of Skyrme Hamiltonian density are then found to constitute the nuclear proximity potential [12, 13, 14]. The sudden densities are defined with or without exchange effects (due to anti-symmetrization), and the one without exchange effects is also referred to as frozen density [15]. In ETF method, the lowest order τ is the Thomas Fermi (TF) kinetic energy density τTF, which already contains a large part of the exchange effects, and that the higher order terms include exchange effects in full. Here we limit τ(⃗r) and J(⃗r) to second order terms for reasons of being enough for numerical convergence [16].

The nucleus-nucleus interaction potential in SEDF, based on semiclassical ETF model, is

\[ V_\mathcal{N}(R) = E(R) - E(\infty) = \int H(\vec{r}) \, d\vec{r} - \left[ \int H_1(\vec{r}) \, d\vec{r} + \int H_2(\vec{r}) \, d\vec{r} \right]. \]  

where the Skyrme Hamiltonian density

\[
H(\rho, \tau, J) = \frac{\hbar^2}{2m} \tau + \frac{1}{2} t_0 \left[ \left( 1 + \frac{1}{2} x_0 \right) \rho^2 - (x_0 + \frac{1}{2}) (\rho_n^2 + \rho_p^2) \right] \\
+ \frac{1}{12} t_1 \rho^2 \left[ \left( 1 + \frac{1}{2} x_3 \right) \rho^2 - (x_3 + \frac{1}{2}) (\rho_n^2 + \rho_p^2) \right] + \frac{1}{4} t_1 \left[ t_1 \left( 1 + \frac{1}{2} x_1 \right) + t_2 \left( 1 + \frac{1}{2} x_2 \right) \right] \rho \tau \\
- \frac{1}{4} \left[ t_1 \left( 1 + \frac{1}{2} \right) - t_2 \left( 1 + \frac{1}{2} \right) \right] \left( \rho_n \alpha_n + \rho_p \alpha_p \right) + \frac{1}{16} \left[ t_1 \left( 1 + \frac{1}{2} x_1 \right) - t_2 \left( 1 + \frac{1}{2} x_2 \right) \right] (\nabla \rho)^2 \\
- \frac{1}{16} \left( 3 t_1 + \frac{1}{2} \right) \left( 2 t_2 + \frac{1}{2} \right) (\nabla \rho_n)^2 + (\nabla \rho_p)^2
\]
\(- \frac{1}{2} W_0 \left[ \rho \nabla \cdot \vec{J} + \rho_n \nabla \cdot \vec{J}_n + \rho_p \nabla \cdot \vec{J}_p \right]. \tag{2}\)

Here, \(\rho = \rho_n + \rho_p\), \(\tau = \tau_n + \tau_p\), \(\vec{J} = \vec{J}_n + \vec{J}_p\) are the nuclear, kinetic energy and spin-orbit densities, respectively. \(m\) is the nucleon mass. \(x_i, t_i \ (i=0,1,2,3)\), \(\alpha\) and \(W_0\) are the Skyrme force parameters, fitted by different authors to ground state properties of various nuclei (see, e.g., [17, 18]). Of the available forces, we use the old, well known SHI and SV forces. Coulomb effects are added directly. Recently, Agrawal et al. [19] modified the Hamiltonian density (2) on two accounts, and obtained a new force GSkI: (i) the third term in (2) is replaced as

\[
\frac{1}{2} \sum_{i=1}^{3} t_{3i} \rho^\alpha \left[ (1 + \frac{1}{2} x_{3i}) \rho^2 - (x_{3i} + \frac{1}{2} (\rho_n^2 + \rho_p^2)) \right],
\tag{3}\]

and (ii) a new term due to tensor coupling with spin and gradient is added as

\[
- \frac{1}{16} (t_1 x_1 + t_2 x_2) \vec{J}^2 + \frac{1}{16} (t_1 - t_2) (\vec{J}_p^2 + \vec{J}_n^2).
\tag{4}\]

We have also used this new GSkI force, with additional six, two each of \(x_3, t_3\) and \(\alpha\), constants.

