ON RESONANCES OF THE $^{24}Mg+^{24}Mg$ SYSTEM ☆

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Resonances in heavy-ion scattering of $^{24}Mg+^{24}Mg$ are studied by the use of a new molecular model, in which collective motions of the system are described in the rotating molecular frame of the di-nuclear system. A stable configuration of the system is found to be a pole–pole one, due to the prolate shape of the ^{24}Mg nuclei. Normal modes at the equilibrium are solved and various molecular levels are obtained, which appear to be in good correspondence with experiment.

1. Introduction

Many isolated resonances have been observed in $^{24}Mg+^{24}Mg$, $^{24}Mg+^{28}Si$ and $^{28}Si+^{28}Si$ scattering in a CM energy around 50 MeV [1–5]. The number of the resonance levels is much larger than the sequence of grazing partial waves. It is therefore expected that not only the relative motion but also many other degrees of freedom participate in formation of these resonance states. The resonances are surprisingly narrow and are observed correlatingly among the elastic and inelastic decay channels. Hence they are considered to be eigenstates of the whole compound system. Recent spin-alignment measurements gave an assignment of high spin, $36\hbar$, for the resonance at a CM energy equal to 45.7 MeV in the $^{24}Mg+^{24}Mg$ system [4]. Strong enhancements are observed in decays up to the 6^+ channel and partial decay widths to these channels amount to approximately 30% of the total width of about 200 keV [3]. It is also noted that fractions of α-transfer channels are an order of magnitude smaller than those of symmetric-mass decay [3] which suggests a di-nucleus configuration in origin of the resonances.

☆ Preliminary results were reported in ref. [1].

The stability of di-nucleus configuration was once studied by Broglia et al. with a macroscopic liquid-drop model [6]. They obtained a stable configuration, for example, for the $^{26}Mg+^{26}Mg$ system with high spins, $40\hbar$–$50\hbar$. For each total spin J, a single state is associated. On the other hand Nilsson–Strutinsky calculations show secondary minima at high spins and at very large deformation, the structure of which appears to correspond to a di-nucleus configuration [7]. Recently Maass and Scheid investigated the stability of a similar configuration with the two-center-shell model, taking into account the liquid-drop energy and shell correction [8]. They also obtained stable di-nucleus configurations for several high spins. Furthermore, they partially considered vibrational modes around the equilibrium. The experiments, however, indicate much dense levels compared with the sequence of grazing angular momenta as stated above. Hence we have to investigate the dynamics of the di-nuclear system, including many degrees of freedom in addition to the relative motion of two nuclei.

As a first step, we take up a system of two identical deformed nuclei with axial symmetry. Collective degrees of freedom of each deformed nucleus are the orientation of the symmetry axis, which is described

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by the Euler angles (α, β). We have to solve the dynamics of these degrees of freedom and of the relative motion. We introduce a rotating molecular frame, the z'-axis of which is parallel to the relative vector of two interacting nuclei. All the degrees of freedom of the system are transformed into the rotational motion of the total system and the internal motions referred to the rotating frame. Thus the Euler angles (α_i, β_i) of each constituent deformed nucleus are no more rotational variables, but are a kind of internal collective variables, which are solved by the method of normal modes around the equilibrium configuration. The new collective modes consisting of the above internal variables give rise to many excited states, which are expected to be responsible for the many sharp resonances observed. The purpose of the present paper is to propose a new molecular model for interacting di-nuclear systems and to clarify the mechanisms of the heavy-ion resonances observed in the heavier systems.

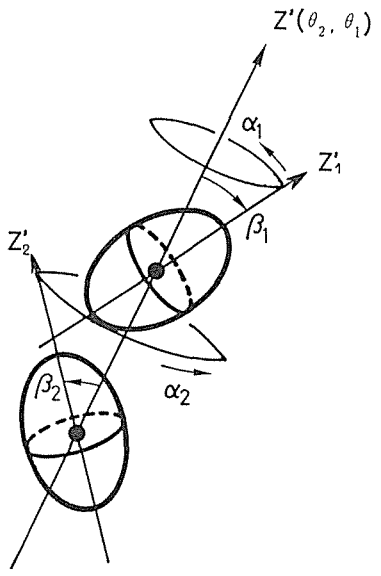

Fig. 1. The coordinates in the rotating molecular frame.

