Temperature and \(n-p\) asymmetry dependencies of the level-density parameter in Ni+Mo fusion reactions

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Properties of evaporation residues and the accompanying light particles have been measured in \(^{60}\text{Ni}+^{92,100}\text{Mo}\) fusion reactions at bombarding energies from \(E/A=5\) to \(9\) MeV. The data indicate that these reactions are essentially complete-fusion reactions with only a small amount of nonequilibrium emission at the highest bombarding energy studied. The measured kinetic-energy spectra of evaporated \(n, p, d, t, ^3\text{He}\), and \(\alpha\) particles are compared to statistical-model predictions. It is found that a constant excitation-energy independent level-density parameter is not able to reproduce these spectra over the compound-nucleus excitation-energy range from \(90\) to \(250\) MeV. The kinetic-energy spectra for all particles were fit using the level-density parameterization \(a=A/(7+1.3\times U/A)\text{ MeV}^{-1}\), where \(U\) is the thermal excitation energy. The Coulomb barriers for charged-particle emission were reduced from the standard values to reproduce the peak energy and multiplicity for \(\alpha\) particles. Using these ingredients, the measured mass, velocity, and angular distributions of evaporation residues are also reproduced. The average \(Z\) and \(N\) of the evaporation residues deduced from the light-particle multiplicities are in agreement with the predicted location of the evaporation attractor line. A neutron-proton asymmetry dependence of the level-density parameter is shown to have potentially important consequences for the neutron or proton richness of the evaporation residues. However, no evidence for such an effect is found in the measured data; and from the experimental multiplicity ratio of \(t/^3\text{He}\) it was deduced that any such dependence in the excitation-energy regime \(100< E^*<250\) MeV is very small.

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I. INTRODUCTION

The nuclear level density is an important quantity in determining both the thermal and decay properties of excited nuclei and for the determination of cross sections used in nuclear astrophysics calculations and applied research. The excitation-energy dependence of the level density is often taken to be the Fermi-gas expression \([1-3]\)

\[
\rho(E^*) = \frac{\sqrt{\pi}}{12} \frac{\exp(2\sqrt{aE^*})}{a^{3/2}(E^*)^{5/4}},
\]

where the level-density parameter \(a\) is proportional to the nucleon number \(A\). For a fixed angular momentum \(J\), the level-density expression becomes \([3,4]\)

\[
\rho(E^*,J) = (2J+1) \frac{\hbar^2}{2I} \frac{3/2}{12} \frac{\sqrt{a}}{a^{3/2}(E^*)^{5/4}} \exp(2aU),
\]

where the thermal excitation energy is \(U=E^*-E_{\text{rot}}(J)\), the rotational energy is

\[
E_{\text{rot}}(J) = \frac{J(J+1)\hbar^2}{2I},
\]

and \(I\) is the moment of inertia. At low excitation energy, the level-density parameter has a value of approximately \(A/8\text{ MeV}^{-1}\) with deviations near closed shells \([5]\). In contrast, the Fermi-gas model gives \(a\sim A/15\text{ MeV}^{-1}\). Some enhancement of the level-density parameter over the Fermi-gas value is expected due to the diffuseness of the nuclear surface \([6-8]\), but to explain the larger experimental value for the level density one must also consider enhancements due to rotational and vibrational collective excitations \([9,10]\). These collective enhancements are expected to wash out with increasing excitation energy \([9,11]\) and thus the level-density parameter would be excitation-energy dependent, decreasing with increasing temperature.
At low excitation energies the effective mass for levels near the Fermi energy is close to the nucleon mass, decreasing for both higher and lower single-particle energies [12]. This enhancement in the effective mass at the Fermi level compared to these other single-particle energies is due to the coupling between the single-particle motion and surface vibrations. It is also expected to decrease with increasing excitation energy again, leading to a decrease in the level-density parameter with increasing excitation energy [13–15]. Experimentally, many analyses of α-particle and proton evaporation spectra at moderate to high excitation energies have confirmed that the level-density parameter is smaller than $A/8\,\text{MeV}^{-1}$ [16–20], however, the excitation energy dependence has not always been determined.

The process of light-particle evaporation moves the position of an excited fragment in the chart of nuclides towards a line called the evaporation attractor line [21,22]. Evaporating systems, on average, approach but never cross this line in standard statistical-model calculations. However, residues detected in $26\leq E/A\leq 50\,\text{MeV}$ Xe + Be, C, and Al reactions [23] were found to lie just on the proton-rich side of this line, while the fusion and incomplete-fusion products were expected to lie on the neutron-rich side [21]. It will be shown in this work that if the level-density parameter is dependent on the neutron-proton ($n$-$p$) asymmetry and not just the total nucleon number $A$, then it is possible to cross the attractor line. Recently, Al-Quraishi et al. [24] have considered two possible $n$-$p$ asymmetry dependencies and have used these to fit experimental level densities for $21\leq A\leq 70$. These fitted $n$-$p$ asymmetry dependencies are quite strong and should have significant consequences for evaporation.

In order to measure the excitation-energy dependence of the level-density parameter and to search for any evidence of a $n$-$p$ asymmetry dependence, we have studied the $E/A = 5–9\,\text{MeV}$ $^{60}$Ni + $^{92,100}$Mo fusion reactions. By utilizing a more symmetric entrance channel, large excitation energies ($E^*\leq 250\,\text{MeV}$) can be induced in a bombarding energy regime ($E/A<10\,\text{MeV}$) where we expect very little non-equilibrium emissions giving rise to essentially complete fusion reactions. The bombarding energies studied and the corresponding excitation energies of the compound systems are listed in Table I. Complete-fusion reactions are most advantageous as the excitation energy is well known, and the angular-momentum distribution of the fused systems that lead to evaporation residues can be constrained from measurements of the residue cross sections.

In order to be confident of the small magnitude of non-equilibrium emissions and to constrain statistical-model parameters, a large set of measurements on the properties of the evaporation residues and the coincidence neutrons and light charged particles were performed using beams of $^{60}$Ni projectiles extracted from the ATLAS accelerator facility at the Argonne National Laboratory. Velocity, and angular and mass distributions were measured for the evaporation residues and these are presented in Secs. II A, II B, and II C, respectively. In Sec. II D the measurements of the light-particle multiplicities, angular distributions, and kinetic-energy spectra are detailed. The discussion of these results and comparisons with statistical-model simulations are made in Sec. III. Finally, in Sec. IV, the conclusions of this work are given.

### II. EXPERIMENTS

#### A. Residue velocity distributions

The velocity distributions of evaporation residues were measured with a 300-$\mu\text{m}$ Si detector of active area 2 cm$^2$. Beams of $E/A = 9.0\,\text{MeV}$ $^{60}$Ni projectiles bombarded thin ($\sim 140\,\mu\text{g/cm}^2$) $^{100}$Mo and $^{92}$Mo foils mounted on 60 $\mu\text{g/cm}^2$ carbon backings. The targets were orientated so that the evaporation residues did not pass through the carbon backings. When using Si detectors to measure particle time-of-flight information, care must be taken as there is a species-dependent plasma delay. In order to avoid introducing systematic uncertainties in the velocity measurement due to the use of some assumed form for this delay, a measurement procedure was utilized, which did not require knowledge of this quantity. The Si detector was mounted on a plunger arrangement allowing measurements at two well-determined time-of-flight distances. The time distribution with respect to the accelerator rf signal was measured for the evaporation residues emitted at $4^\circ$ in the reactions with the $^{100}$Mo target with flight distances of 76.0 and 176.0 cm. From these measurements the zero of the time of flight was determined. Subsequent measurements were performed at $2^\circ$ with the 176.0 cm flight distance.

Experimental invariant velocity distributions $(1/\text{str})\,d\sigma/dv$ of evaporation residues are shown in Fig. 1. The filled arrows indicate the average value $v_{\text{CE}}$ expected for residues produced in complete-fusion reactions, taking into account their average energy loss in the target predicted by the SRIM code [25]. The distributions, practically Gaussian in shape, are centered very near to these indicated complete-fusion velocities. Mean experimental velocities ($v_{\text{UR}}$) were obtained from Gaussian fits to these distributions. The ratio of these velocities, relative to the complete-fusion values, are listed in Table II. These ratios are all within 1% of unity, consistent with complete fusion within the experimental uncertainty, and thus any contribution from nonequilibrium emissions must be small. As a consistency check, the velocity distributions were also determined for elastically scattered projectiles using the same procedure. The ratio of the mean velocity of these distributions to their expected value is

<table>
<thead>
<tr>
<th>$E/A$ (MeV)</th>
<th>$^{100}$Mo</th>
<th>$^{92}$Mo</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>91</td>
<td>76</td>
</tr>
<tr>
<td>6</td>
<td>129</td>
<td></td>
</tr>
<tr>
<td>7</td>
<td>166</td>
<td>145</td>
</tr>
<tr>
<td>8</td>
<td>204</td>
<td>182</td>
</tr>
<tr>
<td>9.1</td>
<td>245</td>
<td>222</td>
</tr>
</tbody>
</table>

TABLE I. Complete-fusion excitation energies in MeV for the $^{60}$Ni + $^{92,100}$Mo reactions studied in this work at the indicated bombarding energies $E/A$. 

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TABLE II. Experimental average velocity of evaporation residues measured at the indicated angles $\theta_{lab}$ over a 176 cm flight distance and for the two targets. The solid arrows indicate the reaction center-of-mass velocity, while the open arrows indicate the expected mean velocity after the nonequilibrium light-particle emission deduced in Sec. II D. The curves show statistical-model predictions (Sec. III B), the height of which has been normalized to reproduce the experimental magnitude.

<table>
<thead>
<tr>
<th>Target</th>
<th>$\theta_{lab}$ (deg)</th>
<th>$\langle v_{ER}/v_{CF} \rangle$</th>
<th>$\langle v_{ER}/v_{CF}^{exp}/v_{CF}^{th} \rangle$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{100}$Mo</td>
<td>4</td>
<td>0.995 ± 0.009</td>
<td>1.003 ± 0.005</td>
</tr>
<tr>
<td>$^{103}$Mo</td>
<td>2</td>
<td>0.992 ± 0.009</td>
<td>1.002 ± 0.005</td>
</tr>
<tr>
<td>$^{92}$Mo</td>
<td>2</td>
<td>0.989 ± 0.009</td>
<td>1.001 ± 0.005</td>
</tr>
</tbody>
</table>

widths are all approximately 0.24 cm/ns (FWHM) and are also much larger than the contribution from variations in energy loss in the target of 0.042 cm/ns predicted by the SRIM code [25]. Therefore it is concluded that these widths are dominated by the recoil motions associated with evaporation of light particles.