The kinetic energy density in ETF method, up to second order [16], for \(q=n\) or \(p\) is

\[
\tau_q(\vec{r}) = \frac{3}{5} (3\pi^2)^{2/3} \rho^{5/3} + \frac{1}{36} \frac{(\nabla \rho)^2}{\rho} + \frac{1}{3} \Delta \rho + \frac{1}{6} \frac{\nabla \rho \cdot \nabla f_q + \rho \nabla \rho}{f_q} \nabla^2 \rho - \frac{1}{12} \rho q \left( \frac{\nabla f_q}{f_q} \right)^2 + \frac{1}{2} \rho q \left( \frac{2 m}{h^2} \right)^2 \left( \frac{W_0 \nabla (\rho + \rho_q)}{f_q} \right)^2,
\tag{5}\]

with \(f_q\) as the effective mass form factor,

\[
f_q(\vec{r}) = 1 + \frac{2 m}{h^2} \left( t_1 (1 + \frac{x_1}{2}) + t_2 (1 + \frac{x_2}{2}) \right) \rho(\vec{r}) - \frac{2 m}{h^2} \left( t_1 (1 + \frac{1}{2}) - t_2 (x_1 + \frac{1}{2}) \right) \rho_q(\vec{r}).
\tag{6}\]

Note that both \(\tau_q\) and \(f_q\) are each functions of \(\rho_q\) and/or \(\rho\) only.

The spin \(\vec{J}\) is a purely quantal property, and hence has no contribution in the lowest (TF) order. However, at the ETF level, the second order contribution gives

\[
\vec{J}_q(\vec{r}) = - \frac{2 m}{h^2} \frac{1}{2} W_0 \frac{1}{f_q} \rho_q \nabla (\rho + \rho_q).
\tag{7}\]

Note, \(\vec{J}_q\) is also a function of \(\rho_q\) and/or \(\rho\) alone.

Next, for the proximity potential we introduce the slab approximation of semi-infinite nuclear matter with surfaces parallel to \(x - y\) plane, moving in \(z\)-direction, and separated by distance \(s\) having minimum value \(s_0\). Then, following Blocki et al. [8] and Gupta et al. [12, 13, 14], the interaction potential \(V_N(R)\) between two nuclei separated by \(R = R_1 + R_2 + s\), is given as

\[
V_N(R) = 2 \pi R \int_{s_0}^{\infty} e(s) ds = 2 \pi R \left\{ H(\rho, \tau, \vec{J}) - [H_1(\rho_1, \tau_1, \vec{J}_1) + H_2(\rho_2, \tau_2, \vec{J}_2)] \right\} dz = 4 \pi R \gamma b \phi(D).
\tag{8}\]

where \(R\) is the mean curvature radius, and \(e(s)\) is the interaction energy per unit area between the flat slabs giving the universal function \(\phi(D)\) in terms of a dimensionless variable \(D = s/b\), with surface width \(b = 0.99\) fm. The nuclear surface energy constant \(\gamma = 0.9517 [1-1.7826(\frac{N-Z}{A})^2] \)
MeV fm$^{-2}$. \(\phi(D)\) can be calculated “exactly” or parameterized in terms of exponential and/or polynomial functions [12]. For axially deformed and oriented nuclei, \(\vec{R}\) is given in terms of the radii of curvature \(R_{i1}\) and \(R_{i2}\) in the principal planes of curvature of each of the two nuclei \((i=1,2)\) at the points of closest approach (defining \(s_0)\), by

\[
\frac{1}{R^2} = \frac{1}{R_{11}R_{12}} + \frac{1}{R_{21}R_{22}} + \left[\frac{1}{R_{11}R_{21}} + \frac{1}{R_{12}R_{22}}\right] \sin^2 \Phi + \left[\frac{1}{R_{11}R_{22}} + \frac{1}{R_{21}R_{12}}\right] \cos^2 \Phi. \tag{9}
\]

Here, \(\Phi\) is the azimuthal angle between the principal planes of curvature of two nuclei (\(\Phi=0\) for co-planar nuclei). The four principal radii of curvature are given in terms of radii \(R_i(\alpha_i)\) and their first and second order derivatives \(R_i'(\alpha_i)\) and \(R_i''(\alpha_i)\) w.r.t. \(\alpha_i\), where

\[
R_i(\alpha_i) = R_0i[1 + \sum_\lambda \beta_{i\lambda}Y^{(0)}(\alpha_i)], \tag{10}
\]

with \(R_{0\lambda}\) as the spherical or half-density nuclear radius, \(\lambda=2,3,4...,\) as multipole deformations, and \(\alpha_i\) as an angle between radius vector \(R_i(\alpha_i)\) and symmetry axis, measured clockwise from symmetry axis. For the estimation of \(s_0\), we refer to [20] for \(\Phi=0\) and to [21] for \(\Phi \neq 0\).