2. Formulation

First, we define a coordinate system of the molecular model and give the kinetic energy operators in the coordinate system. We define the rotating molecular axis z' of the whole system straightforwardly with the direction of the relative vector between two nuclei; the directions of the intrinsic axes of each deformed nucleus are defined referring to the molecular frame. For simplicity, constant deformation of the nuclei is assumed. We thus start with seven degrees of freedom as illustrated in fig. 1, that is, the relative vector $\boldsymbol{R}=(R, \theta_2, \theta_1)$ and the Euler angles of the interacting nuclei (α_1, β_1) and (α_2, β_2). The variables α_1 and α_2 are reconstructed into variables $\theta_3=(\alpha_1+\alpha_2)/2$ and $\alpha=(\alpha_1-\alpha_2)/2$. Then we have $(q_i)=(\theta_1, \theta_2, \theta_3, \alpha, R, \beta_1, \beta_2)$, where the θ_i are the Euler angles of the molecular frame with four other internal variables.

Hess et al. also investigated the molecular spectrum of a di-nuclear system in connection with anomalous $e^{\pm}$ production in $^{238}U+^{238}U$ collisions [9]. They however restricted themselves to certain configurations such as the pole–pole one. In the present study we performed a quantization with full degrees of freedom of the system.

In the laboratory frame the kinetic energy of the system consists of the relative motion between the two nuclei and the rotational energies of the constituent nuclei. It can be easily transformed into the rotational energy of the total system and the kinetic energies of internal motions referred to the molecular frame. After expressing three sets of the angular velocities with the time derivatives of the corresponding Euler angles, we obtain a classical kinetic energy expression $\frac{1}{2}\sum g_{ij}\dot{q}_i\dot{q}_j$, and then quantize the system by the usual prescription.

The kinetic energy operator obtained can be divided into three parts, that is, a rotational operator $\frac{1}{2}\hbar^2 \cdot \sum \mu_{ij}\hat{J}'_i\hat{J}'_j$ for the Euler angles θ_i, operators for the internal variables, and Coriolis coupling terms. As is shown later, the dependence of an interaction potential on the α-degree of freedom is fairly weak and thus we consider the motion related to this variable as *an internal rotation*. Accordingly, we obtain T_{rot} for the Euler angles θ_i and the variable α. The remaining parts are the operators T_{vib} and T_{C} for the internal variables R, β_1, β_2 and the Coriolis coupling terms. Since the kinetic energy operators are not simple yet, the nondiagonal parts $\hat{J}'_i\hat{J}'_j$ of T_{rot} with *small* coefficients are rearranged into T_{C}, that is, cross terms between $\hat{J}'_x$, $\hat{J}'_y$ and $\hat{J}'_z$, $\hat{J}'_\alpha$ are moved to T_{C} but the term $\hat{J}'_z \cdot \hat{J}'_\alpha$ is kept in T_{rot}. Then, $T=T'+T'_{\text{C}}$ with $T'=$

$T'_{\rm vib}+T'_{\rm rot}$. In the following, we restrict our discussion to rotation and vibration terms expressed by T'. The effects of the remaining terms $T'_{\rm C}$ will be discussed elsewhere. Consequently, with the coordinate system, we introduce a rotation–vibration type wave function as basis one,

$$\Psi_\lambda \sim D^J_{MK}(\theta_i)\exp(\mathrm{i}\nu\alpha)\chi_K(R,\beta_1,\beta_2)\,. \tag{1}$$