B. Residue angular distributions

The angular distributions of the evaporation residues were measured in the $E/A = 9$ MeV bombardments discussed in the preceding section as well as for bombardments at $E/A = 8.0$ and 5.0 MeV. Evaporation residues and scattered beam particles were detected in six $5 \times 5$ mm$^2$ windowless Si photodiodes. These detectors were located 36 cm from the target with fixed angular separations of 1.2° and were each collimated to ±0.3°. Measurements were performed at a number of angular settings to cover the range from −13° to 9°. Particle identification was obtained from time-of-flight and energy measurements. Evaporation-residue cross sections were normalized with respect to the yields of elastically scattered beam particles measured in the same detectors and in fixed monitor counters located at ±15°. The measured evaporation-residue angular distributions $d\sigma/d\theta$ obtained with the $^{100}$Mo and $^{92}$Mo targets are shown in Figs. 2 and 3, respectively. The rms multiple-scattering angle of residues in the target estimated from the SRIM code is small (0.3° for the $E/A = 9.0$ MeV). Therefore, the measured angular distributions are also dominated by the recoil motions associated with the evaporation of light particles.

The curves in Figs. 2 and 3 are statistical-model predictions assuming complete fusion (see Sec. III B). They have been normalized to the experimental data and are used to extrapolate the yield to smaller angles. Integration of these angular distributions yielded the evaporation-residue cross sections $\sigma_{ER}$ listed in Table III. The range of $\ell$ waves contributing to residue formation is estimated from the sharp-cutoff model, $\sigma_{ER} = \pi\lambda^2\ell_{max}(\ell_{max}+1)$, where $\ell_{max}$ characterizes the largest $\ell$ waves associated with residue formation. Values of $\ell_{max}$ are plotted as a function of bombarding energy in Fig. 4 and indicate that residue formation is restricted to $\ell$ waves of $\sim 60\hbar$ and less, with only a small dependence on bombarding energy. This result is not surprising as fission competition is expected to be the limiting factor.

C. Residue mass distributions

Mass distributions of the evaporation residues formed in fusion reactions of $E/A = 9.2$ MeV $^{66}$Ni projectiles on $\sim 530 \mu g/cm^2$ self-supporting $^{100}$Mo and $^{92}$Mo targets were obtained with the fragment mass analyzer (FMA) [26,27]. This device is an 8-m-long recoil mass spectrometer consisting of an achromatic symmetric combination of two electric quadrupoles and one magnetic dipole. At the focal plane, the reaction products are dispersed by $A/Q$ (mass/charge) and are detected as they pass through a position-sensitive parallel-grid avalanche counter (PGAC). After a further distance of 40 cm, the products are stopped in $5 \times 5$ cm$^2$ double-sided Si strip detector (DSSD) with sixteen strips a-
ranged in orthogonal directions on either side of the detector. The angular acceptance of the FMA was defined by a 4.5° × 4.5° entrance aperture. A 20 μg/cm² carbon foil was placed before this aperture to reset the equilibrium charge distribution of the residues before they entered the FMA.

Evaporation residues were separated from other reactions products that reached the focal plane from time-of-flight and energy measurements. The position along the focal plane measured by the PGAC is a measure of the $A/Q$ ratio of the residue. However, as the residue mass distribution is wide and distributions associated with more than four charge states pass through the PGAC, the $A/Q$ information is not sufficient by itself to isolate any $A, Q$ combination. To determine the residue mass distribution, a moderate-resolution determination of a detected fragment’s relative mass is necessary. This was derived from the time-of-flight measurement of the residue through the FMA to the PGAC and the energy measured with the DSSD. Corrections to the residue energy measured by the DSSD were applied to account for the pulse-height deficit in Si and the energy lost in the gas volume and windows of the PGAC.

The electric and magnetic fields of the FMA were set such that the species with charge state $Q_0$, mass $A_0$, and energy $E_0$ passed through the center of the device with a time of flight of $t_0$. If

### Table III. Experimental evaporation-residue cross sections measured with the two Mo targets.

<table>
<thead>
<tr>
<th>$E/A$ (MeV)</th>
<th>$\sigma_{ER}^{100\text{Mo}}$ (mb)</th>
<th>$\sigma_{ER}^{92\text{Mo}}$ (mb)</th>
</tr>
</thead>
<tbody>
<tr>
<td>5.0</td>
<td>282 ± 22</td>
<td>185 ± 35</td>
</tr>
<tr>
<td>8.0</td>
<td>201 ± 22</td>
<td>184 ± 15</td>
</tr>
<tr>
<td>9.0</td>
<td>212 ± 22</td>
<td>172 ± 15</td>
</tr>
</tbody>
</table>

FIG. 2. Evaporation-residue angular distributions measured for the $^{60}\text{Ni} + ^{100}\text{Mo}$ reactions at the indicated bombarding energies. The data points represent the experimental results and the curves are statistical-model predictions normalized to fit these data. The dashed lines indicate the maximum and minimum acceptance angles of the parallel-plate avalanche counter (PPAC) used for the measurements of coincident light particles (Sec. II D).

FIG. 3. As for Fig. 2, but now for the $^{60}\text{Ni} + ^{92}\text{Mo}$ reaction.

FIG. 4. Maximum $\ell$ waves contributing to evaporation-residue production as a function of $E/A$. Results determined for the $^{60}\text{Ni} + ^{100}\text{Mo}$ reaction are indicated by the open data points, where the highest energy point comes from Ref. [32]. The solid data points are for the $^{60}\text{Ni} + ^{92}\text{Mo}$ reaction. The curves show statistical-model fits to these data (Sec. III B).
the values of the constants used in Eq. (7) for trajectories in the FMA and for a constant flight distance.

<table>
<thead>
<tr>
<th>Constant</th>
<th>FMA</th>
<th>Constant flight distance</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K_c$</td>
<td>-0.54</td>
<td>-0.50</td>
</tr>
<tr>
<td>$K_m$</td>
<td>0.44</td>
<td>0.50</td>
</tr>
<tr>
<td>$K_{ee}$</td>
<td>0.46</td>
<td>0.38</td>
</tr>
<tr>
<td>$K_{em}$</td>
<td>0.18</td>
<td>-0.13</td>
</tr>
<tr>
<td>$K_{en}$</td>
<td>-0.13</td>
<td>-0.25</td>
</tr>
</tbody>
</table>

we express deviations from these selected values in terms of $\delta m$, $\delta e$, and $\delta t$ as

\[
\frac{A}{Q} = \frac{A_0}{Q_0} (1 + \delta m),
\]

\[
E = \frac{E_0}{Q_0} (1 + \delta e),
\]

\[
t = t_0 (1 + \delta t),
\]

then, to second order, the time of flight for the trajectories followed by particles in the FMA is

\[
\delta t = K_c \delta e + K_m \delta m + K_{ee} (\delta e)^2 + K_{em} (\delta m)^2 + K_{en} \delta e \delta m.
\]

The values of the various $K$ parameters are listed in Table IV. They are not too different from the values expected for a constant flight distance, which are also listed. Using the measured time of flight and energy, Eq. (7) was solved to give a mass that we will call $A_1$.

The correlation between $A/Q$, measured with the PGAC, and $A_1$ is shown in Fig. 5 for products detected in $^{68}$Ni + $^{100}$Mo reactions when the axis of the FMA was positioned at $\theta_{FMA} = 0^\circ$. The most striking feature in this plot is the array of ridges, each of which is oriented predominately along the vertical mass dimension. For identification purposes, these ridges are numbered relative to the ridge located at the center of the focal plane, which is assigned a value of zero. As examples, the ridges assigned the values $-3$, $0$, $+3$ are indicated by the arrows in the figure. For this measurement, the fields of the FMA were set up for the species with $Q_0 = 32$, $A_0 = 141$, and $E_0 = A_0 v_C^2/2$. The origin of the ridges can be understood by noting that, by taking any $(A/Q)$ pair in the vicinity of these selected values, then the $A/Q$ ratio for this pair will be changed very little by reducing $A$ by 4 and $Q$ by one and similarly reducing $A$ by 8 and $Q$ by 2, etc. In the correlation between $A_1$ and $A/Q$, all such combinations line up, and, due to the larger resolution of the $A_1$ measurement, merge to form the ridges. However, the distinction between adjacent masses on the ridges is not lost, the yield along the crest of the ridges has structures that can be attributed to these masses. This is most evident in the figure for the ridges located at the largest values of $A/Q$ where the neighboring masses become better resolved. By grouping the data into these ridge structure, adjacent mass peaks are separated by

![FIG. 5. Distribution of relative mass versus $A/Q$ for evaporation residues detected near $0^\circ$ in the $^{68}$Ni+$^{100}$Mo reaction with the fragment mass analyzer. The contour interval is constant. The sloping dashed lines indicate the mean location of the residues with the same charge state. The vertical dashed line indicates the $A/Q$ value associated with the nominal center of the focal plane. The vertical ridge structures clearly evident in the data have been labeled by numbers, three examples are indicated by the arrows.](image)

four mass units and the mass distributions can be extracted without requiring a higher-resolution mass measurement.

When the data are gated on a particular ridge, the $A/Q$ measurements can also be used to give a second independent determination of the mass, which will be designated as $A_2$. The resolution of this second mass measurement is best for the positive-numbered ridges where their orientation deviates most from vertical. A final mass is determined from a weighted average of the two mass measurements, $A = wA_1 + (1-w)A_2$, where $w = \sigma_2^2/(\sigma_1^2 + \sigma_2^2)$ and $\sigma_1$ and $\sigma_2$ are the standard deviations associated with $A_1$ and $A_2$, respectively.