For nuclear density \(\rho_i\) of each nucleus \((i=1,2)\), we use the temperature T-dependent, two-parameter Fermi density (FD) distribution, which for the slab approximation is given by

\[
\rho_i(z_i) = \rho_{0i}(T) \left[1 + \exp \left(\frac{z_i - R_i(T)}{\alpha_i(T)}\right)\right]^{-1} \quad -\infty \leq z \leq \infty \tag{11}
\]

with \(z_2 = R - z_1 = (R_1(\alpha_1)+R_2(\alpha_2)+s) - z_1\), and central density \(\rho_{0i}(T) = \frac{4A_i}{\pi R_i^2(T)}[1 + \frac{\pi^2 \alpha_i^2(T)}{R_i^2(T)}]^{-1}\). Then, since \(\rho_i = \rho_{ni} + \rho_{pi}\), following our earlier work [12], for nucleon density we define

\[
\rho_{ni} = (N_i/A_i)\rho_i \quad \text{and} \quad \rho_{pi} = (Z_i/A_i)\rho_i, \tag{12}
\]

with half-density radii \(R_{0i}\) and surface thickness parameters \(a_{0i}\) in Eq. (11) at \(T=0\), obtained by fitting the experimental data to respective polynomials in nuclear mass region \(A=4\text{–}238\), as

\[
R_{0i}(T = 0) = 0.9543 + 0.0994A_i - 9.8851 \times 10^{-4}A_i^2 + 4.8399 \times 10^{-6}A_i^3 - 8.4366 \times 10^{-9}A_i^4 \tag{13}
\]

\[
a_{0i}(T = 0) = 0.3719 + 0.0086A_i - 1.1898 \times 10^{-4}A_i^2 + 6.1678 \times 10^{-7}A_i^3 - 1.0721 \times 10^{-9}A_i^4. \tag{14}
\]

The T-dependence in the above formulae are then introduced as in Ref. [22],

\[
R_{0i}(T) = R_{0i}(T = 0)[1 + 0.0005T^2], \quad a_{0i}(T) = a_{0i}(T = 0)[1 + 0.01T^2]. \tag{15}
\]

Also, the surface width \(b\) is made T-dependent [23], \(b(T) = 0.99(1+0.009T^2)\), where \(T\) is related to the incoming center-of-mass energy \(E_{c.m.}\) or the compound nucleus excitation energy \(E_{CN}\) via the entrance channel \(Q_{in}\)-value, as \(E_{CN} = E_{c.m.} + Q_{in} = \frac{3}{2}AT^2 - T\) (\(T\) in MeV).

Next, for the composite system, \(\rho = \rho_1 + \rho_2\), and the \(\tau(\rho)\) and \(\vec{J}(\rho)\) are added as per prescription used. For sudden approximation (with exchange effects),

\[
\tau(\rho) = \tau(\rho_{1n} + \rho_{2n}) + \tau(\rho_{1p} + \rho_{2p}), \quad \text{and} \quad \vec{J}(\rho) = \vec{J}(\rho_{1n} + \rho_{2n}) + \vec{J}(\rho_{1p} + \rho_{2p}), \tag{16}
\]

and for the frozen approximation (equivalently, sudden without exchange effects),

\[
\tau(\rho) = \tau_1(\rho_1) + \tau_2(\rho_2), \quad \text{and} \quad \vec{J}(\rho) = \vec{J}_1(\rho_1) + \vec{J}_2(\rho_2), \tag{17}
\]
with \( \rho_i = \rho_{in} + \rho_{ip} \), \( \tau_i(\rho_i) = \tau_{in}(\rho_{in}) + \tau_{ip}(\rho_{ip}) \), and \( J_i(\rho_i) = J_{in}(\rho_{in}) + J_{ip}(\rho_{ip}) \). In the following, we consider only frozen densities (sudden without exchange), since \( \ell \)-summed Wong formula within sudden-plus-exchange does not fit the \( ^{64}\text{Ni} \)-based reactions data for use of SIII force [24].

Finally, adding the Coulomb and centrifugal interactions to the nuclear interaction potential \( V_N(R) \), we get the total interaction potential for deformed and oriented nuclei [4, 25], as

\[
V_\ell(R) = V_N(R, A_i, \beta_{\lambda_i}, T, \theta_i, \Phi) + V_C(R, Z_i, \beta_{\lambda_i}, T, \theta_i, \Phi) + V_\ell(R, Z_i, \beta_{\lambda_i}, T, \theta_i, \Phi),
\]

with non-sticking moment-of-inertia \( I_{NS} (=\mu R^2) \) for \( V_\ell \). Eq. (18) gives the barrier height \( V_B^\ell \), position \( R_B^\ell \), and the curvature \( \hbar \omega_\ell \) for each \( \ell \), to be used in extended \( \ell \)-summed Wong’s formula [10], discussed in the following section.