The kinetic energy operator T' is again divided into three parts, $T'=T'_{\rm vib}+T'_{\rm rot}=T^0_{\rm vib}+T^0_{\rm rot}(J,K,\nu)+T_{\rm mode\text{-}mode}$ with

$$T^0_{\rm vib}=-\frac{\hbar^2}{2}\left[\frac{1}{\mu}\frac{\partial^2}{\partial R^2}+\left(\frac{1}{I}+\frac{1}{\mu R^2}\right)\left(\frac{\partial^2}{\partial\beta_1^2}+\frac{\partial^2}{\partial\beta_2^2}\right)\right], \tag{2}$$

$$T^0_{\rm rot}(J,K,\nu)=\frac{\hbar^2}{2}\left[\frac{J(J+1)-\frac{3}{2}K^2-\frac{1}{2}\nu^2}{\mu R^2}+\left(\frac{1}{I}+\frac{1}{\mu R^2}\right)\left(\frac{(K+\nu)^2-1}{4\sin^2\beta_1}+\frac{(K-\nu)^2-1}{4\sin^2\beta_2}-\frac{1}{2}\right)\right], \tag{3}$$

$$T_{\rm mode\text{-}mode}=\frac{\hbar^2}{2\mu R^2}\left[[\exp(2\mathrm{i}\alpha)+\exp(-2\mathrm{i}\alpha)]\times\left(-\frac{\partial^2}{\partial\beta_1\,\partial\beta_2}+\tfrac{1}{4}(\hat J'^2_z+1)\cot\beta_1\cot\beta_2\right)-\tfrac{1}{4}\hat J'_\alpha[\exp(2\mathrm{i}\alpha)+\exp(-2\mathrm{i}\alpha)]\hat J'_\alpha\cot\beta_1\cot\beta_2\right], \tag{4}$$

where μ denotes the reduced mass of the two nuclei and I denotes the moment of inertia of a single constituent nucleus. In the expression $T^0_{\rm rot}(J,K,\nu)$ of eq. (3), we use the eigenvalues K and ν instead of $\hat J'_z$ and $\hat J'_\alpha$. Terms corresponding to the additional potential in the nuclear collective model are included in $T^0_{\rm rot}(J,K,\nu)$ and $T_{\rm mode\text{-}mode}$ for convenience. Accordingly, the wave function $\chi_K(R,\beta_1,\beta_2)$ is defined so as to contain the square root of the volume element $R\sqrt{\sin\beta_1\sin\beta_2}$.

The potential energy $V_{\rm int}$ between two deformed nuclei is calculated by the double-folding model [10,11]. We expand the density of each deformed nucleus in the body-fixed frame as $\rho^{\rm B}(r)=\sum\rho_{lm}(r)Y^*_{lm}(\hat r)$, with $\rho(r)=\sum\rho_{lm}(r)D^l_{\nu m}Y^*_{l\nu}(\hat r)$ in the laboratory frame. $V_{\rm int}$ is expressed in the molecular frame as

$$V_{\rm int}(R,\beta_1,\beta_2,\alpha)=\sum_{l'l''l}(2\pi)^{-3}\mathrm{i}^{l'-l''-l}\hat l'\hat l''(l'l''00|l0)\times F_{l'l''l}(R)G_{l'l''l}(\beta_1,\beta_2,\alpha)\,, \tag{5}$$

$$F_{l'l''l}(R)=\int \mathrm{d}q\,q^2 j_l(qR)\tilde v(q)\tilde\rho_{l'}(q)\tilde\rho_{l''}(q)\,, \tag{6}$$

$$G_{l'l''l}(\beta_1,\beta_2,\alpha)=\sum_m(l'l''m-m|l0)d^{l'}_{m,0}(\beta_1)d^{l''}_{-m,0}(\beta_2)\times\exp(2\mathrm{i}m\alpha)\,, \tag{7}$$

where $\tilde\rho_l(q)$ and $\tilde v(q)$ denote the Fourier transforms of the nuclear density and of the nucleon–nucleon interaction, and $\hat l=\sqrt{(2l+1)}$. A density profile of $\rho^{\rm B}(r)$ is assumed to be a Fermi distribution, and the radius of the deformed nucleus is taken to be $R_{\rm N}=R_0(1+\beta Y_{20})$, R_0 being $1.1\,A_i^{1/3}$ fm. The deformation parameter β of ^{24}Mg is determined to be 0.535 with the $B(\mathrm{E2})$ value of the ground-state band. The diffuseness parameter $a_{\rm N}$ is adjusted to be 0.385 fm by reproducing the RMS radius of the ground state. The M3Y potential is used as a folded nucleon–nucleon interaction with a knock-on exchange term of strength -262 MeV fm^3 [11].