The standard deviation $\sigma_1$ is $\approx 2.2$. For ridge numbered zero, the mass resolution obtained from the $A/Q$ measurement is $\sigma_2 = 1.7$, giving a weighting factor of $w = 0.38$. The value of $\sigma_2$ increases for lower-numbered ridges and becomes similar to $\sigma_1$ for ridge numbered $-3$ where $w = 0.47$. Distributions of the average mass obtained for ridges $-3$ to +4 are shown in Fig. 6. These distributions can be well fit with a combination of equally spaced Gaussians. The fits are indicated by the solid curves and each of the Gaussian components are indicated by a dashed curve. The fitted standard deviations of these Gaussians vary from 1.7 for ridge numbered $-3$ to 1.3 for ridge numbered +4.

From these fits, mass distributions are obtained for each charge state. These are plotted in Figs. 7 and 8 for measurements with the $^{100}$Mo and $^{92}$Mo targets, respectively. The focal plane of the FMA is not large enough to detect the full mass distribution for any one of the charge states, but portions of the mass distribution from each of the detected charge states are connected together in the figures allowing the full mass distribution to be estimated. This merging of the various fragments of the distribution seems to work very well, as the yields of the same $A$ in different charge states are similar, especially for the data obtained with the $^{100}$Mo target.
shown in Fig. 7. Values of $Q_0$ were selected which correspond to the expected peak in the charge-state distribution and therefore no large dependence on the mass yield on $Q$ are expected. This was checked experimentally for the 100 Mo target for which the residue yield was briefly measured for other regions of $A/Q$ and the selected $A/Q$ region was found to be located at the maximum in the yield. Therefore, distortions to the mass distributions due to the charge-state distributions are not expected to be too serious.

In assigning absolute masses to all of the peaks, it is important to identify which of the peaks on the central ridge, numbered zero in Fig. 5, is associated with the selected values of $Q_0$ and $A_0$ which should be located at the center of the focal plane. The absolute calibration of $A/Q$ on the $x$ axis of this figure was obtained using the central position of the focal plane as measured with the PGAC in past experiments. The $A/Q$ ratio associated with this central position is indicated by the vertical dashed line, and unfortunately it passes between the two most prominent peaks on the central ridge. The uncertainty as to which of these two peaks is actually in the central location gives rise to an ambiguity of four mass units in assigning a mass to all observed peaks. Of the two possibilities, the lower mass solution has been selected for the distributions plotted in Figs. 7 and 8. The other solution, which corresponds to the same distributions shifted up by four mass units, was discarded as those distributions are incompatible with the average residue masses deduced from coincident light-particle multiplicities presented in Sec. II D.

D. Light-particle evaporation

Particle evaporation in coincidence with evaporation residues was measured in reactions of $E/A = 5.0, 6.0, 7.0, 8.0$, and $9.1$ MeV $^{60}$Ni beam particles on $770 \mu g/cm^2$ $^{100}$Mo targets and $E/A = 7.0, 8.0$, and $9.1$ MeV beam particles on $716 \mu g/cm^2$ $^{92}$Mo targets. A schematic of the experimental apparatus is shown in Fig. 9. The evaporation residues were detected in an annular parallel-plate avalanche counter (PPAC) which was centered on the beam axis and subtended angles from $2.2^\circ$ to $7.8^\circ$. These limits are indicated on the evaporation-residue angular distributions in Figs. 2 and 3. From these angular distributions we expect to intercept $60\%$ and $77\%$ of the residues at $E/A = 5$ and $9$ MeV, respectively. While one might expect that this selection does not introduce a significant bias, we find that for $E/A = 5$ MeV it does, as

![FIG. 7. Mass distributions measured for each of the indicated charge states for evaporation residues detected in the $^{60}$Ni+$^{100}$Mo reaction at $\theta_{FMA} = 0^\circ$ and $5.5^\circ$. Data for odd and even charge states are plotted as the solid and open data points, respectively. The solid bar indicates the mean residue mass estimated from the multiplicities of evaporated H and He isotopes. The width of the bar reflects the uncertainties of the multiplicity measurements.](image)

![FIG. 8. Same as for Fig. 7, but now for the $^{60}$Ni+$^{92}$Mo reaction.](image)
FIG. 9. A schematic of the experimental apparatus used to measure the evaporation spectra. The locations of the PPAC, CsI(Tl) counters, and the Si telescope inside the scattering chamber are indicated. The arrangement of the neutron counters are also shown. However, in reality the neutron counters were located in a plane perpendicular to that containing CsI(Tl) and Si detectors. This avoids any shadowing of the neutron counters.

the peak of the angular distribution is at very small angles and is missed by the PPAC. This issue will be discussed later.

Coincident light charged particles were detected in a Si $E - \Delta E$ telescope and four $2 \times 2$ cm square CsI(Tl) scintillator detectors, all located approximately 18 cm from the target. The Si telescope, consisting of a 100-$\mu$m-thick $\Delta E$ detector followed by a 5-mm-thick $E$ detector, was placed at 39.4° from the beam axis. The CsI(Tl) detectors of thickness 1.0, 1.0, 1.0, and 0.4 cm were located at angles of 19.4°, 36.9°, 60.0°, and 163.2° from the beam axis, respectively. Light-particle identification with these CsI(Tl) detectors was achieved from pulse-shape information as described in Ref. [29].

The energy calibration of the $E$ and $\Delta E$ Si detectors was determined from the measured spectra of $\alpha$ particles emitted from a $^{228}$Th source and the energies at which $p, d, t$, and $\alpha$ particles punch through these detectors. For the CsI(Tl) detectors, the light output is a nonlinear function of deposited energy. In this work, the scintillation light from the CsI(Tl) crystals was read out by photomultiplier tubes and the energy calibration is based on the integrated light output in the first

![FIG. 10. Experimental neutron kinetic-energy spectra (center-of-mass frame) measured in coincidence with evaporation residues produced in the $^{60}$Ni+$^{100}$Mo reaction at the indicated bombarding energies. The curves are labeled by the laboratory angle that the neutrons were detected at with the solid data points representing the most forward angle detector. Each data point is the average from two detectors located on opposite sides of the beam axis.](image)

$\sim 400$ ns [29]. Calibration data for crystals of similar Tl doping have been obtained in Ref. [30]. The $p, d, t$ and $\alpha$ particle-through energies measured in this work were used to normalize the light output $L$ and match to these established calibration points. In addition, for $\alpha$ particles, calibration points were also obtained using the $^{228}$Th $\alpha$ source. For the energy range of interest in this work, the energy calibration can be parametrized as

$$E = mL + c[1 - \exp(-L/b)]$$

(8)

where the constants $m$, $c$, and $b$ are nuclide dependent and the light output is normalized so that $m = 1$ MeV for protons. For large $L (>3b)$, the calibrations are essentially linear, but with an offset $c$. For $d, t$, and $\alpha$ particles, all three constants were well determined from a fit to the calibration data. As a consistency check, the energy spectra of these particles for the 36.9° detector were compared to those obtained from the neighboring Si telescope at 39.4° (see Figs. 11–13 for comparison). Due to the small difference in angle, the comparison was made after the transformation into the reaction center-of-mass frame. Consistency between the two spectra were found for all three particle types. For protons, the energy calibration of the light output is usually assumed to be linear [29], i.e., $c = 0$. However, this assumption was re-evaluated as full consistency between spectra from the CsI(Tl) detector, and its neighboring Si telescope was not

![TABLE V. Values of the constants in Eq. (8) for the energy calibration of the CsI(Tl) light output.](image)
obtained. The proton calibration points are confined to large energies (~25 MeV) and only constrain the parameter \( c \) to a small value. Therefore the constant \( c \) was chosen so as to make the exponential tail of the proton spectra coincide with the result from the Si telescope. It was then found possible to obtain an extremely good match around the Coulomb barrier region by adjusting the parameter \( b \). Table V lists the values of the parameters \( m \), \( c \), and \( b \) used for the different particle types.

The target and PPAC were mounted in a thin-walled (3.2-mm-Al) spherical scattering chamber of 40 cm radius. The CsI(Tl) detectors and the Si telescope were also located inside the chamber and were mounted in the vertical plane. Outside the chamber, in the horizontal plane, were placed 16 cylindrical neutron detectors of 7.6 cm radius and 7.6 cm thickness, each containing BC501 liquid scintillator. Pairs of detectors were arranged symmetrically at either side of the beam axis at angles of 13°, 29°, 47°, 70°, 105°, 125°, 142°, and 160°. Their separation from the target ranged from 80 cm at the most forward angles to 60 cm at the most backward angles. Pulse-shape information was used to differentiate between detected neutron and \( \gamma \) rays and the detection efficiency was determined from a \(^{252}\text{Cf}\) source at low energies (<5 MeV) and from the SCINFUL code \(^{31}\) at higher energies as in Ref. \(^{32}\). Neutron time of flight was measured with respect to the accelerator rf with a total timing resolution of ~1 ns (FWHM).

Examples of the center-of-mass kinetic-energy spectra of neutrons, protons, and \( \alpha \) particles emitted in coincidence with evaporation residues are shown in Figs. 10–12, respectively, for reactions on the \(^{100}\text{Mo}\) target at three bombarding energies (\( E/A = 5, 7, \) and 9 MeV). A few examples are also shown for deuterons and tritons in Fig. 13. At the lower energy, the proton spectra are shifted upwards compared to the spectra of other particles.
bombarding energies, the spectral shapes for each particle are independent of angle, as is expected for emission from a compound nucleus traveling at the reaction center-of-mass velocity. The spectral shapes are typical of evaporation spectra with Coulomb peaks for the charged particles and then exponentially decreasing high-energy tails. However, at \( E/A = 7 \) MeV, the high-energy tails for the neutrons in Fig. 10 show a slight enhancement at forward angles indicating a small contribution from nonequilibrium emissions. This becomes more prominent at \( E/A = 9 \) MeV where also a nonequilibrium component is now clearly visible for \( \alpha \) particles in Fig. 12. Protons, deuterons, and tritons also show indications of a small contributions at this bombarding energy. However, the characteristic of the neutron and charged-particle nonequilibrium components are different. The nonequilibrium neutron component is still clearly visible in the detector at \( \theta_{\text{lab}} = 70^\circ \), while for charged particles, this component is only observed in the most forward detector, and spectral shapes for the other detectors are all consistent within the experimental error.