3. Extended \( \ell \)-summed Wong Formula

According to Wong [9], in terms of \( \ell \) partial waves, the fusion cross-section for two deformed and oriented nuclei colliding with \( E_{c.m.} \) is

\[
\sigma(E_{c.m.}, \theta_i, \Phi) = \frac{\pi}{k^2} \sum_{\ell=0}^{\ell_{max}} (2\ell + 1) P_\ell(E_{c.m.}, \theta_i, \Phi), \quad k = \sqrt{\frac{2\mu E_{c.m.}}{\hbar^2}}
\]

(19)

with \( \mu \) as the reduced mass. Here, \( P_\ell \) is the transmission coefficient for each \( \ell \) which describes the penetration of barrier \( V_\ell(R, E_{c.m.}, \theta_i, \Phi) \). Using Hill-Wheeler [26] approximation, the penetrability \( P_\ell \), in terms of its barrier height \( V_B^\ell(E_{c.m.}, \theta_i, \Phi) \) and curvature \( \hbar \omega_\ell(E_{c.m.}, \theta_i, \Phi) \), is

\[
P_\ell = \left[ 1 + \exp \left( \frac{2\pi(V_B^\ell(E_{c.m.}, \theta_i, \Phi) - E_{c.m.})}{\hbar \omega_\ell(E_{c.m.}, \theta_i, \Phi)} \right) \right]^{-1},
\]

(20)

with \( \hbar \omega_\ell \) evaluated at the barrier position \( R = R_B^\ell \) corresponding to \( V_B^\ell \). Note, the \( \ell \)-dependent potentials are required here, given by Eq. (18). Carrying out the \( \ell \)-summation in Eq. (19) empirically for a best fit to measured cross-section [10], and on integrating over the angles \( \theta_i \) and \( \Phi \), we get the fusion cross-section \( \sigma(E_{c.m.}) \).

4. Calculations

We have made our calculations for the \( ^{64}\text{Ni}+^{100}\text{Mo} \) reaction, using SIII, SV and GSkI forces, with frozen densities. Fig. 1(a) shows for one \( E_{c.m.} \) and fixed \( (\theta_i, \Phi) \), a comparison of interaction potentials for the three forces, illustrating their barrier height and position to be force-dependent. This is an interesting property, which we use to fit the fusion-evaporation cross-section in Fig. 1(b). Interesting enough, the data [1] for the above said \( ^{64}\text{Ni}+^{100}\text{Mo} \) reaction fit the \( \theta_i \)-integrated \( (\Phi=0^0) \) \( \ell \)-summed Wong formula for only the new force GSkI. The other forces (SIII and SV) would apparently need additional barrier modification effects to be added empirically [10]. From the deduced \( \ell_{max} \)-values, presented in Fig. 1(c), we notice that \( \ell_{max} \) as a function of \( E_{c.m.} \) vary smoothly only for GSkI force, achieving zero value at sub-barrier energies and a tendency to saturate at an above-barrier energy. For the other forces (SIII and SV), not-fitting the data, the \( \ell_{max} \) varies with \( E_{c.m.} \) erratically, in particular at sub-barrier energies, but could also be smoothed by adding further “barrier lowering” or “barrier narrowing” empirically [10].

5. Summary and discussion

Concluding, the \( \ell \)-summed Wong expression, using the barriers calculated in frozen-density approximation in semiclassical extended Thomas Fermi method, based on Skyrme energy density formalism, describes the fusion-evaporation cross-section data for \( ^{64}\text{Ni}+^{100}\text{Mo} \) reaction nicely with the new Syrme force GSkI only. The variation of deduced \( \ell_{max} \) with \( E_{c.m.} \) is found smooth. Other Skyrme forces (SIII and SV) demand additional barrier modifications at sub-barrier
energies for this reaction. However, the same calculation when applied \cite{KumarR2005} to $^{58,64}$Ni+$^{58,64}$Ni data result in a similar good fit, with a smooth dependence of $\ell_{\text{max}}$ on $E_{\text{c.m.}}$, for only the Skyrme force SIII. Thus, barrier-modification or no barrier-modification in $\ell$-summed Wong expression depends on the choice of Skyrme force in semiclassical ETF method for frozen densities.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure1.png}
\caption{(a) Interaction potentials for one $E_{\text{c.m.}}$ and fixed ($\theta_i$, $\Phi$), (b) fusion-evaporation cross-section as a function of $E_{\text{c.m.}}$ for $^{64}$Ni+$^{100}$Mo reaction, calculated by using $\ell$-summed Wong formula integrated over $\theta_i$ ($\Phi=0^0$) and compared with experimental data \cite{JiangCL2004}, and (c) deduced $\ell_{\text{max}}$ values vs. $E_{\text{c.m.}}$, for the Skyrme forces SIII, SV and GSkI, using frozen densities.}
\end{figure}

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