It should be noted that the simple folding-model potential is not accurate in the fully overlapping region. We must take into account the repulsive effect by the Pauli principle, or the density overlap going beyond the normal density. In addition to the potential $V_{\rm int}$, we introduce a repulsive potential due to the higher overlapping density, which is estimated by the binding energy loss of nuclear matter, i.e., by the equation of state. The energy loss per unit volume (fm^3) is taken to be $\rho_0^2\cdot 500$ MeV fm^3 for twice the normal density ρ_0. The overlapping volume with twice the normal density is simply calculated by a similar expression to eq. (5), but with the δ-function instead of the interaction, and by introducing the density distribution of constituent nuclei with a very small diffuseness parameter $a_{\rm P}$ of 0.1 fm.

3. Internal collective modes

In order to solve the normal modes for three variables R, β_1 and β_2, $V_{\rm int}$ is expanded into a quadratic form at an equilibrium configuration. The energy surface of $V_{\rm int}(R, \beta, \beta, 0)$ is displayed in fig. 2a, to which we add the dominant centrifugal potential $J(J+1)\hbar^2/2\mu R^2$ to give the correct radial dependence of the energy surface with spin J being 36. We find a local minimum point at $\beta_1=\beta_2=0$ and $R=8.24$ fm, namely at the pole–pole configuration. In fig. 2b, the α-dependence is shown with $\beta_1=\beta_2=\beta$, and in fig. 2c, the K- and ν-dependence of the β_i-energy curves of $T^0_{\rm rot}+V_{\rm int}(R, \beta, \beta, 0)$. The α-dependence is seen to be very weak with small values of β, which justifies the present choice of the *internal-rotational mode* for the variable α. By expanding $V_{\rm int}+J(J+1)\hbar^2/2\mu R^2$ at the equilibrium point we obtain $V_{\rm Q}=V_0+V_{\rm mode-mode}$ with $V_0=E_0+J(J+1)\hbar^2/2\mu R_{\rm e}^2+\frac{1}{2}k_R(R-R_{\rm e})^2+\frac{1}{2}k_\beta(\beta_1^2+\beta_2^2)$ and $V_{\rm mode-mode}=V_\alpha(R_{\rm e})\beta_1\beta_2\cos 2\alpha$, where only the $m=0$, 1 terms in eq. (7) contribute up to the second order. $V_{\rm mode-mode}$ represents an α-dependence of $V_{\rm int}$ and gives rise to mode–mode couplings in addition to $T_{\rm mode-mode}$ from the kinetic energy T'. The total hamiltonian is written as follows:

$$H=H_0+H_{\rm mode-mode}+T'_{\rm C}\,, \tag{8}$$

$$H_0=T^0_{\rm vib}+T^0_{\rm rot}+V_0\,, \tag{9}$$

$$H_{\rm mode-mode}=T_{\rm mode-mode}+V_{\rm mode-mode}\,. \tag{10}$$

In $T^0_{\rm vib}$ and $T^0_{\rm rot}$ we have an R-dependent mass parameter $(1/I+1/\mu R^2)^{-1}$, but the dependence is very weak because $1/\mu R^2$ is about one tenth of $1/I$ in the contact region. We therefore insert the equilibrium distance $R_{\rm e}$ into $(1/I+1/\mu R^2)$. Accordingly the hamiltonian H_0 is separable in three variables;