In order to determine the evaporative light-particle multiplicities and estimate the contributions from nonequilibrium emissions, the laboratory spectra were each fit with contributions from two thermal-like moving sources, one associated with compound-nucleus emissions and the other with the nonequilibrium component. For neutrons, both sources in their moving frames were taken as isotropic, with “volume” emission spectra parametrized as

\[
\frac{d^2m}{dE d\Omega} = \frac{m}{2(\pi T)^{3/2}} \exp(-E/T),
\]

where \( m \) is the multiplicity and \( T \) is the apparent temperature. Based on the residue-velocity measurements, the compound-nucleus source was constrained to move with the reaction center-of-mass (c.m.) velocity. At the highest bombarding energy, the velocity of the nonequilibrium source was initially left as a fit parameter and we obtained a value that is consistent, within statistical error, to the beam velocity. Subsequently, the nonequilibrium source velocity was fixed to the beam velocity for all other bombarding energies. Thus with two sources there are four fit parameters: \( m_{\text{CN}}, m_{\text{neq}}, T_{\text{CN}}, T_{\text{neq}} \). An example of the fit to the spectra obtained with \( E/A = 9 \) MeV \(^{60}\text{Ni} + ^{100}\text{Mo} \) reaction is displayed in Fig. 14. The solid curves indicate the fitted contributions from both sources, while the dashed curves show the compound-nucleus contribution only. Overall, the two-source parametrization allows a good reproduction of the data. However, the magnitude of the nonequilibrium emissions is not fully accounted for at high kinetic energy for \( \theta_{\text{lab}} = 70^\circ \) and \( 105^\circ \), but the fits are still reasonable estimates of the multiplicity of this component. The fitted multiplicities for the two components are listed in Table VI. For all reactions, the nonequilibrium component represents at most 5% of the total neutron multiplicity.

For charged particles, the compound-nucleus source was taken as “surface” emission with a Gaussian distribution of Coulomb barriers, i.e.,

\[
\frac{d^2m}{dE d\Omega} = \frac{m}{(2\pi)^{3/2} T^2 \sigma_B} \left[ 1 + a_2 P_2[\cos(\theta)] \right] \int_0^\infty (E-B) \times \exp[-(E-B)/T] \exp[-(B-B_0)^2/2\sigma_B^2] dB
\]

in the moving frame where \( P_2 \) is the Legendre polynomial of second order. The mean Coulomb barrier is \( B_0 \) and the standard deviation is \( \sigma_B \). For \( \alpha \) particles only, the parameter \( a_2 \) was permitted to be nonzero to allow for the expected angu-

<table>
<thead>
<tr>
<th>Particle</th>
<th>( m_{\text{CN}} ) (MeV)</th>
<th>( T_{\text{CN}} ) (MeV)</th>
<th>( a_2 )</th>
<th>( m_{\text{neq}} )</th>
<th>( T_{\text{neq}} ) (MeV)</th>
<th>( B_0 ) (MeV)</th>
<th>( \sigma_B ) (MeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>n</td>
<td>8.2 ± 0.9</td>
<td>2.93 ± 0.05</td>
<td>NU</td>
<td>0.28 ± 0.13</td>
<td>4.0 ± 1.0</td>
<td>NU</td>
<td>NU</td>
</tr>
<tr>
<td>p</td>
<td>2.9 ± 0.3</td>
<td>3.0 ± 0.2</td>
<td>NU</td>
<td>0.09 ± 0.05</td>
<td>2.2 ± 1.3</td>
<td>5.4 ± 0.1</td>
<td>1.4 ± 0.1</td>
</tr>
<tr>
<td>d</td>
<td>0.39 ± 0.06</td>
<td>3.5 ± 0.5</td>
<td>NU</td>
<td>0.015 ± 0.012</td>
<td>2.0 ± 1.0</td>
<td>6.8 ± 0.2</td>
<td>1.3 ± 0.1</td>
</tr>
<tr>
<td>t</td>
<td>0.13 ± 0.03</td>
<td>3.7 ± 0.6</td>
<td>NU</td>
<td>0.005 ± 0.005</td>
<td>2.0 ± 1.5</td>
<td>6.6 ± 0.2</td>
<td>1.3 ± 0.2</td>
</tr>
<tr>
<td>(^3\text{He})</td>
<td>0.04 ± 0.01</td>
<td>4.8 ± 1.0</td>
<td>NU</td>
<td>NU</td>
<td>14.0 ± 0.5</td>
<td>2.0 ± 0.5</td>
<td></td>
</tr>
<tr>
<td>( \alpha )</td>
<td>2.1 ± 0.2</td>
<td>3.7 ± 0.2</td>
<td>0.2</td>
<td>0.04 ± 0.03</td>
<td>1.7 ± 1.1</td>
<td>13.5 ± 0.2</td>
<td>2.0 ± 0.1</td>
</tr>
</tbody>
</table>
lar dependence to these emissions, but also to take into account a bias induced by the restricted range for evaporation-residue detection. This bias is discussed extensively in Ref. [19], but briefly, as \( \alpha \)-particle emission imparts the largest recoil kick to a compound system, its emission is most strongly correlated with the angle of the evaporation residue. At the lowest bombarding energy where the residue angular distribution is most focused (Sec. II A), the emission of an \( \alpha \) particle near \( \theta_{\text{c.m.}} = 90^\circ \) recoils the residue to larger angles increasing its probability of being detected and thus enhances the angular distribution of coincident \( \alpha \) particles at this angle. This bias gives rise to a negative value of \( a_2 \), which can be seen for the \( E/A = 5 \) MeV spectra in Fig. 12. In this case, all spectra have the same shape, but the \( \theta_{\text{lab}} = 60^\circ \) \( (\theta_{\text{c.m.}} \approx 90^\circ) \) spectrum is largest in absolute magnitude. At larger bombarding energies where this bias is reduced [19], positive values of \( a_2 \) are obtained from the fit, reflecting the expected unbiased emission pattern.

For the nonequilibrium source we use a form similar to that from Ref. [33], consisting of an isotropic “surface”-type moving source traveling at the beam velocity, but which experiences a Coulomb repulsion from the compound nucleus travelling at the center-of-mass velocity. For the Coulomb repulsion, the same distribution of Coulomb barriers as in Eq. (10) is used. With both sources there are six fit parameters; \( m_{\text{CN}} \), \( m_{\text{neq}} \), \( T_{\text{CN}} \), \( T_{\text{neq}} \), \( B_0 \), and \( \sigma_p \) except for \( \alpha \) particles where the seventh parameter \( a_2 \) is added. For \( ^3\text{He} \) fragments, sufficient particle identification was not obtained with the most forward and most backward detectors to obtained useful spectra. Hence only a compound-nucleus source was used in the fit to the spectra from the remaining detectors. Fit parameters obtained from the \( E/A = 9 \) MeV \( ^{68}\text{Ni} + ^{100}\text{Mo} \) data are listed in Table VI and, as an example, the fitted spectra obtained for \( \alpha \) particles are shown in Fig. 15. The dotted line here indicates the nonequilibrium contribution obtained for the most forward angle CsI(Tl) detector. This component makes a negligible contribution in the spectra from the other detectors. Finally, the compound-nucleus multiplicities extracted for all particles are plotted as a function of bombarding energy in Fig. 16 and listed in Table VII along with the values for nonequilibrium neutrons.

For the lowest bombarding energy, there was some difficulty in obtaining the absolute normalization of the light-particle spectra as there was contamination of the residue gate with multiple random-coincident elastic-scattering events. For the higher bombarding energies, the extracted neutron multiplicity increases linearly with \( E/A \) with high precision, as expected from the statistical-model simulations. Thus the normalization at \( E/A = 5 \) MeV was chosen so that the deduced neutron multiplicity followed this linear trend.

The multiplicities and spectra of coincident light particles were also measured by Gonin et al. for the \( E/A = 9 \) and 11 MeV \( ^{68}\text{Ni} + ^{100}\text{Mo} \) reactions [16]. For the \( E/A = 9 \) MeV reaction, the shapes of the \( n, p, d, t, \) and \( \alpha \) spectra are similar to those measured in this work. However, there are significant differences between the multiplicities deduced in the two studies. The measured multiplicities are 2–4 times smaller for charged particles and 1.7 times larger for neutrons in the work of Gonin et al. Similar discrepancies with the \( E/A = 11 \) MeV data of Gonin et al. were also found in Ref. [32]. It is also very difficult to understand the mechanism for the creation of the extremely neutron-deficient evaporation residues, which are implied by the work of Gonin et al.

Coincidence charged particles were studied in the \( E/A = 5–10 \) MeV \( ^{64}\text{Ni} + ^{100}\text{Mo} \) reactions in Ref. [19]. According to statistical-model simulations, the shape of the kinetic-energy spectra in this and the present work should be very similar as the projectiles differ by only four neutrons. We find some differences between the two sets of spectra and in particular we note that, contrary to the present study, the measured angular distributions for deuterons and tritons in Ref. [19] were not found to be consistent with statistical emission. The high-energy tails of the kinetic-energy spectra were more pronounced at the forward angles even at the lowest bombarding energy. We suspect this observation might be erroneous as the deuteron and triton energy calibrations of the CsI(Tl) detectors in Ref. [19] were not directly determined. Due to the use of a Si telescope and a larger set.
of calibration points for the CsI(Tl) detectors, the energy calibrations for all charged particles are believed to be superior in present study, and therefore the kinetic-energy spectra presented in the present work are preferred to those of Ref. [19].

III. DISCUSSION

A. Characterization of reactions

The experimental data indicate that the evaporation residues are formed in fusion reactions where practically all the nucleons from the projectile and target nuclei are present in the compound nucleus. However, a very small amount of nonequilibrium emissions is present. The most important of these are the nonequilibrium neutrons that have approximately the same multiplicity as the evaporated deuterons (Table VII). Also as indicated in Sec. II D, the nature of the nonequilibrium source may be different for neutron and charged particles. This is further emphasized when one looks at the fitted apparent temperatures in Table VI. For neutrons $T_{neq} > T_{CN}$, while for charged particles $T_{neq} < T_{CN}$. The low temperature for nonequilibrium charged particles indicates these particles are more localized around the beam velocity, suggesting an incomplete-fusion interaction where the nonequilibrium component represents the unfused portions of the projectile which continue traveling forward with velocity close to the beam value. The larger temperature for the nonequilibrium neutrons may be more indicative of emission from a hotspot or emission as Fermi jets as in the models of Refs. [34, 35].

