Dielectric anomaly in coupled rotor systems

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The correlated dynamics of coupled quantum rotors carrying electric dipole moment is theoretically investigated. The energy spectra of coupled rotors as a function of dipolar interaction energy are analytically solved. The calculated dielectric susceptibilities of the system show a peculiar temperature dependence different from that of isolated rotors.

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I. INTRODUCTION

With the advent of nanotechnologies, quantum rotors have attracted much attention in relevancy to a fundamental element of molecular scale machinery.1–3 Arrays of surface mounted quantum rotors with electric dipole moments are of particular interest because dipole-dipole interactions can be controlled and even designed to yield specific behavior, such as ferroelectricity. Ordered two-dimensional arrays of dipole rotors yield either ferroelectric or antiferroelectric ground states, depending on the lattice type, while disordered arrays are predicted to form a glass phase.4,5

Besides technological problems, the microscopic dynamics of quantum rotors have been extensively studied from the point of physical and chemical interest. The idea of quantum rotors is applicable to interstitial oxygen impurities in crystalline germanium, where oxygen atoms are quantum-mechanically delocalized around the bond center position.6 The rotation of oxygen impurities around the Ge-Ge axis has been experimentally observed by phonon spectroscopy.7 While the rotation of oxygen impurities in Ge is weakly hindered by an azimuthal potential caused by the host lattice, several materials are known to show a free rotation of molecules. An example is ammonia groups in certain Hofmann clathrates M(NH₃)₂M’(CN)₄–G,8–10 usually abbreviated as M·M’-G, where M and M’ are divalent metal ions and G is a guest molecule. Nearly free uniaxial quantum rotation of NH₃ has been observed for the first time in Ni-Ni-(C₆D₆)₂ by inelastic neutron scattering.8 Recently, a surprising variation of the linewidth has been observed for Ni-Ni-(C₁₂H₁₀)₂,11 which has been interpreted by a novel line broadening mechanism based on rotor-rotor coupling.12 It is also known that the β phase of solid methane13 as well as methane hydrate14 shows almost free rotation of CH₄ molecule. The linewidths of methane in clathrates show inhomogeneous broadening owing to the dipolar coupling with water molecules.15 It is therefore expected that new interesting phenomena will be found by investigating the influence of dipolar interaction between quantum rotors.

In the present paper, we study the correlated dynamics of coupled quantum rotors carrying electric dipole moments. We give the exact solution of eigenvalue problem of interacting rotors with arbitrary configurations. It is revealed that coupled rotors show a peculiar dielectric response at low temperatures, which can be interpreted by taking account of the selection rule of dipolar transition for coupled rotors.

II. HAMILTONIAN

Suppose two dipole rotors qᵣ₁ and qᵣ₂ are separated by the vector R. The Hamiltonian for the system is given by

\[ H = H_K + W_D, \]

where the kinetic term is

\[ H_K = -\frac{\hbar^2}{2I} \left( \frac{\partial^2}{\partial \Theta_1^2} + \frac{\partial^2}{\partial \Theta_2^2} \right) \]

and the interaction term becomes

\[ W_D = \frac{q^2}{4\pi \varepsilon} \left( \frac{1}{|R|} + \frac{1}{|R+r_1-r_2|} - \frac{1}{|R+r_1|} - \frac{1}{|R-r_2|} \right). \]

Here I is the moment of inertia for dipole rotors and ε is the dielectric constant. Figure 1 shows a configuration of two dipoles rotors under consideration. We assume that rotors do not feel any potential variation along a ring of radius r. In the Jacobi coordinate, the vectors r₁, r₂, and R are given by

\[ r₁ = (\cos \Theta_1, \sin \Theta_1 \cos \alpha_1, \sin \Theta_1 \sin \alpha_1), \]

\[ r₂ = (\cos \Theta_2 \cos \beta - \sin \Theta_2 \cos \alpha_2 \sin \beta), \]

FIG. 1. Schematic configuration of coupled rotors. Each rotor represented by r₁ and r₂ rotates along a ring of a radius r, and carries dipole moment qᵣ₁(2) .
The higher-order term $W_D^{(4)}=W_D-W_D^{(3)}$ is $O(r^3/R^4)$, which can be negligible for the case $R \gg r$. Actually we have confirmed that the calculated results presented in this paper change very little by taking into account the term $W_D^{(4)}$.