$$H_0=H_R+H_{\beta_1}+H_{\beta_2}\,, \tag{11}$$

$$H_R=-\frac{\hbar^2}{2\mu}\frac{\partial^2}{\partial R^2}+\frac{k_R}{2}(R-R_{\rm e})^2+E_0 + \frac{\hbar^2}{2}\left[\frac{J(J+1)-\frac{3}{2}K^2-\frac{1}{2}\nu^2}{\mu R_{\rm e}^2} - \frac{1}{2}\left(\frac{1}{I}+\frac{1}{\mu R_{\rm e}^2}\right)\right], \tag{12}$$

$$H_{\beta_i}=\frac{\hbar^2}{2}\left(\frac{1}{I}+\frac{1}{\mu R_{\rm e}^2}\right) \times\left(-\frac{\partial^2}{\partial\beta_i^2}+\frac{(K\pm\nu)^2-1}{4\sin^2\beta_i}\right)+\frac{k_\beta}{2}\beta_i^2\,, \tag{13}$$

where plus and minus signs of $(K\pm\nu)$ are associated with β_1 and β_2, respectively. Eq. (12) represents the simple one-dimensional harmonic oscillator, but eq. (13) does not. By the approximation of $\sin\beta_i\cong\beta_i$, eq. (13) is reduced to a hamiltonian of the kind of a spherical oscillator and we know the eigenfunctions:

$$\varphi_{n\kappa}(\beta)=N_{n\kappa}\beta^{(|\kappa|+1)/2} \times{}_1F_1(-n, \tfrac{1}{2}|\kappa|+1; \sigma\beta^2)\exp(-\tfrac{1}{2}\sigma\beta^2)\,, \tag{14}$$

where $\kappa=K\pm\nu$. In the above approximations the normal modes turn out to be those associated with the variables R, β_1 and β_2, respectively. The latter two variables describe the so-called butterfly and anti-butterfly motions. The energy eigenvalues of the system are given as follows:

$$E^J(n, n_1, n_2, K, \nu)=E_0 + \frac{\hbar^2}{2}\left[\frac{J(J+1)-\frac{3}{2}K^2-\frac{1}{2}\nu^2}{\mu R_{\rm e}^2} - \frac{1}{2}\left(\frac{1}{I}+\frac{1}{\mu R_{\rm e}^2}\right)\right]+(n+\tfrac{1}{2})\hbar\omega + (2n_1+\tfrac{1}{2}|K+\nu|+1)\hbar\omega_{\beta_1} + (2n_2+\tfrac{1}{2}|K-\nu|+1)\hbar\omega_{\beta_2}\,, \tag{15}$$

where $\hbar\omega$, $\hbar\omega_{\beta_1}$ and $\hbar\omega_{\beta_2}$ are vibrational energy quanta for the radial, β_1- and β_2-modes, respectively. E_0 is the minimum value of the potential energy at the equilibrium. Concerning the intrinsic excitations, not only the above vibrational modes but also the precessional rotation modes exist, the quantum states of which are specified by K and ν. K is a projection of the angular momentum onto the z'-axis which specifies coincident rotation of the constituent nuclei. ν specifies the motion which is associated with the variable $\alpha=(\alpha_1-\alpha_2)/2$ and we call this motion the twisting-rotation mode, where the constituent nuclei rotate in an opposite direction with each other.

There is a selection rule, $K\pm\nu=$even, for precessional angular momentum. Because of the parity and boson symmetry, n_1 can be taken to be larger than n_2.

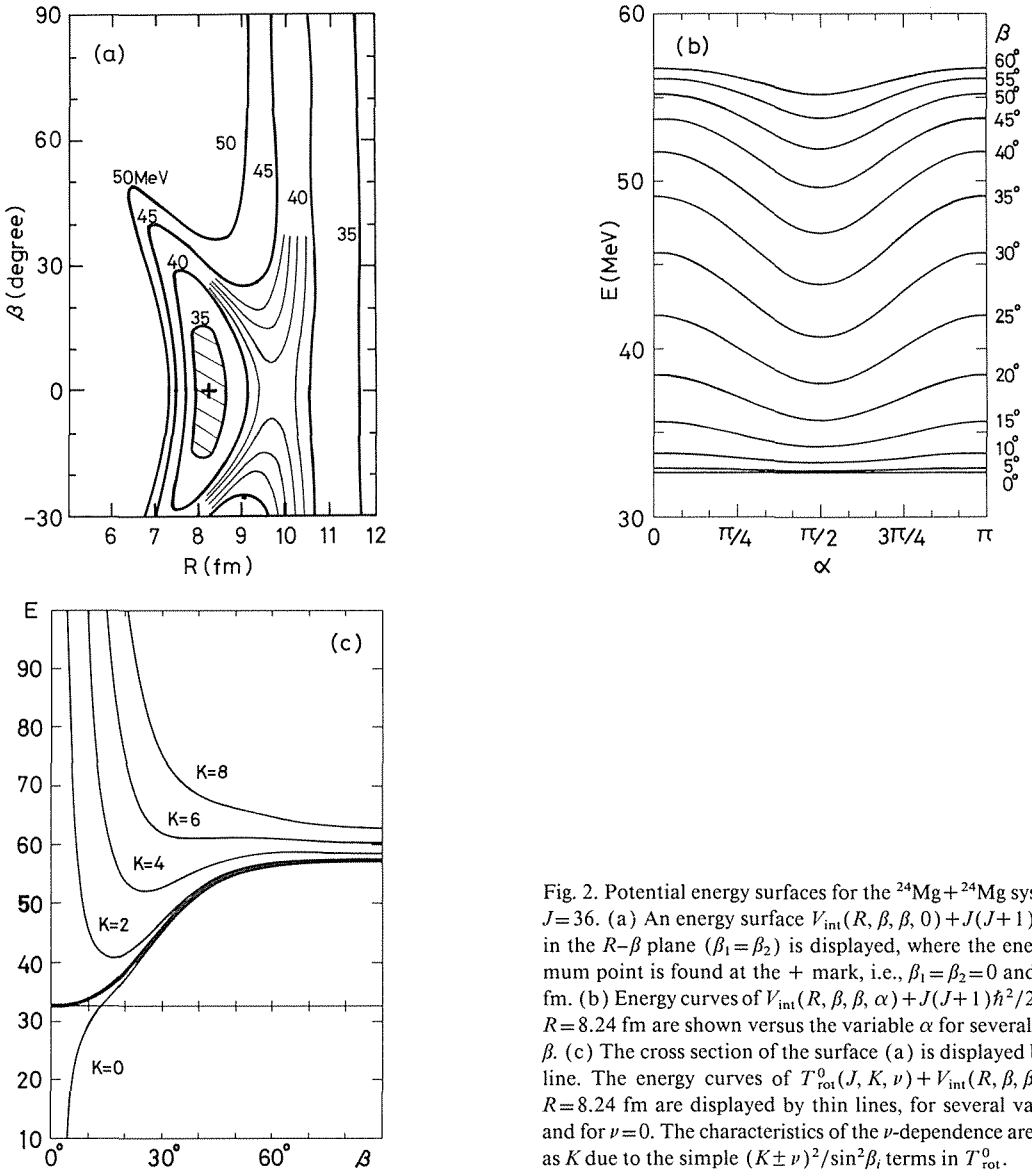

Fig. 2. Potential energy surfaces for the $^{24}Mg+^{24}Mg$ system with $J=36$. (a) An energy surface $V_{\rm int}(R,\beta,\beta,0)+J(J+1)\hbar^2/2\mu R^2$ in the R–β plane ($\beta_1=\beta_2$) is displayed, where the energy minimum point is found at the + mark, i.e., $\beta_1=\beta_2=0$ and $R=8.24$ fm. (b) Energy curves of $V_{\rm int}(R,\beta,\beta,\alpha)+J(J+1)\hbar^2/2\mu R^2$ with $R=8.24$ fm are shown versus the variable α for several values of β. (c) The cross section of the surface (a) is displayed by a thick line. The energy curves of $T^0_{\rm rot}(J,K,\nu)+V_{\rm int}(R,\beta,\beta,0)$ with $R=8.24$ fm are displayed by thin lines, for several values of K and for $\nu=0$. The characteristics of the ν-dependence are the same as K due to the simple $(K\pm\nu)^2/\sin^2\beta_i$ terms in $T^0_{\rm rot}$.

Accordingly, ν is allowed to have negative values if $n_1\neq n_2$ and $K\neq 0$.

We furthermore diagonalize $H_{\rm mode-mode}$ with the basis states obtained above, i.e. with eigenstates of H_0. Since we neglect the Coriolis terms $T'_{\rm C}$, the K-quantum number persists as a good quantum number. The quantum number n of the radial mode also persists by the approximation of $R=R_{\rm e}$ in $T_{\rm mode-mode}$, which is reasonable due to the good localization of the two nuclei around $R_{\rm e}$. The factor $\cot\beta_i$ in $T_{\rm mode-mode}$ is also approximated by β_i^{-1}, consistent with the approximation $\sin\beta_i\cong\beta_i$. With these approximations we can obtain analytic expressions of the matrix elements of the hamiltonian $H_{\rm mode-mode}$. Their magnitudes are found to be very small, namely, to be around 1 MeV. It, however, should be noted

that the energies of several basis states with pure excitations of the normal modes are almost degenerate, if the sum of the quantum numbers $2(n_1+n_2)+\frac{1}{2}|K+\nu|+\frac{1}{2}|\kappa-\nu|$ is equal among them. These almost degenerate levels with the same K-value become separate when $H_{\text{mode-mode}}$ is taken into account. The resulting levels for $J=36$ obtained after diagonalization of $H_{\text{mode-mode}}$ are shown in fig. 3 together with the zeroth-order levels labeled with A–L (hereafter we call them the A-level, etc.). The degenerate C-, D-... levels are seen to be split into separate levels by $H_{\text{mode-mode}}$. The energy gains or losses in them originate mainly from the mutual-orientation-dependent interaction $V_{\text{mode-mode}}$. But the non-degenerate levels such as the A-level are not affected so much, as expected from the weakness of $H_{\text{mode-mode}}$. Intuitive pictures are given for several intrinsic excitation modes at the bottom of the columns. The A-level is the lowest state with $J=36$, without any excitations, i.e., with all the quantum numbers being zero. The B-level is a radially excited state with $n=1$ and the other quantum numbers being zero. The lower level stemming from C-levels corresponds to a butterfly motion and the upper partner to an anti-butterfly motion. D-levels correspond to higher excitations of butterfly and anti-butterfly motions. These excitation modes are expected to reflect their own characteristics in observable quantities such as partial decay widths, not only in their excitation energies (ref. [12], preliminary results are reported in ref. [13]).

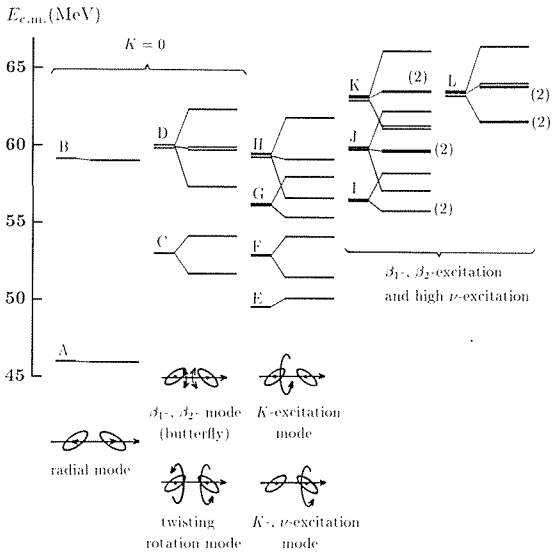

Fig. 3. Molecular normal modes for the $^{24}Mg+^{24}Mg$ system for $J=36$. Degeneracies are given in parentheses on the right-hand side of the levels. Energy levels of normal modes of H_0 are displayed on the left-hand side of the associated levels and are labeled by A–L. The quantum states (n, n_1, n_2, K, ν) of them are as follows; A: (0, 0, 0, 0, 0), B: (1, 0, 0, 0, 0), C: (0, 1, 0, 0, 0) (0, 0, 0, 0, 2), D: (0, 2, 0, 0, 0) (0, 1, 1, 0, 0) (0, 1, 0, 0, 2) (0, 0, 0, 0, 4), E: (0, 0, 0, 1, 1), F: (0, 0, 0, 2, 0), (0, 0, 0, 2, 2), G: (0, 0, 0, 3, 1) (0, 0, 0, 3, 3), H: (0, 0, 0, 4, 0) (0, 0, 0, 4, 2) (0, 0, 0, 4, 4), I: (0, 1, 0, 1, ±1) (0, 0, 0, 1, 3), J: (0, 1, 0, 2, 0) (0, 1, 0, 2, ±2) (0, 0, 0, 2, 4), K: (0, 1, 0, 3, ±1) (0, 1, 0, 3, ±3) (0, 0, 0, 3, 5), L: (0, 2, 0, 1, ±1) (0, 1, 1, 1, 1) (0, 1, 0, 1, ±3) (0, 0, 0, 1, 5).

The molecular states obtained above are intrinsic states of the system, with which rotational bands are associated. Hence many molecular bands are expected to exist. It, however, should be noted that all of them cannot be observed as resonances in $^{24}Mg+^{24}Mg$ collisions, because states with $K\neq 0$ and/or $\nu\neq 0$ carry no partial width of the elastic channel due to the selection rule in angular momentum and therefore they cannot be excited through the $^{24}Mg+^{24}Mg$ entrance channel. Of course some of them may have $K=\nu=0$ components due to the Coriolis coupling T'_C and have some resonance strengths. In fig. 4, the energy spectrum with $K=0$ is shown in the resonance energy region. Experimental resonance levels are shown for comparison at the rightmost column of fig. 4. We have found that the number of the levels is reproduced well in the present calculations.

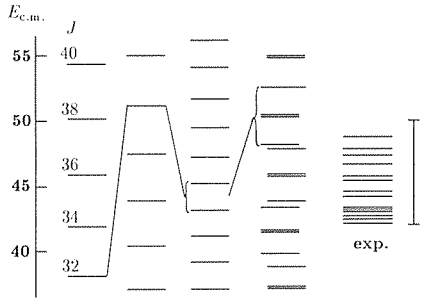

Fig. 4. Di-nucleus spectrum with $K=0$ for the $^{24}Mg+^{24}Mg$ system in the resonance energy region. States with the same total spins are connected by thin lines to guide the eye. We select resonance levels from the elastic and inelastic excitation functions [3], if two prominent peaks are seen to be correlated with each other. They are shown on the right-hand side, with a bar for indicating the available energy region of the data.

4. Concluding remarks

We have proposed a new molecular model, which takes into account many degrees of freedom of the dinuclear system. It describes the dynamics of the interacting system as a sum of normal modes at the equilibrium configuration.

We have applied the model to the $^{24}Mg+^{24}Mg$ system. Various molecular states with excitations of the normal modes have been obtained. Since $H_{\text{mode-mode}}$ depends on the mutual orientation of the ^{24}Mg nuclei, we have diagonalized it and obtained molecular states which embody the intuitive pictures such as butterfly motion, etc. Calculated molecular states with several-J values relevant to the experiment have been found to be in good correspondence with the observed resonances. It is strongly desired that the experiments explore a wider energy region.

The effects of the Coriolis coupling terms T'_C are now under investigation. A closer comparison with experiment is being made with analyses of partial decay widths into various inelastic channels [13]. Experimental data to be compared are strongly desired.

The results obtained above are extremely encouraging for further development of the model. Up to now, a change of deformation of the constituent nuclei themselves is not taken into account, neither is the axial symmetry of ^{24}Mg. The extension of the model which includes such degrees of freedom seems to be very interesting. Applications to systems with oblate deformation, such as $^{28}Si+^{28}Si$, are also interesting.

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