The mean total linear momentum of all the fitted nonequilibrium sources represents \( \sim 1\% \) of the beam momentum at \( E/A = 9 \) MeV. In Fig. 1, the corresponding velocity of the compound nucleus, indicated the open arrowheads, is consistent with the experimental velocity distributions. The actual error of this mean velocity is difficult to estimate as it will depend on whether the assumed form for the nonequilibrium source in the moving source fits can describe it accurately. This is difficult to ascertain, as most of the nonequilibrium is buried under of the compound-nucleus component. In any case, it is clear that these emissions are small. Similarly, the excitation energy removed by these fitted nonequilibrium emissions is at most 9 MeV in the \( E/A = 9 \) MeV \(^{60}\)Ni + \(^{100}\)Mo reaction. This is out of a total of 245 MeV expected for complete fusion. These losses of excitation energy are small enough that we will ignore them from now on in the analysis of the data.

The nucleon loss by light-particle emissions was totaled, and in Fig. 7, the corresponding mean mass of the evaporation residue is indicated by the filled bar, the width of which represents the uncertainty in this determination. This result should be compared to the evaporation-residue mass distribution measured by the FMA at \( \theta_{FMA} = 5.5° \) as the angular coverage of the residues in this case is commensurate with the PPAC coverage used for the light-particle measurements. A 0.5 nucleon correction was made to account for a small difference in the mean center-of-target beam energy in the two measurements. We have assumed that the mass loss by the emission of heavier particles is negligible. For the \( E/A = 11 \) MeV reaction studied in Ref. [32], the mean mass loss from Li and Be fragments accounts for only 0.49 ± 0.06 nucleons and a smaller number is expected at this lower bombarding energy. The comparison with the mass distribution in Fig. 7 was used to select the correct mass solution in the analysis of the FMA data (Sec. II C).

B. Statistical-model calculations

In order to extract information about the level-density parameter and other properties of the compound system, statistical-model calculations were performed with the Monte Carlo computer code GEMINI [36]. The evaporation of light particles (\( Z \leq 4 \)) was calculated with the Hauser-Feshbach formalism using transmission coefficients obtained from the incoming-wave-boundary-condition model [37]. A more detailed discussion of the statistical-model calculations is described in Refs. [19,32]. As the multiplicity and spectral shape of \( \alpha \) and other heavy evaporated particles are dependent on the compound-nucleus angular momentum, it is important to constrain the distribution of this quantity associated with evaporation-residue formation. Experimentally, this distribution extends up to around the \( \ell_{max} \) values plotted in Fig. 4. In the statistical-model simulations these values are not determined by the maximum angular momentum leading to fusion (which is much larger), but by competition with statistical fission. In the calculations, the fission width was calculated using the transition-state formalism with Sierk’s fission barriers [38]. The parameter \( \alpha_f / \alpha_0 \), the ratio of level-

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### Table VII. Compound-nucleus multiplicities of light particles obtained from fitting the experimental spectra with two moving sources.

<table>
<thead>
<tr>
<th>E/A</th>
<th>Target</th>
<th>( n )</th>
<th>( n_{neq} )</th>
<th>( p )</th>
<th>( d )</th>
<th>( t )</th>
<th>( ^3)He</th>
<th>( \alpha )</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>(^{106})Mo</td>
<td>4.1 ± 0.4</td>
<td>0.03 ± 0.03</td>
<td>1.0 ± 0.1</td>
<td>2.3 ± 1.4 \times 10^{-2}</td>
<td>4.4 ± 2.2 \times 10^{-3}</td>
<td>0.85 ± 0.13</td>
<td></td>
</tr>
<tr>
<td>6</td>
<td>(^{106})Mo</td>
<td>5.1 ± 0.5</td>
<td>0.10 ± 0.07</td>
<td>1.6 ± 0.2</td>
<td>7.6 ± 1.3 \times 10^{-2}</td>
<td>1.7 ± 0.5 \times 10^{-2}</td>
<td>4.1 ± 2.1 \times 10^{-3}</td>
<td>1.06 ± 0.11</td>
</tr>
<tr>
<td>7</td>
<td>(^{106})Mo</td>
<td>6.2 ± 0.7</td>
<td>0.15 ± 0.05</td>
<td>2.2 ± 0.2</td>
<td>0.17 ± 0.03</td>
<td>4.7 ± 1.2 \times 10^{-2}</td>
<td>9.6 ± 3.5 \times 10^{-3}</td>
<td>1.35 ± 0.14</td>
</tr>
<tr>
<td>8</td>
<td>(^{106})Mo</td>
<td>7.1 ± 0.8</td>
<td>0.20 ± 0.09</td>
<td>2.6 ± 0.3</td>
<td>0.26 ± 0.05</td>
<td>7.5 ± 1.9 \times 10^{-2}</td>
<td>2.2 ± 0.8 \times 10^{-2}</td>
<td>1.64 ± 0.18</td>
</tr>
<tr>
<td>9.1</td>
<td>(^{100})Mo</td>
<td>8.2 ± 0.9</td>
<td>0.28 ± 0.13</td>
<td>2.9 ± 0.3</td>
<td>0.39 ± 0.06</td>
<td>0.13 ± 0.03</td>
<td>3.9 ± 1.1 \times 10^{-2}</td>
<td>2.1 ± 0.2</td>
</tr>
<tr>
<td>7</td>
<td>(^{208})Mo</td>
<td>2.92 ± 0.4</td>
<td>0.05 ± 0.05</td>
<td>3.9 ± 0.4</td>
<td>0.13 ± 0.04</td>
<td>1.6 ± 0.8 \times 10^{-2}</td>
<td>1.4 ± 0.8 \times 10^{-2}</td>
<td>1.4 ± 0.1</td>
</tr>
<tr>
<td>8</td>
<td>(^{208})Mo</td>
<td>3.74 ± 0.5</td>
<td>0.09 ± 0.07</td>
<td>4.3 ± 0.5</td>
<td>0.20 ± 0.08</td>
<td>3.1 ± 1.6 \times 10^{-2}</td>
<td>3.6 ± 1.9 \times 10^{-2}</td>
<td>1.7 ± 0.3</td>
</tr>
<tr>
<td>9.1</td>
<td>(^{208})Mo</td>
<td>4.65 ± 0.5</td>
<td>0.25 ± 0.1</td>
<td>4.7 ± 0.5</td>
<td>0.35 ± 0.06</td>
<td>6.9 ± 1.8 \times 10^{-2}</td>
<td>6.8 ± 1.6 \times 10^{-2}</td>
<td>2.2 ± 0.2</td>
</tr>
</tbody>
</table>
density parameters for the saddle-point and the ground-state configurations, was adjusted to \( a_T/a_n = 1.03 \) in order to fit the experimental residue cross sections. The corresponding \( \ell_{\text{max}} \) values determined from these fitted cross sections are shown by the curves in Fig. 4. Also of concern for \( \alpha \) and other heavy fragments is the bias imposed by the restricted angular range of detected residues. In order to include this bias in the calculations as well, simulated events where the residues lie outside this angular range were discarded.

Comparison of calculations to experimental spectra is made in Figs. 17–21. Before we discuss these comparisons, a few words are in order concerning the data to which the comparisons should be made. For neutrons, the nonequilibrium component is evident in the high-energy tails even at large angles. Therefore the spectra from the two most backward angles (142° and 160°) were considered most free of this component, and subsequently they were averaged to produce a single spectrum with smaller statistical errors. These averaged neutron spectra are plotted as the open square-shaped data points in Figs. 17–19. Of course these spectra do not extend to the smallest center-of-mass energies where backward emission in the laboratory frame is not possible. For these energies, we turn to the detectors at the two most forward angles (13° and 29°) and their average spectra (low energies only) are plotted as the open circular points. For charged particles, the spectra obtained from the Si telescope were considered most trustworthy as there are fewer uncertainties about the energy calibration (which is linear and particle independent) and these spectra also have no significant nonequilibrium component in the fits. However, the high-energy protons punch through this telescope, and so for these particles we also plot the averaged spectrum from the 37° and 60° CsI(Tl) detectors. The high-energy tail of these proton spectra are in the linear region of the energy calibration, so we have very high confidence in their accuracy. For protons the spectra from the Si telescope and the CsI(Tl) detectors are indicated by the circular and square solid data points in Figs. 17–19, respectively.

In order to reproduce these experimental spectra, both the Coulomb barriers and the level-density parameter had to be adjusted. Let us deal with the Coulomb barriers first. The calculations indicated in Fig. 17 by the dotted curves were obtained with a standard set of transmission coefficients and with a level-density prescription (to be discussed later) that reproduces the exponential slopes in the high-energy region. By standard set, it is meant that the Coulomb barrier is obtained using optical-model parameters derived from global fits to elastic-scattering data. It is well known that these standard calculations cannot reproduce the yield of low-energy \( \alpha \) and heavier particles in reactions like these [18,19,32,39,40]. Specifically, the predicted peak energy of the \( \alpha \)-particle spectrum is too high compared to the experimental data. This can be clearly observed in Fig. 17 for \( E/A = 5 \) and 9 MeV, and is also true for intermediate bombarding energies. To reduce the peak energy in the calculations and obtain better fits to the \( \alpha \)-particle spectra, the radius parameter associated with the nuclear potential was increased by 10% at all excitation energies. This expansion of the system decreases the height of the Coulomb barriers by \( \sim 7\% \) and, to be consistent, the adjustment was made for all particles including neutrons. The Coulomb barrier was reduced by 1.3 and 0.6 MeV for He and H isotopes, respectively. The transmission coefficients were recalculated and the predicted spectra are now indicated by the solid curves. The agreement with the experimental \( \alpha \) particle spectra at low energies is significantly better for both bombarding energies in Fig. 17; however, the changes in the predicted \( n \) and \( p \) spectra are much smaller. We will postpone a discussion of the meaning of this "expansion," however, we will continue using these new transmission coefficients for all statistical-model calculations.