### III. EIGENVALUES AND EIGENFUNCTIONS

The Schrödinger equation for the Hamiltonian $H_0=H_K+W_D^{(3)}$ has analytic solutions as shown below. Transforming variables to $\theta_1=(\Theta_1+\Theta_2)/2$ and $\theta_2=(\Theta_1-\Theta_2)/2$, Eqs. (1) and (4) yield

$$H_K=-\frac{\hbar^2}{4I} \left( \frac{\partial^2}{\partial \theta_1^2} + \frac{\partial^2}{\partial \theta_2^2} \right).$$

The parameters $c_i$ and $\gamma_i(i=1,2)$ are functions of angles $\alpha_1, \alpha_2$, and $\beta$ defined in Fig. 1, whose explicit forms are given by

$$c_i=\frac{1}{2}\sqrt{x^2+y^2}, \quad \gamma_i=\frac{1}{2}\tan^{-1}\left(\frac{-x_i}{y_i}\right),$$

with the definitions

$$x_1=\sin \beta(\cos \alpha_1-\cos \alpha_2),$$

$$x_2=\sin \beta(\cos \alpha_1+\cos \alpha_2),$$

$$y_1=\cos \beta(1-\cos \alpha_1\cos \alpha_2)+2 \sin \alpha_1 \sin \alpha_2,$$

$$y_2=\cos \beta(1+\cos \alpha_1\cos \alpha_2)-2 \sin \alpha_1 \sin \alpha_2.$$
Even if the higher-order term and these two localized eigenstates are nearly degenerate, the eigenfunctions are strongly localized around the two minima, the unchanged values of Fig. 3 slightly shift upwards and/or downwards except for the unchanged values of $E$ at $E_D = 0$.

IV. DIELECTRIC SUSCEPTIBILITIES

Let us consider the dielectric response of dipole rotors coupled via dipolar interaction. The real part of the frequency-dependent dielectric susceptibility is expressed as

$$
\chi_{\mu\nu}(\omega, T) = -\frac{2}{\varepsilon Z} \sum_{j, \sigma_j} |\langle E_j| p_\mu |E_i\rangle|^2 \times \frac{E_j - E_i}{(E_j - E_i)^2 - (\hbar \omega)^2} \exp \left( -\frac{E_j}{k_B T} \right),
$$

where $Z = \sum_j \exp(-E_j/k_B T)$ is the partition function and $|E_j\rangle$ is the eigenvector belonging to the eigenvalue $E_j$. The quantity $p_\mu$ is the $\mu$ component of the total dipole moment $p = q(r_1 + r_2)$, which depends on the relative orientation with respect to the external field. We should note that the susceptibility depends on the selection rules for dipole transitions between different eigenstates. In Fig. 3, allowed dipole transitions for $p_\mu$ are indicated in parts by solid arrows. Note that only a part of allowed transitions are shown in the figure, which are dominant for the dielectric susceptibility $\chi(\omega, T)$ at temperatures $T = E_K / k_B$. The rest of the allowed dipolar transitions do not contribute to the susceptibility given by Eq. (11), because the energy difference $|E_j - E_i|$ is so large and/or the Boltzmann factor $\exp(-E_j/k_B T)$ becomes much smaller than unity. The interpretation of three labels (a)–(c) shown in Fig. 3 will be given later.

We have calculated the temperature dependence of the dielectric susceptibility $\chi_{\mu\mu}(\omega, T)$ for various $E_D$. Figure 4 shows the calculated results of dc susceptibility $\chi_{\mu\mu}(T)$ normalized by a factor $q^2 r^2/(\varepsilon E_K)$. We have taken four values of $E_D/E_K$: the solid line ($E_D = 0$), the dashed one ($E_D/E_K = 0.01$), the dash-dotted one ($E_D/E_K = 0.1$), and the dotted one ($E_D/E_K = 1.0$). The angles are set to be $(\alpha_1, \alpha_2, \beta) = (\pi/4, 0, 0)$ for all $E_D$. For the case of $E_D = 0$, the susceptibility monotonically increases with decreasing temperature and becomes constant at lower temperatures. The crossover temperature between the steady increase and the almost constant value in Fig. 4 is determined by the minimum-energy difference of eigenstates at $E_D = 0$ that are allowed for dipole transition, namely, indicated as (a) in Fig. 3. For the case of $E_D/E_K > 1$, the strong Coulomb interaction prevents from the relative motion of rotors so that the magnitude of the susceptibility $\chi(T)$ decreases with increasing $E_D$.

It is noteworthy that, for relatively weak interaction $E_D/E_K < 0.1$, a bump appears in the susceptibility at about $E_K / (k_B T) = 5.0$. The kinetic energy $E_K = \hbar^2 / (2l)$ for actual rotating molecules is of the order of 1 meV, indicating that the characteristic temperature $T^* = E_K / k_B \times 0.2$ corresponding to the bump is estimated as about 1 K. We made sure that
the bump can be observed for any angles \((\alpha_1, \alpha_2, \beta)\) when \(E_D/E_K\) is less than 0.1. This anomaly stems from the correlated rotation of paired rotors via the dipolar interaction, and can be interpreted by the argument on the selection rule for dipolar transition.

To understand the origin of the bump, we decompose the total susceptibility \(\chi_{xx}(T)\) given in Eq. (11) as

\[
\chi_{xx}(T) = \sum_{(i,j+l)} \eta_{i,j}(T),
\]

where \(\eta_{i,j}(T)\) is the summation over all possible combinations of \((i,j)\) under the condition \(l \neq j\). Note the fact that only three components of \(\eta_{i,j}(T)\) are responsible for the total susceptibility (12) around the characteristic temperature \(T^*\).

We note these components by \(\eta^a\), \(\eta^b\), and \(\eta^c\), which are characterized by the eigenfunction \(\Psi_j = \langle \theta_1, \theta_2 | E_i \rangle \) and \(\Psi_i = \langle \theta_1, \theta_2 | E_j \rangle \) as follows:

\[
\eta^a; \quad \Psi_j = c_0 c_0(\theta_1)c_0(\theta_2), \quad \Psi_i = c_0 c_0(\theta_1)c_0(\theta_2),
\]

(14)

\[
\eta^b; \quad \Psi_j = c_0 c_0(\theta_1)c_0(\theta_2), \quad \Psi_i = c_0 c_0(\theta_1)c_0(\theta_2),
\]

(15)

\[
\eta^c; \quad \Psi_j = c_0 c_0(\theta_1)c_0(\theta_2), \quad \Psi_i = c_0 c_0(\theta_1)c_0(\theta_2).
\]

(16)

The alphabets superscribed on \(\eta\) correspond to three dipolar transitions labeled by (a)–(c) shown in Fig. 3. For example, the solid arrow of (a) in Fig. 3 connects the eigenstates \(\Psi_j\) and \(\Psi_i\) in Eq. (14).

For weak coupling \(E_D \ll E_K\), the solution of the Mathieu equation (9) is easily solved. In the lowest order of the perturbation theory, the eigenvalues \(E_i(i=1,2)\) read

\[
E_i = \frac{E_K}{2} n^2 + a E_D \quad (n = 0, \pm 1, \pm 2, \ldots)
\]

(17)

with a constant \(a\). The solution (17) gives the eigenenergies of the states \(\langle E_i \rangle\) and \(\langle E_i \rangle\) relevant to the three components as follows:

\[
\eta^a; \quad (E_j,E_i) = (0,E_K + \delta E_a),
\]

(18)

\[
\eta^b; \quad (E_j,E_i) = \left(\frac{E_K}{2} \pm \delta E_b, \frac{E_K}{2} \pm \delta E_b\right),
\]

(19)

\[
\eta^c; \quad (E_j,E_i) = \left(\frac{5E_K}{2} \pm \delta E_c, \frac{5E_K}{2} \pm \delta E_c\right).
\]

(20)

The small corrections \(\delta E\) stem from the small interaction energy \(E_D \ll E_K\). Substituting these eigenenergies into Eq. (13), we find that the three components are approximated by

\[
\eta^a(u) = \frac{2p_a^2}{E_K} \left(1 - e^{-u}\right)/Z(u),
\]

(21)

\[
\eta^b(u) = \frac{2p_b^2}{E_K} \left(1 - e^{-u}\right)/Z(u),
\]

(22)

\[
\eta^c(u) = \frac{2p_c^2}{E_K} \left(1 - e^{-u}\right)/Z(u),
\]

(23)

where we define \(u = E_K/(k_B T)\). The quantities \(p_a\), \(p_b\), and \(p_c\) become equal to the value of \(\langle |E_j| |p_a| E_i \rangle\) for the case of \(\eta^a\), \(\eta^b\), and \(\eta^c\), respectively. The explicit form of the partition function \(Z(u)\) is

\[
Z(u) = 1 + 4 e^{-u} + 4 e^{-u} + 4 e^{-u} + 8 e^{-5u/2} + 4 e^{-4u} + 4 e^{-5u/2} + 8 e^{-5u} + \ldots,
\]

(24)

which monotonically decreases with rising \(u\) and reaches unity for the limit \(u \to \infty\). This means that the component \(\eta^a(u)\) is a monotonically increasing function of \(u\). On the other hand, the components \(\eta^b(u)\) and \(\eta^c(u)\) are convex functions giving a maximum at finite \(u\). The conditions of \(u\) for the maximum of \(\eta^b\) and \(\eta^c\) are expressed by

\[
1 - \frac{u}{2} - \frac{Z'(u)}{Z(u)} = 0, \quad \text{for} \ \eta^b,
\]

(25)

\[
1 - \frac{5}{2} \frac{u - Z'(u)}{Z(u)} = 0, \quad \text{for} \ \eta^c.
\]

(26)

The solutions of Eqs. (25) and (26) are estimated as \(u \approx 4\) for \(\eta^b\) and \(u \approx 0.5\) for \(\eta^c\). Since the total susceptibility \(\chi(T)\) is given by the summation \(\eta^a + \eta^b + \eta^c\), it is expected that the convex features of \(\eta^a(u)\) and \(\eta^c(u)\) cause the bump of the total susceptibility at \(u \approx 5\) shown in Fig. 4.

The argument is clarified by the numerical results shown in the inset of Fig. 4, where the \(u\) dependence of the components for \(E_D/E_K = 0.1\) are displayed: \(\eta^a\) (dashed-dotted), \(\eta^b\) (dotted), \(\eta^c\) (dashed-dotted-dotted), together with that of the total susceptibility \(\chi(u)\) that is negligible due to the factor \(e^{-u}\) in Eq. (23). As a result, the summation \(\eta^a(u) + \eta^b(u)\) shows a bump at \(u = 5.0\), which is the origin of the anomalous bump of the total susceptibility \(\chi(T)\) at the characteristic temperature \(T^* = E_K/k_B \times 0.2\). We should note here that, if quantum rotors are not interacting at all, the component \(\eta^b\) exactly vanished due to the degeneracy \(E_j = E_i = E_K^2/2\) [see Eq. (19)], and only the component \(\eta^c\) is dominant for the total susceptibility \(\chi(T)\). This means that the total susceptibility is a monotonic function same as \(\eta^c\) so that the bump does not emerge. The anomalous bump of the susceptibility, therefore, manifests the relevance of the dipolar interaction to the dielectric response of quantum rotors.

V. CONCLUSIONS

It is important to recall experiments reported in Ref. 18 for the dielectric susceptibility of KCl crystals with Li defects. It has been found that the susceptibility does not scale linearly with the Li concentration, and even becomes smaller with increasing concentration \((\approx 1000 \text{ ppm})\), where the in-
teraction between defects becomes relevant. In addition, a bump of the susceptibility is observed at about 200 mK for concentrations of 200–1000 ppm. These temperature dependences of the susceptibility together with the bumps are recovered well by our results shown in Fig. 4. Noting that defects in both systems move along closed loops and correlate each other, it is natural to assume that the similar picture holds. For a quantitative discussion, of course, one should take into account the effect of potential variation. For a quantitative discussion, of course, one should take into account the effect of potential variation hindering the free rotation of Li$^+$, which is caused by the Coulomb interaction between a mobile Li$^+$ ion and the host atoms K$^+$ and Cl$^-$. The problem has been theoretically investigated in Refs. 19 and 20 based on the two-level tunneling model.

In conclusion, we have investigated the quantum dynamics of two dipole rotors coupled via dipolar interaction. By solving analytically the eigenvalue problem of coupled rotors, we have demonstrated the energy spectra of coupled rotors as a function of dipolar interaction. The anomalous temperature dependence of dielectric susceptibility is also shown. Our model is so general that it should be applicable in a variety of physical context relevant to quantum rotors.

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1 J.K. Gimzewski, C. Joachim, R.R. Schlitvler, V. Langlais, H. Tang, and I. Johanssen, Science 281, 531 (1998).
2 N. Koumura, R.W.J. Zijlstra, R.A. van Delden, N. Harada, and B.L. Feringa, Nature (London) 401, 152 (1999).
3 V. Bermudez, N. Kapron, T. Gase, F.G. Gatti, F. Kajzar, D.A. Leigh, F. Zerbetto, and S. Zhang, Nature (London) 406, 608 (2000).
4 V.M. Rozenbaum and V.M. Ogenko, Fiz. