

TOWARD A MICROSCOPIC DESCRIPTION OF THE SUBCOULOMB $^{12}\text{C} + ^{12}\text{C}$ RESONANCES

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A properly antisymmetrized, SU_3 -coupled $^{12}\text{C} + ^{12}\text{C}$ cluster model basis is used with simplified $^{12}\text{C} + ^{12}\text{C}$ interactions of $Q \cdot Q$ type to investigate the question: Is a simple quasimolecular basis sufficient to account for the observed density of low- J carbon resonances?

The phenomenon of the so-called quasimolecular resonances constitutes one of the oldest problems [1] in the interaction of heavy ions. Although much theoretical progress has been made and a vast amount of experimental data has been accumulated, we are far from a complete understanding of the resonances in the excitation functions of the $^{12}\text{C} + ^{12}\text{C}$ and similar heavy ion systems. A large number of theoretical studies of such resonances seem to confirm the usefulness of the quasimolecular picture. Yet, an interpretation of the gross structure of the $^{12}\text{C} + ^{12}\text{C}$ resonances seems to be possible without invoking the existence of molecular pockets in the potentials for the relative motion of the two ions [2]. Hard evidence for the quasimolecular picture must thus be sought in the fine structure of these resonances. The large number of closely spaced, narrow fine structure components are the distinctive feature of the $^{12}\text{C} + ^{12}\text{C}$ resonances, compared with the simpler fine structure of the $^{16}\text{O} + ^{16}\text{O}$ resonances, for example. This may be related to the nonspherical shape of the two ^{12}C nuclei, or in the language of SU_3 symmetry, the much richer SU_3 structure of the $^{12}\text{C} + ^{12}\text{C}$ cluster model wave functions. The $^{12}\text{C} + ^{12}\text{C}$ system may thus furnish the strictest test of the molecular interpretation. In addition, the subcoulomb and Coulomb barrier region of the $^{12}\text{C} + ^{12}\text{C}$ system is one of the most thoroughly studied [3]. The present note focuses on this region of the $^{12}\text{C} + ^{12}\text{C}$ system, since an interpretation in terms of quasibound states of molecular character should be most relevant for this energy range. A recent analysis of the subcoulomb $^{12}\text{C} + ^{12}\text{C}$ resonances by Erb and Bromley [4] seems to indicate that the observed fine structure in this region can be systematized remarkably well by the energy pattern,

$$E = -D + a(v + \frac{1}{2}) - b(v + \frac{1}{2})^2 + cJ(J + 1),$$

characteristic of diatomic molecular phenomena, although no simple molecular vibration-rotation interpretation was necessarily implied.

Most recent theoretical interpretations of the $^{12}\text{C} + ^{12}\text{C}$ resonances make use of some variant of the Nogami-Imanishi model [5], based on an internal excitation of the ^{12}C nuclei, or its generalization by Greiner et al. [6].

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The band crossing model of Kondō et al. [7], and the recent calculations of Koennecke et al. [8], in particular, reproduce the degree of complexity of the fine structure observed in particular exit channels. However, none of these models are sufficiently microscopic to be capable of examining in detail the large amount of recent experimental information involving specific states in $\alpha + {}^{20}\text{Ne}$, $p + {}^{23}\text{Na}$, $n + {}^{23}\text{Mg}$, ${}^{16}\text{O} + {}^8\text{Be}$, as well as ${}^{12}\text{C} + {}^{12}\text{C}$ exit channels. A microscopic treatment has recently been undertaken by Tohsaki–Suzuki [9] who couples ${}^{12}\text{C} + {}^{12}\text{C}$, $\alpha + {}^{20}\text{Ne}$, and ${}^{16}\text{O} + {}^8\text{Be}$ channels, however, in terms of intrinsic internal wave functions where the details of the angular momentum structure of the various channels is suppressed. This treatment predicts a subcoulomb band and three bands in the Coulomb barrier region of the ${}^{12}\text{C} + {}^{12}\text{C}$ system but is as yet unable to examine the detailed rotational fine structure of such bands.

Although a fully microscopic multichannel resonating group or generator coordinate treatment of this challenging problem is perhaps not yet within reach, recent detailed calculations of norm and overlap matrix elements for the ${}^{12}\text{C} + {}^{12}\text{C}$ system [10,11], may perhaps be combined with simplified ${}^{12}\text{C} + {}^{12}\text{C}$ interactions of the $Q \cdot Q$ type to examine the question: Is a simple ${}^{12}\text{C} + {}^{12}\text{C}$ molecular picture sufficient to account for the observed density of low- J resonant states in the subcoulomb region, and does such a picture lead to strength functions of the type observed to the various $\alpha + {}^{20}\text{Ne}$, $p + {}^{23}\text{Na}$, ... exit channels?

For the closely spaced, narrow, isolated resonances observed in the excitation functions of the ${}^{12}\text{C} + {}^{12}\text{C}$ reaction the T matrix may be approximated by a sum over resonant states ν

$$T_{cc'} = i \exp(i\delta_c + i\delta_{c'}) \sum_{\nu} \frac{\Gamma_{\nu c}^{1/2} \Gamma_{\nu c'}^{1/2}}{E - E_{\nu} - i\Gamma_{\nu}/2}, \quad (1)$$

where the entrance channel c consists of ${}^{12}\text{C}(0^+) + {}^{12}\text{C}(0^+)$. The important exit channels include $\alpha + {}^{20}\text{Ne}$, $p + {}^{23}\text{Na}$, $n + {}^{23}\text{Mg}$, ${}^8\text{Be} + {}^{16}\text{O}$, with contributions from many low-lying excited states of ${}^{20}\text{Ne}$ and ${}^{23}\text{Na}$, in particular. It will be assumed that the resonant states $|\nu\rangle$ can be built from a basis of quasibound states $|i\rangle$, $|\nu\rangle = \sum |i\rangle \times \langle i|\nu\rangle$, which are then to be coupled to continuum states of the various exit channels. No attempt is made to calculate widths, Γ_{ν} , or the partial width amplitudes $\Gamma_{\nu c}^{1/2}$. We concentrate instead on the question: is a basis of molecular quasibound states $|i\rangle$ sufficient to account for the observed density of states $|\nu\rangle$.

The states $|i\rangle$ are to be built from a “molecular” ${}^{12}\text{C} + {}^{12}\text{C}$ cluster function basis

$$|\Psi_{\text{cluster}}\rangle = \mathcal{A} | \{ [\phi_{I_1}({}^{12}\text{C})^{(04)} \times \phi_{I_2}({}^{12}\text{C})^{(04)}]_{I_c} \times \Phi(\mathbf{R}_{\text{relative}})_L \} JM_J \rangle, \quad (2)$$

where \mathcal{A} is the antisymmetrization operator. The internal wave function ϕ of the two ${}^{12}\text{C}$ nuclei is assumed to have a good Wigner supermultiplet and hence SU_3 symmetry, given by the Elliott quantum $(\lambda\mu) = (04)$, with $I_1 (I_2) = 0, 2, 4$. That is, in our molecular wave function each ${}^{12}\text{C}$ nucleus can change its state of orientation, but no other intrinsic excitations of the ${}^{12}\text{C}$ nuclei are included. The rotation of ${}^{12}\text{C}$ about ${}^{12}\text{C}$ contained in the relative motion function Φ can take place with many states of relative orientation of the two nonspherical ${}^{12}\text{C}$ nuclei, giving rise to a large number of states for each J . The ${}^{12}\text{C}-{}^{12}\text{C}$ relative motion function should couple properly to the continuum functions of entrance and exit channels. An expansion in oscillator functions $\chi^{(Q0)}(\mathbf{R})$ should thus contain a large number of terms, including oscillator excitations of very high Q , $Q = 2N + L$. The basis of quasibound states $|i\rangle$, however, can be expected to have a large overlap with shell model functions of $2\hbar\omega$ oscillator excitation. Quasimolecular basis states of this type correspond to ${}^{12}\text{C} + {}^{12}\text{C}$ cluster functions in which the relative motion function is restricted to a single oscillator function with $Q = 14 [=Q(\text{minimum Pauli allowed}) + 2]$, that is, a projection of the ${}^{12}\text{C} + {}^{12}\text{C}$ cluster function onto the shell model space of $2\hbar\omega$ excitation

$$|\Psi\rangle = \mathcal{A} | \{ [\phi({}^{12}\text{C})^{(04)} \times \phi({}^{12}\text{C})^{(04)}]^{(\lambda_c\mu_c)} \times \chi(\mathbf{R})^{(Q0)=(14,0)} \}^{(\lambda\mu)_K} JM_J \rangle. \quad (3)$$

This cluster function is here given in SU_3 coupled form. The ^{12}C internal functions of SU_3 symmetry (04) are first coupled to SU_3 symmetry $(\lambda_c \mu_c)$. For this system of two identical nuclei only the symmetrically coupled internal functions with $(\lambda_c \mu_c) = (08), (24), (40)$ can couple with relative motion functions of even Q , hence even L and positive parity. The functions, (3), with all possible $(\lambda_c \mu_c) (\lambda \mu) \kappa JM_J$ are to constitute our molecular basis. If the antisymmetrization operator were to be ignored, and only the identity of the two ^{12}C nuclei were to be taken into account, this basis would contain 14 0^+ states, 40 2^+ , and 57 4^+ states, for example. The Pauli-allowed states, however, are restricted to 7 0^+ , 20 2^+ , and 26 4^+ states. Pauli-allowed states are the eigenvectors of \mathcal{A} in the above basis, with eigenvalues Λ_i different from zero

$$\mathcal{A} |(\lambda \mu)_i, \kappa JM_J\rangle = \Lambda_{(\lambda \mu)_i} |(\lambda \mu)_i, \kappa JM_J\rangle, \tag{4a}$$

$$|(\lambda \mu)_i, \kappa JM_J\rangle = \sum_{(\lambda_c \mu_c)} [(\lambda_c \mu_c) \times (14,0)] (\lambda \mu) \kappa JM_J \langle [(\lambda_c \mu_c) \times (14,0)] (\lambda \mu) | (\lambda \mu)_i \rangle. \tag{4b}$$

The fully antisymmetrized, normalized state vectors $|i\rangle$

$$|i\rangle \equiv \Lambda_{(\lambda \mu)_i}^{-1/2} \mathcal{A} |(\lambda \mu)_i, \kappa JM_J\rangle, \tag{4c}$$

constitute our $^{12}C + ^{12}C$ quasimolecular basis from which the quasibound states $|\nu\rangle, |\nu\rangle = \sum |i\rangle \langle i|\nu\rangle$, are to be built. The eigenvectors, (4), have been calculated by Bargmann-integral-transform, SU_3 -recoupling techniques and have been tabulated in ref. [10], (see table 6 of ref. [10] for the transformation coefficients $\langle [(\lambda_c \mu_c) \times (14,0)] \times (\lambda \mu) | (\lambda \mu)_i \rangle$). Overlap matrix elements with different cluster decompositions, (see table 8 of ref. [10]) have also been calculated, making it possible to calculate spectroscopic amplitudes to most of the important states in $\alpha + ^{20}Ne$, $p + ^{23}Na$, $n + ^{23}Mg$, and $^{16}O + ^8Be$ exit channels. Although no attempt has been made to calculate partial width amplitudes $\Gamma_{\nu c}^{1/2}$, it is expected that the spectroscopic amplitudes $A_{\nu c}$ would give some measure of the relative importance of the states $|\nu\rangle$ in the observed excitation functions of specific $^{12}C + ^{12}C \rightarrow A + B$ reactions. The observed resonances of the same J^π fall within an energy interval of ~ 3 MeV, and in our model individual resonant states $|\nu\rangle$ are made up of quasibound states of very similar character. The penetrability factors associated with a product of partial width amplitudes $(\Gamma_{\nu c}^{1/2} \cdot \Gamma_{\nu c'}^{1/2})$ can thus be expected to be approximately independent of ν , and the relative magnitudes of the corresponding product of spectroscopic amplitudes $(A_{\nu c} A_{\nu c'})$ can be expected to give some measure of the relative magnitudes of the excitation functions for specific channels c' .

Our main aim, however, is to determine whether the quasimolecular basis, (4), can give rise to the rich fine structure of resonant states ν observed in the subcoulomb region. The hamiltonian is decomposed into

$$H = H_{i_1} + H_{i_2} + H_{relative} + H_{interaction}, \tag{5}$$

where the excitations of the ^{12}C internal degrees of freedom, specified by H_{i_1}, H_{i_2} , are restricted to $0^+, 2^+, 4^+$ internal ^{12}C energies of 0 MeV, 4.44 MeV, 14.08 MeV (taken from experiment). The relative motion hamiltonian will have a centrifugal term, leading to an $L(L + 1)$ -dependence.

The splitting of the quasibound states into different fine structure components must com from $H_{interaction}$. A simple $Q \cdot Q$ interaction between the ^{12}C ions can be expected to play a dominant role in this interaction. A transformation from single particle to internal and relative motion coordinates of the two ^{12}C gives the quadrupole operator in the form $Q = Q_{i_1} + Q_{i_2} + Q_{relative}$ and leads to interaction terms of the type $(Q_{i_1} \cdot Q_{i_2}), (Q_{relative} \cdot (Q_{i_1} + Q_{i_2}))$. In the SU_3 -coupled basis of eq. (3) these interactions can be approximated by Elliott type $Q \cdot Q$ interactions, with eigenvalues

$$(Q_{i_1} \cdot Q_{i_2}) = \frac{1}{2} \{ C_{(\lambda_c \mu_c)} - 2C_{(04)} - \frac{3}{4} [I_c(I_c + 1) - I_1(I_1 + 1) - I_2(I_2 + 1)] \}, \tag{6}$$

$$(Q_{rel} \cdot (Q_{i_1} + Q_{i_2})) = \frac{1}{2} \{ C_{(\lambda \mu)} - C_{(\lambda_c \mu_c)} - C_{(Q0)} - \frac{3}{4} [J(J + 1) - L(L + 1) - I_c(I_c + 1)] \}, \tag{7}$$

where the SU_3 Casimir invariants are given by $C_{(\lambda \mu)} = \lambda^2 + \lambda \mu + \mu^2 + 3\lambda + 3\mu$. The relative motion hamiltonian can be expected to contribute an additional $L(L + 1)$ dependence. As a slight generalization of the above, we

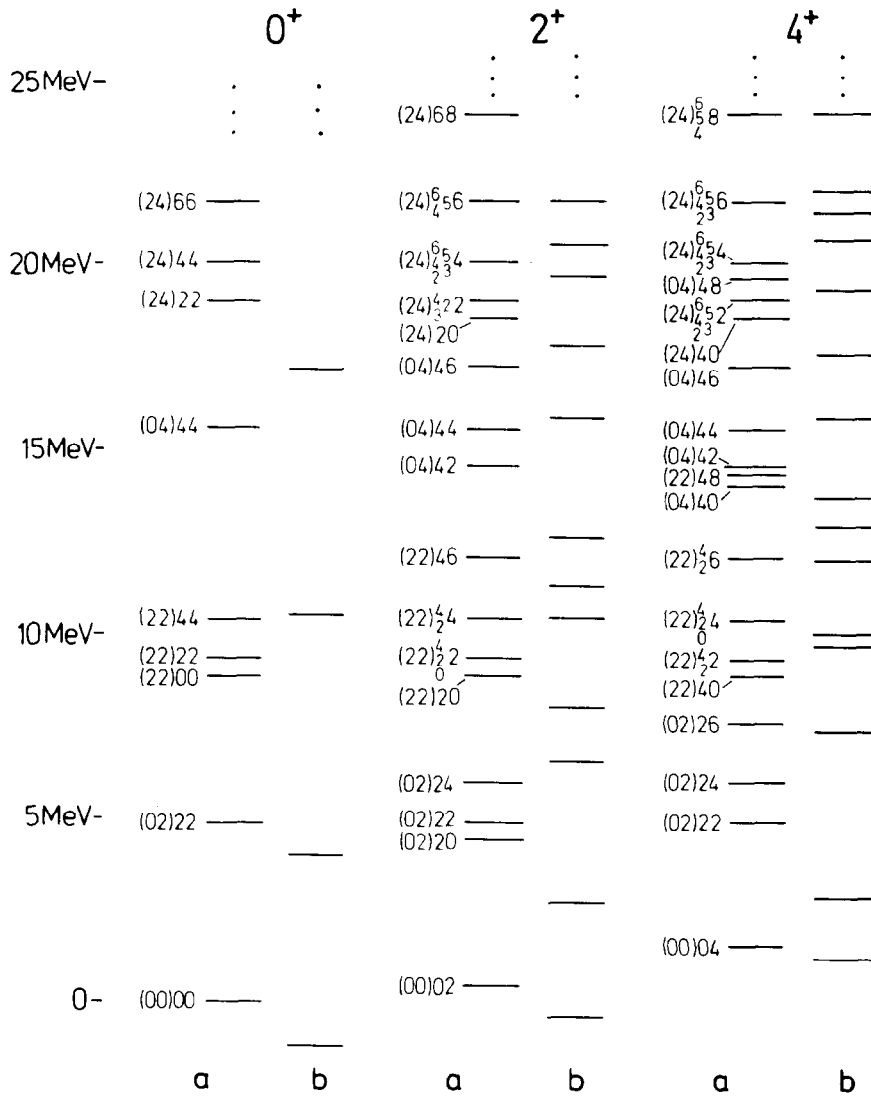


Fig. 1. The effect of antisymmetrization. Energies are plotted for $H = E_{\text{internal}} + \delta L(L + 1)$, with $E_{\text{internal}} = E_{i_1} + E_{i_2}$, (with $E_0^+ = 0$ MeV, $E_2^+ = 4.44$ MeV, $E_4^+ = 14.08$ MeV), and with $H_{\text{interaction}}$ "turned off". The rotational constant $\delta = 76$ keV is taken from ref. [4]. Columns (a) show the 0^+ , 2^+ , 4^+ energies with antisymmetrization effects neglected completely. The angular momenta $(I_1 I_2) I_c L$ are indicated on each level. Columns (b) show the energies for the properly antisymmetrized basis, see eqs. (4).

choose the schematic hamiltonian

$$H = E_{i_1} + E_{i_2} + \alpha [C_{(\lambda\mu)} - \frac{3}{4} J(J + 1)] + \frac{3}{4} \beta [I_1(I_1 + 1) + I_2(I_2 + 1)] + \gamma I_c(I_c + 1) + \delta L(L + 1) + \nu C_{(\lambda_c \mu_c)}. \quad (8)$$

The diagonalization of this hamiltonian in the antisymmetrized basis, (4), constitutes our model for the resonant states $|\nu\rangle$. Matrix elements of this model H are given by

$$\begin{aligned}
\langle (\lambda' \mu') i', \kappa' J M_J | H | (\lambda \mu) i, \kappa J M_J \rangle = & \alpha [C_{(\lambda \mu)} - \frac{3}{4} J(J+1)] \delta_{(\lambda \mu)(\lambda' \mu')} \delta_{i i'} \delta_{\kappa \kappa'} \\
& + \sum_{(\lambda_c \mu_c)} \sum_{(\lambda'_c \mu'_c)} \langle [(\lambda_c \mu_c) \times (14, 0)] (\lambda \mu) | (\lambda \mu) i \rangle \langle [(\lambda'_c \mu'_c) \times (14, 0)] (\lambda' \mu') | (\lambda' \mu') i' \rangle \\
& \times \sum_{\kappa_c \kappa'_c} \sum_{I_c L} \langle (\lambda_c \mu_c) \kappa_c I_c; (14, 0) L | | (\lambda \mu) \kappa J \rangle \langle (\lambda'_c \mu'_c) \kappa'_c I_c; (14, 0) L | | (\lambda' \mu') \kappa' J \rangle \\
& \times \left[\gamma I_c (I_c + 1) + \delta L (L + 1) + \nu C_{(\lambda_c \mu_c)} \right] \delta_{(\lambda_c \mu_c)(\lambda'_c \mu'_c)} \delta_{\kappa_c \kappa'_c} \\
& + \sum_{I_1 I_2} \langle (04) I_1; (04) I_2 | | (\lambda_c \mu_c) \kappa_c I_c \rangle \langle (04) I_1; (04) I_2 | | (\lambda'_c \mu'_c) \kappa'_c I_c \rangle \\
& \times \left(\frac{3}{4} \beta [I_1 (I_1 + 1) + I_2 (I_2 + 1)] + \sum_{j=1,2} (0\delta_{I_j 0} + 4.44\delta_{I_j 2} + 14.08\delta_{I_j 4}) \right) \Big], \tag{9}
\end{aligned}$$

Where the double-barred coefficients are $SU_3 \supset R_3$ reduced Wigner coefficients. Since Pauli-forbidden components are excluded from the basis $|i\rangle$, the \mathcal{A} operator plays an important role in determining the positioning and spacing of fine structure components, $|\nu\rangle$. This is illustrated in fig. 1. Columns marked (a) show $J^\pi = 0^+, 2^+, 4^+$ states with both $H_{\text{interaction}}$ and the Pauli principle “turned off”, (that is, with antisymmetrization effects neglected completely). These are shown here because such $(I_1 I_2)L$ -dependent $^{12}C + ^{12}C$ energies have been taken as the starting point in many earlier treatments of the carbon resonances. Columns (b) show the $0^+, 2^+, 4^+$ states in the same basis, with $H_{\text{interaction}}$ still “turned off”, but with the antisymmetrizer “turned on”; that is, in the properly antisymmetrized basis spanning the same $(I_1 I_2)LJ$ values for $Q = 14$. The Pauli principle not only reduces the density of levels but is responsible for large effective energy shifts. The exclusion of the Pauli-forbidden states is thus of prime importance in the treatment of the low- J carbon resonances. It should, however, be emphasized that this remark applies to the quasibound states in the subcoulomb region. At much higher excitation energies the $^{12}C + ^{12}C$ relative motion functions will be dominated by oscillator functions with $Q > 14$, and antisymmetrization effects will be less important. In the high-energy regime a weak-coupled cluster function of the form of eq. (2) should give the best approximation, and in this case I_1 and I_2 are good quantum numbers, even in the presence of \mathcal{A} . In the limit in which the quasibound states are dominated by cluster functions of a single Q , the SU_3 strong-coupled functions of eqs. (3) and (4), with definite $(\lambda \mu)$, become eigenfunctions of \mathcal{A} and hence the most convenient basis for the diagonalization of the hamiltonian, (8). The observed rotational spacing of the carbon resonances sets fairly stringent limits on the parameters of the schematic hamiltonian. Many parameter sets have been tried, often with qualitatively similar results. A characteristic example is shown in fig. 2 which gives the positions of the eigenvalues, E_ν , of the hamiltonian (9) for the parameter set

$$\alpha = \beta = -0.04 \text{ MeV}, \quad \gamma = 0 \text{ MeV}, \quad \delta = 0.15 \text{ MeV}, \quad \nu = -0.10 \text{ MeV}.$$

To give some measure of the relative importance of each quasibound state for the reaction cross section, we plot a strength factor, \bar{S} , defined in terms of spectroscopic amplitudes

$$\bar{S} = A_{\nu c} \left(\sum_{c'} A_{\nu c'}^2 \right)^{1/2}, \tag{10}$$

where $A_{\nu c}$ is the spectroscopic amplitude for the entrance channel. The sum over spectroscopic amplitudes $A_{\nu c'}$ is taken over those exit channels, c' , included in the experimental excitation function. (Spectroscopic amplitudes for most of the open exit channels in $\alpha + ^{20}\text{Ne}$, $p + ^{23}\text{Na}$, $n + ^{23}\text{Mg}$, and $^8\text{Be} + ^{16}\text{O}$ have been calculated, in SU_3 approximation, in ref. [10].) Our model makes no attempt to calculate the absolute energy positions of the resonances. However, the lower states can be expected to fall in the subthreshold region or be effectively eliminated by very small Coulomb penetrability factors. States of very high energy (not shown) will be eliminated by the reaction mechanism. The states of fig. 2 encircled by the square box are the expected candidates for the subcoulomb reso-

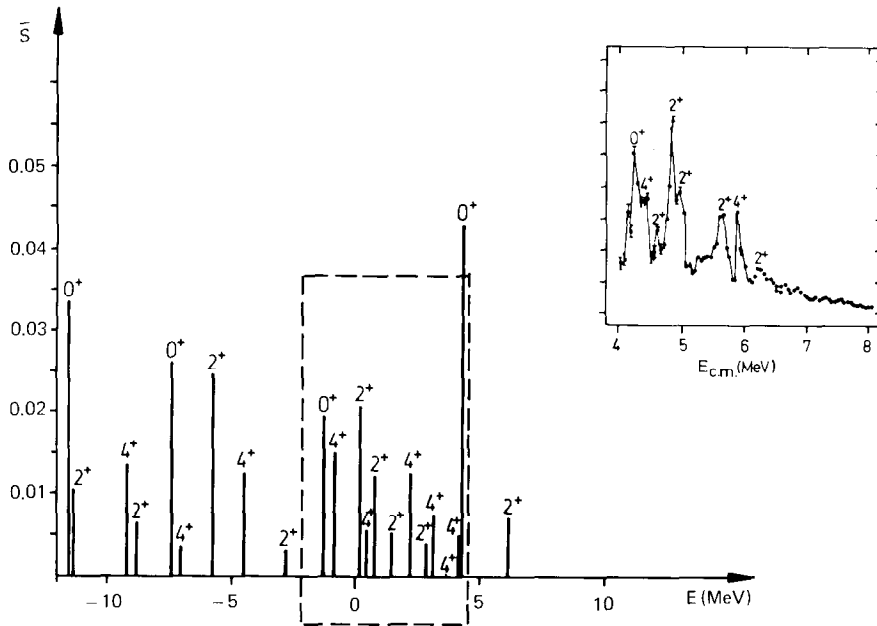


Fig. 2. Energies of the hamiltonian (8) with $\alpha = \beta = -0.04$ MeV, $\gamma = 0$ MeV, $\delta = +0.15$ MeV, $\nu = -0.1$ MeV. No attempt is made to calculate the absolute position of the "0" of the energy scale. The states encircled by the square box are to be compared with the experimental excitation function of the insert, taken from ref. [12] for the sum of the transition 0-5 of the reaction $^{12}\text{C}(^{12}\text{C}, \alpha)^{20}\text{Ne}$. The vertical lines give the \bar{S} factors defined by eq. (10), where the sum includes the 0^+ , 2^+ , 4^+ , 2^- , and 3^- states in ^{20}Ne .

nances. The insert shows an experimental excitation function for the subcoulomb region. The results of the experimental excitation functions in the subcoulomb region are usually displayed in terms of nuclear structure or astrophysical S -factors. The insert of fig. 2 is taken from ref. [12] and shows the experimentally measured S -factor for the summed transitions 0-5 in the $^{12}\text{C}(^{12}\text{C}, \alpha)^{20}\text{Ne}$ reaction. The sum over c' for the \bar{S} -factor for this figure is thus taken to include the 0^+ , 2^+ , 4^+ members of the ground-state band of ^{20}Ne , approximated by good SU_3 symmetry $(\lambda\mu) = (80)$ and the 2^- , 3^- members of the $\kappa = 2$ band at 4.97 MeV in the SU_3 approximation, $(\lambda\mu) = (82)$. Except for the upper 0^+ state the predicted model quasibound states and \bar{S} factors have qualitatively many of the features of the experimentally observed resonances and nuclear structure factors. (It is not clear whether the high 0^+ state is out of place or whether its large amplitude leads to a broader resonance which may be partially lost in the background.)

Our main purpose, however, is to answer the question; does our simplified quasimolecular basis lead to the observed density of resonances? The observed excitation function contains at least six isolated, narrow 2^+ resonances in a 3 MeV interval of the subcoulomb region. Our model predicts at most four 2^+ quasibound states in a comparable interval. Although a more detailed treatment may lead to a compression of the predicted density of resonances, it seems more likely that coupling to other, essentially nonmolecular degrees of freedom must be included to gain a detailed description of the subcoulomb resonances in the $^{12}\text{C} + ^{12}\text{C}$ system. These degrees of freedom may of course involve more complicated intrinsic excitations in ^{12}C itself, such as the 0^+ state at 7.66 MeV.

We conclude that the properly antisymmetrized $^{12}\text{C} + ^{12}\text{C}$ molecular basis, involving many states of different relative orientation of two rotating nonspherical ^{12}C nuclei, as described by the SU_3 coupled $^{12}\text{C} + ^{12}\text{C}$ cluster function, can account for much of the observed fine structure in the subcoulomb resonance region, but that coupling to nonmolecular degrees of freedom or more complicated excitations in the $A = 24$ system are likely to play a role in accounting for all the details of the observed rich fine structure.

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