For a Fermi gas, the level-density parameter relates three thermodynamical quantities, the entropy \( S \), the thermal excitation energy \( U \), and the temperature \( T \), as \( S = 2aT \), \( S = 2\sqrt{aU} \), and \( U = aT^2 \). Now if the level-density parameter is not constant, but temperature or excitation-energy dependent, then these three relations cannot be satisfied by a single parameter \( a \). Therefore it is useful to define three level-density parameters: \( a_{ST} = S/2T \), \( a_{SU} = S^2/4U \), and \( a_{UT} = U/T^2 \). In the statistical-model calculations, one requires knowledge of the level density, which is approximately proportional to

![FIG. 17. Comparison of experimental neutron, proton, and \( \alpha \)-particle spectra (data points) with statistical-model calculations (curves) for the \( E/A = 5 \) and 9 MeV \( { }^{60}\text{Ni}^+{ }^{100}\text{Mo} \) reactions. The dotted curves were obtained using transmission coefficients obtained from optical-model fits to elastic scattering data, while for the solid curves the radius parameter for the nuclear potential was increased by a factor of 1.1. The square open data points for neutrons come from the detector at the most backward angle, while the circular open points are from the detectors at the most forward angles. The circular solid data points for protons are obtained from the Si telescope while the square solid data points from the 37° and 60° CsI detectors. The \( \alpha \)-particle data are from the Si telescope.](044611-12)
\[ \exp(S), \] as a function of excitation parameters of \( A/8 \) and \( A/12 \text{ MeV}^{-1} \).

We will generally use \( a_{SU} \) in the following discussion. The temperature can then be determined from

\[ \frac{1}{T} = \frac{d\ln(r)}{dU} \simeq \frac{1}{a_{SU}} \left( 1 + \frac{U}{a_{SU}} \right) \sqrt{\frac{U}{a_{SU}}} \]  

and it characterizes the exponential slope of the emission spectra at high energies for first-chance emission. As the final spectrum is composed of contributions from multiple-chance emissions as the system cools, the apparent temperature characterizing its slope will be smaller than the initial temperature of the system.

The need for a level-density parameter that is dependent on temperature or excitation energy is apparent in Fig. 18. Calculations using two constant level-density parameters \( A/8 \) and \( A/12 \text{ MeV}^{-1} \) are compared to the \( n, p, \) and \( \alpha \) spectra measured at the lowest and highest bombarding energies. At the lowest bombarding energy \( E/A = 5 \text{ MeV} \), all three spectra are reasonably well reproduced with \( \alpha = A/8 \text{ MeV}^{-1} \), but not with \( A/12 \text{ MeV}^{-1} \). However, at the highest energy \( E/A = 9 \text{ MeV} \), the reproduction of the exponential slopes of the \( p \) and \( \alpha \) spectra is poor for \( A/8 \text{ MeV}^{-1} \). For \( \alpha \) particles at \( E/A = 9 \text{ MeV} \), the \( A/12 \text{ MeV}^{-1} \) calculation is significantly better in this regard; however, this calculation is still poor for neutrons. Note that the choice of the level-density parameter has a much greater effect on the energy at which the proton spectrum peaks than the shift induced by changing the Coulomb barrier in Fig. 17.
tation of the spectral slopes at high energy. In this work, both $K$ and $\kappa$ were varied to obtain the best fit to the $n$, $p$, and $\alpha$-particle spectra for all excitation energies and for both the targets used in this study. The value $K$, which determines the value of $a$ at low excitation energies, is most constrained from the fits to the low-energy portions of the proton and neutron spectra. This is not surprising as predominantly low-energy neutrons and protons are emitted at low excitation energies. For protons in particular, the peak energy of the kinetic-energy spectrum is very sensitive to the value of $K$ (see Fig. 18). Subsequently, the value of $\kappa$ is constrained from fitting, at all bombarding energies, the high-energy exponential tails of the $n$, $p$, and in particular, the $\alpha$-particle spectra.

The fitted values are $K = 7$ MeV and $\kappa = 1.3$ and the fitted spectra are compared to the experimental data in Fig. 19. The agreement is very good and it should be stressed that in Fig. 19, we are not just comparing spectral shapes, but absolute differential-multiplicity spectra; i.e., there is no artificial normalization between the experimental data and calculations to bring their magnitudes into agreement.

Calculated total multiplicities (curves) are also compared to the experimental data in Fig. 16. The agreement for neutrons, protons, and $\alpha$ particles is excellent. The energy dependence of the $\alpha$-particle multiplicity, at low bombarding energy, is different from other particles. In the simulations, this is due to the bias associated with detecting residues in a restricted angular range. At $E/A = 5$ MeV, this bias accounts for a 40% increase in the $\alpha$-particle multiplicity, but it is only a 1% effect at $E/A = 9$ MeV. For the rarer evaporated particles $d$, $t$, and $^3$He, the agreement is poorer, the predictions are 50% too large, at least at the higher bombarding energies where the statistical errors are small. The yield of these expensive particles is itself quite sensitive to the level-density parameter prescription, and the inability to reproduce simultaneously the slopes of the $p$ and $\alpha$-particle spectra, and the $d$ and $t$ multiplicities, was also found in the decay of A = 110 compound systems formed in incomplete-fusion reactions [17]. To reproduce the deuteron multiplicities with a constant level-density parameter, a value of $A/7$ MeV$^{-1}$ was required. A value of $A/8$ MeV$^{-1}$ works for $t$ and $^3$He particles. However, not only are these values incompatible with the $n$, $p$, and $\alpha$-particles spectra, the shapes of $d$, $t$, and $^3$He spectra also favor the fitted $a(U)$ dependence. Examples of $d$, $t$, and $^3$He spectra are displayed in Figs. 20 and 21 and compared to calculations with $a = A/8$ MeV$^{-1}$ (dashed curves) and the fitted $U$-dependent value (solid curves). Since at this point we wish to concentrate on the shapes of the spectra, in these figures we have normalized the calculated curves to reproduce the peak height of the experimental data. Although the statistical uncertainties are larger than for the $n$, $p$, and $\alpha$ data, the fitted $a(U)$ calculation is clearly a better reproduction of the $d$, $t$, and $^3$He spectral shapes. Therefore it is not believed that a high level-density parameter such as $A/7$ MeV$^{-1}$ can account for the experimental multiplicities of these rare particles and that the true explanation lies elsewhere, possibly in a preformation factor, or due to breakup of these weakly bound particles in the Coulomb field of the compound nucleus after they are emitted [19,41].
Not only is it important to decrease the Coulomb barriers to reproduce the \(\alpha\)-particle spectral shape, it is also important for the total \(\alpha\) multiplicity. The change in barrier increased the predicted multiplicity by 60\% at \(E/A=5\) MeV and by 40\% at \(E/A=9\) MeV, and allowed good reproduction with the experimental values in Fig. 16. The residue mass distribution is also quite sensitive to the \(\alpha\)-particle multiplicity as \(\alpha\) emission is very efficient at removing mass. Because of the large \(\alpha\) binding energy, it costs much less excitation energy to remove four nucleons in an \(\alpha\) particle than to emit them independently. The residue mass distribution can therefore be used as an independent check on the \(\alpha\) multiplicity. Figure 22 compares the experimental mass distributions measured with the FMA to the statistical-model calculations. The compound-nucleus masses are indicated by the arrows. The dashed curves show the predictions with the original Coulomb barriers. The average evaporative mass loss in these calculations is clearly too small. In the calculations with the reduced barriers, shown by the solid curves, the \(\alpha\) multiplicity is increased by \(-0.6\), thus shifting down the mean residue mass. The predicted means and shapes of the distributions are now in reasonably good agreement with the experimental data. For the \(^{100}\)Mo target, both the experimental and predicted mass distributions show a small downshift in average mass in going from \(\theta_{\text{FMA}}=0^\circ\) to \(5.5^\circ\). Due to the larger recoil kicks from \(\alpha\) particles, residues emitted at larger angles have larger \(\alpha\) multiplicities, on average, and thus the mean mass loss is larger.

The striking dip in the experimental mass at \(A=140\) for \(\theta_{\text{FMA}}=0^\circ\) in the reaction on the \(^{100}\)Mo target (shown most clearly in Fig. 7) is not seen in the statistical-model calculations and may not be real. We suspect this dip is caused by isomers that decay by electron conversion during their transit through the FMA, changing their charge state and thus no longer reaching the focal plane. The \(^{140}\)Sm system has 6.2 and 22.3 ns isomers [42] and may be responsible for the dip as 26\% and 68\%, respectively, of the population of these isomers decay after passing through the reset foil in front of FMA. Both isomers have significant probabilities for decay via the conversion electron process.

As \(\alpha\) and heavier fragments provide the greatest recoil kick to the residues, they are also important for understanding the widths of the residue velocity and angular distributions. The curves in Figs. 1–3 show the statistical-model predictions (with reduced barriers) that have been normalized in height to the same magnitude as the experimental data. The predicted widths of the velocity and angular distributions are clearly in accord with the experimental data. As in Ref. [32], the statistical-model calculations also considered the evaporation of clusters with \(A>4\), including stable and unstable He, Li, and Be fragments. At the highest excitation energy, the total multiplicity of these particles obtained in the simulations is 0.41, however, 87\% of these are unstable and decay to lighter fragments and thus are already included in the predicted multiplicities shown in Fig. 16. The inclusion of these heavier clusters makes small changes to the statistical-model predictions, but it does lead to somewhat better agreement with the data. The \(\alpha\) multiplicity is increased by 16\%, the widths of the residue velocity distributions is increased by 9\%, and the total evaporative mass loss is increased by 1.0 nucleon. The secondary \(\alpha\) particles from the decay of some of these fragments are expected to give rise to a small shoulder in the \(\alpha\) spectrum at extreme sub-barrier energies [32]. Some indication of this feature can be seen in Fig. 15, especially for the 19° and 39° spectra at energies below 25 and 20 MeV, respectively.

The temperature as a function of excitation energy for \(A=160\) derived from Eq. (11) is shown in Fig. 23. The temperature increases faster with excitation energy for \(\alpha(U)\) than for either of the two constant parameters, \(A/8\) or \(A/12\) MeV\(^{-1}\). At \(U=70\) MeV, the temperature is comparable to the \(A/8\) MeV\(^{-1}\) result, while at \(U=220\) MeV, the temperature is now equal to the \(A/12\) MeV\(^{-1}\) value. This parallels the change in constant level-density parameter in Fig. 18 required to fit the \(\alpha\)-particle spectra as these two excitation energies are typical of the values at which these particles are emitted at \(E/A=5\) and 9 MeV. The other level-density parameters \(a_{ST}\) and \(a_{UT}\) have stronger excitation energy dependences as the temperature also depends on the derivative of \(a_{SU}\) as well as its absolute value [Eq. (11)]. These two parameters can be approximately described by \(\kappa=3.0\) and \(\kappa=5.4\), respectively. The fitted \(\alpha(U)\) dependence is in reasonable agreement with the calculations of Shlomo and Natowitz [43] for \(A=160\). The values of \(a_{UT}\) obtained at \(T=1, 2, 3, 4,\) and 5 MeV from Fig. 1 in Ref. [43], were used to determine the solid data points plotted in Fig. 23. For \(T\leq4\) MeV, the \(a_{SE}\) obtained from these calculations of
excitation energies. However, as one moves away from the attractor line on the neutron-rich (proton-rich) side, the neutron (proton) decay width quickly dominates the other at low excitation energies, and thus the decaying systems evaporate towards the line. At higher excitation energies, the evaporation of nucleons and clusters remove similar number of neutrons and protons and the initial trajectories of the decaying system in the chart of nuclides is approximately parallel to the attractor line. Movement towards the attractor line only occurs at the end of the decay sequence when the system is cooler. The compound nuclei $^{152}$Yb and $^{160}$Yb formed with the $^{92}$Mo and $^{100}$Mo targets, respectively, lie on either side of the expected position of the attractor line. Thus, the location of the residues should approach the line from either side, allowing us to constrain its location. The average $N$ and $Z$ of evaporation residues formed in the reactions studied in this work were deduced from the multiplicities of the light particles (both compound-nucleus and nonequilibrium sources). Their locations on the chart of nuclides are displayed as the data points in Fig. 24. A data point for the $E/A = 4.3$ MeV $^{58}$Ni + $^{94}$Mo reaction (same $^{152}$Yb compound system as the $^{60}$Ni + $^{92}$Mo reaction) is also plotted. This point was determined from the intensity of evaporation-residue channels measured by Nolte et al. using $\gamma$-ray spectroscopy [44]. The mean locations of the evaporation residues from the two compound nuclei converge towards a region of similar $n$-$p$ asymmetry consistent with the indicated attractor line. The location of the attractor line shown in this plot was calculated using the statistical-model parameters deduced in Sec. III B. Compared to Ref. [21], its location in this mass region has shifted by 0.5 of a neutron along the $N$ axis to the neutron-rich side. This shift is mostly due to changes in the proton Coulomb barriers between the two calculations, but experimentally one cannot differentiate between the two. The solid curves in Fig. 24 are the loci of the predicted, average locations of the evaporation residues with increasing excitation energy for the two compound nuclei. Again, these are consistent with the experimental data. From the indicated experimental errors bars in Fig. 24, the true location of the attractor line can be shifted from the predicted value by at most $\pm 2$ units along the $N$ axis.

Hanold et al. have measured the $N$-$Z$ distributions of evaporation residues formed in $E/A = 26–50$ MeV $^{129}$Xe + Be, C, and Al incomplete fusion reactions [23]. The average $A$ for each residue $Z$ was found to be 2 mass units more neutron deficient than expected in GEMINI simulations starting with distributions of compound nuclei predicted by various incomplete-fusion models. If the reduced Coulomb barriers from this work were utilized, then the difference would be even larger. In Ref. [21] it was pointed out that the mean locations of the evaporation residues and the assumed compound systems in the incomplete-fusion models, lie on opposite sides of the evaporation attractor line. The residues lie on the proton-rich side, while the compound nuclei lie on the neutron-rich side. As, on average, residues approach but do not cross, the line in statistical-model simulations, then this suggests we must find an explanation to get both the residues and the compound systems on the same side. Maybe the actual location of the attractor line is on the proton-rich side.

**FIG. 23.** Temperature $T$ as function of thermal excitation energy $U$ for $A = 160$. The dashed curves are associated with constant level-density parameters of $A/8$ and $A/12$ MeV$^{-1}$. The solid curve is obtained from the fitted excitation-energy dependent level-density parameter. The solid data points are theoretical predictions from Shlomo and Natowitz [43]. The open data point is the experimental result of Gonin et al. [16].

Shlomo and Natowitz can be well reproduced using Eq. (12) with $K = 8.2$ MeV and $\kappa = 0.9$, but this fit deviates at higher temperatures. Although we disagree with the absolute light-particle multiplicities measured by Gonin et al. [16], the temperature they extracted by spectra subtraction for $E^* = 236$ MeV (indicated by the open data point in Fig. 23) is in agreement with this work.

Fineman et al. [18] and Caraley et al. [20] have fitted $\alpha$-particle and proton data in coincidence with heavier compound systems using the same form for $a_{SU}(U)$ [Eq. (12)]. Due to the smaller range of excitation energies investigated in these studies and the absence of neutron data, the excitation-energy dependence of $a$ could not be extracted. For example, with the $A = 200$ systems studied in Ref. [20], the $\alpha$-particle data could be reproduced with $K = 12$ MeV, $\kappa = 0$, and also with $K = 8.2$ MeV, $\kappa = 2–3$. In Ref. [18] for the $A = 224$ compound system, reproduction of proton and $\alpha$-particle data was obtained with both $K = 15$ MeV, $\kappa = 0$, and with $K = 8.2$ MeV, $\kappa = 8.5$. Although there is uncertainty as to the excitation-energy dependence of $a_{SU}$, it is clear that the absolute value of $a_{SU}/A$ is smaller for these heavier systems. Also, if we assume that the value of $K$ should be near 8 MeV [5] then this would imply a significant increase of $\kappa$ from the value of 1.3 obtained in the present investigation of $A \approx 160$ systems to the larger values in these two studies of heavier systems.

**C. Evaporation attractor line and the $n$-$p$ asymmetry dependence of level-density parameter**

The evaporation attractor line is a line on the chart of nuclides towards which an evaporating system moves, on average, as it cools [21,22]. At the attractor line, neutron and proton partial decay widths are approximately equal at all
level densities at their respective Fermi energies. In Thomas-Fermi calculations of $N$ fermions of mass $m$ in a box of volume $V$, 

$$g \propto V^{2/3} N^{1/3} m,$$  

and thus for a nucleus of $N$ neutrons and $Z$ protons we would expect

$$a \propto A^{2/3}(N^{1/3} m_n + Z^{1/3} m_p).$$  

If we take $m_n = m_p$ and expand about $N - Z = 0$, one obtains

$$a \propto A \left[ 1 - \frac{1}{9} \left( \frac{N-Z}{A} \right)^2 \right] = A \left[ 1 - \frac{4}{9} \left( \frac{a}{A} \right)^2 \right].$$  

This dependence on isospin projection $t_3 = (Z-N)/2$ is small, and when implemented in statistical-model simulations, the effect could not be discerned within the statistical uncertainties of the Monte Carlo simulations. Larger dependences may be obtained if the proton and neutron effective masses are different [45], i.e., $m_n \neq m_p$. Also at large excitation energies, or for systems close to a drip line, the treatment of the continuum region of the single-particle level density becomes important and this may lead to a $n\cdot p$ asymmetry dependence when the proton and neutron separation energies are different.

Recently, Al-Quraishi et al. [24] have suggested two dependences based on fitting low-energy level-density data for 20$\leq A \leq 70$. In their case $B$, justified from the counting of nuclear levels of different isospin,

$$a_B = \frac{A}{9.009 \text{ MeV}} \exp[6.41 \times 10^{-4} (N-Z)^2].$$  

For their case $C$, associated with continuum corrections,

$$a_C = \frac{A}{8.787 \text{ MeV}} \exp[4.93 \times 10^{-2} (Z-Z_\beta(A))^2].$$  

where $Z_\beta(A)$ is the $Z$ of the $\beta$-stable isotope of mass $A$. While an extrapolation of this fit to $A \approx 160$ and high excitation energies and spins is questionable, these two dependences nevertheless serve as useful examples of the effect of different $n\cdot p$ asymmetry dependences. The variation of $a_B$ and $a_C$ with neutron number is plotted in Fig. 25 for $A = 160$ isobars. For the $^{160}$Yb compound system, both level-density parameters show large dependences on $N$, especially $a_C$. The slopes of the $N$ dependences are opposite in sign for the two cases. Statistical-model calculations incorporating these two level-density parameters were performed and the loci of the average location of the evaporation residues as a function of bombarding energy are plotted as the solid curves in Fig. 26. Compared to the calculations of Fig. 24, the mean locations of the residues have been shifted towards the regions where the level-density parameter has its maximum; $N=Z$ for case $B$ and the $\beta$ valley of stability for case $C$.

To understand what is happening in the statistical-model simulations, consider the relative emission rates for neutrons and protons. In this case the two daughter nuclei have the

FIG. 24. This figure illustrates a portion of the chart of nuclides showing the location of the evaporation attractor line (dotted line) relative to the stable nuclei (filled squares), the limits of known nuclei, and the $N=Z$ line. The location of the two compound nuclei ($^{152}$Yb and $^{160}$Yb) studied in this work are indicated. The mean location of the evaporation residues deduced from the experimental multiplicities of light particles are indicated by the data points; open-square points for $^{152}$Yb and solid circular points for $^{160}$Yb. The crossed-squared datum point was obtained by Nolte et al. [44] and it also is associated with the $^{152}$Yb compound nucleus. The smooth solid curves are statistical-model predictions of the mean location of the residues with increasing excitation energy.

An alternative approach might be to consider a possible dependence of the level-density parameter on the neutron-proton asymmetry. Only an $A$ dependence of the level-density parameter is generally assumed in statistical-model calculations, however, some dependence on $N$ and $Z$ should be expected. For instance, in a two-component Fermi-gas model, the level-density parameter is given by

$$a = \frac{\pi^2}{6} (g_n + g_p),$$  

where $g_n$ and $g_p$ are the neutron and proton single-particle
same $A$, but $N$ differs by one unit between them. The level-density parameters of the daughter following neutron and proton emission can be written as

$$a_n = \bar{a} \left( 1 + \frac{1}{2\bar{a}} \frac{\partial a}{\partial t_3} \right), \quad a_p = \bar{a} \left( 1 - \frac{1}{2\bar{a}} \frac{\partial a}{\partial t_3} \right),$$

where $\bar{a}$ is the average of $a_n$ and $a_p$ and the derivative is for constant $A$. As the mass and spin of $n$ and $p$ are similar, the angular-momentum effects on the statistical decay widths will also be similar for both particles. The main difference in decay widths will come from the ratio of relative level densities of the daughter nuclei, which can be written approximately as

$$\frac{\Gamma_n}{\Gamma_p} \approx \exp\left[ 2 \sqrt{a_n(U - E_n^c) - 2\sqrt{a_p(U - E_p^c)}} \right]$$

where $E_n^c$ is the energy cost of emitting the particle. This cost is the separation energy for neutrons and, for protons, it also includes the Coulomb barrier. If $\partial a/\partial t_3$ is small, this can be expanded as

$$\frac{\Gamma_n}{\Gamma_p} = \exp\left[ 2 \sqrt{a(U - E_n^c)} \right] 
- 2\sqrt{a(U - E_p^c)} \sqrt{\frac{U - \langle E^c \rangle}{\bar{a}}} \right],$$

where $\langle E^c \rangle$ is the average of the two costs. The first exponential factor is the standard expression for the case where $\partial a/\partial t_3 = 0$, and thus the second factor contains the isospin-projection or asymmetry dependence. This second asymmetry factor is largest at high excitation energies. For the $^{160}$Yb compound system with case $B$ in Fig. 26, the asymmetry factor causes the decaying system to move towards the $N = Z$ line at high excitation energies, crossing the attractor line. As the system cools, the asymmetry factor becomes less important and, on average, the system moves back towards the attractor line. Similarly for case $C$, the asymmetry effect resulted in a crossing of the attractor line for the $^{152}$Yb compound system. Clearly, an $n,p$ asymmetry dependence can cause the decaying system to cross the attractor line if the initial excitation energy is high enough. The asymmetry dependence in case $C$ is extreme in this mass region and the predicted residues in Fig. 26 are much too neutron rich compared with experiment, while the curves for case $B$ just pass close to the farthest edge of the experimental error bars.

Besides the location of the residues, the multiplicity ratio $t^3$He is another very useful quantity to consider as it is quite sensitive to asymmetry effects at high excitation energies. These two costly particles are only emitted during the first few evaporation steps before the systems cool significantly. For first-chance emission, Eq. (21) can also be used for the $t^3$He ratio as, like $n$ and $p$ emissions, the daughters have the same $A$, but differ by one unit in $N$. However, when multiple-chance emission is considered, a somewhat smaller effect is found in statistical-model simulations. The asymmetry dependence of the level-density parameter may itself be temperature dependent and thus the $t^3$He ratio will probe the high-temperature asymmetry dependence directly. The experimental ratio is plotted as a function of bombarding en-

FIG. 25. Dependence of level-density parameter $a$ on neutron number $N$ for $A = 160$. The dotted lines are for constant values of $A/8$ and $A/12$ MeV$^{-1}$. The solid and dashed curves correspond to cases $B$ and $C$ from Al-Quraishi et al. [24], which peak at $N = Z$ and the $\beta$ valley of stability, respectively. The neutron number of the $^{160}$Yb compound nucleus is indicated by an arrow.

FIG. 26. Portions of the chart of nuclides as in Fig. 24, but now the smooth curves indicate the loci of the predicted mean residue position with increasing excitation energy, which were calculated with the asymmetry-dependent level-density parameters; cases $C$ and $B$ from Ref. [24].
energy in Fig. 27 and compared to predictions with the two cases (dashed curves) and with the excitation-energy dependent, but asymmetry independent, parameterization from Sec. III B (thick solid curve). Although this latter parametrization overpredicts both the absolute $t$ and $^3$He multiplicities, it reproduces the ratio quite well. In contrast, case B is roughly a factor of 5 too large and case C is over a factor of 100 too small. Clearly, the asymmetry dependence at high excitation energies is considerably smaller than either case B or C. To constrain the magnitude of the asymmetry dependence, calculations were performed using the fitted temperature-dependent parametrization and adding a first-order isospin-projection correction

$$a_{SU}(U) = \frac{A}{K + \kappa \frac{U}{A}} \left[ 1 + \frac{\partial a}{\partial t_3}(t_3 - t_3^{CN}) \right],$$

where $t_3^{CN}$ is the compound nucleus value of $t_3$. As $t$ and $^3$He particles are emitted in the first few decay steps, we are only probing the value of $\partial a/\partial t_3$ in the immediate vicinity of the compound nucleus. The solid curves in Fig. 27 indicate the predicted multiplicity ratio for $\partial a/\partial t_3 = \pm 0.15$ and $\pm 0.3 \text{ MeV}^{-1}$. For all four data points which probe the excitation-energy range $100 < U < 240 \text{ MeV}$, we can restrict $|\partial a/\partial t_3| < 0.2$. For comparison, $\partial a/\partial t_3 = 0.7 \text{ MeV}^{-1}$ in case B. At much lower excitation energies where the asymmetry dependence of the level density would be of interest for $r$ and $rp$ processes [24], we have little sensitivity. As far as the data of Hanold et al. is concerned, an asymmetry dependence at low excitation energy will not cause one to cross the attractor line, thus we still do not have a complete understanding of these data.

### D. Coulomb barriers

In Sec. III B the Coulomb barrier for $\alpha$-particle emission was reduced from the value deduced in optical-model fits to elastic-scattering data in order to reproduce the peak energies in the $\alpha$-particle kinetic-energy spectra. This reduction was achieved by considering an expansion of the compound nucleus; increasing its radius by 10%. This is not the only means by which a fit to the $\alpha$-particle peak energy could be accomplished. For example, in Ref. [32], the experimental peak energies for He, Li, and Be fragments produced in the $E/A=11 \text{ MeV} \ 60\text{Ni}+\text{Mo}$ reaction were fit by assuming the compound nucleus had a large prolate deformation (ratio of major to minor axes of 1.6). Apart from expansion and deformation, one can also reduce the Coulomb barrier and fit the data by increasing the surface diffuseness by 30% or increasing both the radius and diffuseness by 7%. It is not clear whether one, or some combination, of these mechanisms is responsible. In elastic scattering, one is probing the Coulomb barrier for cold ground-state systems located in the $\beta$ valley of stability. In this work, the Coulomb barrier is associated with emission from hot, rapidly rotating systems located far from the valley of stability on the protons-rich side. The roles of excitation energy, angular momentum, and proton richness can be considered.

Hartree-Fock-Bogoliubov calculations [46] suggest that the surface diffuseness for the systems of interest are not expected to be significantly different from $\beta$-stable values and the radii scale approximately as $A^{1/3}$ as assumed in the optical-model fits. The surface diffuseness and, to a smaller extent, the radius are expected to increase with increasing temperature [47–49]. The Coulomb barrier would then decrease with increasing temperature [50]. However, the data in this work were reproduced with a constant temperature-independent reduced barrier, and even at the highest temperature obtained in these reactions, the predicted decrease in barrier height is too small [50].

Deformed compound nuclei can arise through thermal shape fluctuations giving rise to a distribution of shapes, which is dependent on the excitation energy and the angular momentum of the system. At low excitation energies $E^* \approx 100 \text{ MeV}$, statistical-model calculations incorporating these fluctuations were shown to account approximately for the magnitude of the barrier reduction, but at higher excitation energies the predicted reduction was insufficient to fit the experimental data [51,52]. Larger nonthermal deformations induced by the fusion dynamics might be able to account for the barrier reduction, however, in that case we should expect to see the reduction dependent on the entrance channel. Experimentally, no such dependence has been confirmed in this mass region [39,40]. Possibly, some combination of thermal-shape fluctuations at low excitation energies and thermal expansions and increased surface diffuseness at higher temperatures may be the correct explanation.

No matter which of these scenarios, or which combination of them, is the correct answer, there are consequences for the...
level density. In the Fermi-gas model, the density (represented as $\rho$) dependence of the level-density parameter is $a \propto \rho^{-2/3}$. Therefore by expanding the nucleus, $a$ increases. One also expects increases in $a$ due to deformation or increased surface diffuseness as in both cases there is an increasing amount of low-density surface material. However, at the same time there is an energy cost associated with expanding, deforming, or increasing the diffuseness of the system, which decreases the level density. If the effect is thermally driven, then the system will only expand, deform, etc. if the entropy increases and thus, on net, the level density must increase due to these effects. The decrease in the level-density parameter with excitation energy extracted in Sec. III B must therefore be judged relative to an enhanced value expected from these effects. If corrected for these effects, the true decrease with excitation energy may be larger.

IV. CONCLUSIONS

The properties of evaporation residues and the accompanying light particles have been studied in the $E/A = 5-9$ MeV $^{60}$Ni+$^{106}$Mo, $^{92}$Mo reactions. Kinetic-energy spectra of evaporated $n$, $p$, $d$, $t$, $^3$He, and $\alpha$ particles and the mass, velocity, and angular distributions of evaporation residues have been measured. These data indicate that the residues are formed in essentially complete-fusion reactions with very little nonequilibrium emissions. The kinetic-energy spectra of $n$, $p$, and $\alpha$ particles cannot be reproduced in statistical-model calculations with the standard prescription for the level-density parameter with no excitation-energy dependence. A dependence of the form $a = A/(7 + 1.3 \times U/A)$ MeV$^{-1}$ was found to reproduce the exponential slopes of all the light-particle spectra simultaneously. This is consistent, in the excitation region studied, with the predictions of Shlomo and Natowitz [43]. In order to fit the spectra, it was also necessary to reduce the Coulomb barriers for charged particle emission by 7%. With these ingredients the calculations were also able to reproduce the measured mass, velocity, and angular distributions of the residues and the light-particle multiplicities as functions of bombarding energy. The biggest disagreement was a 50% overprediction of the multiplicities of the rarer particles: $d$, $t$, and $^3$He.

From the evaporative multiplicities, the mean location of the evaporation residues in the chart of nuclides was deduced and found to be consistent with the evaporation attractor line of Ref. [51]. It was shown that if the level-density parameter has a $n-p$ asymmetry or isospin-projection dependence, then this could lead to the production of more neutron-rich or proton-rich evaporation residues. Also an asymmetry dependence at large excitation energies will modify the ratio of tritons to $^3$He fragments evaporated. From the experimental $t/^3$He ratio, the magnitude of any asymmetry dependence was tightly constrained in the excitation-energy window $100 < E^* < 245$ MeV.

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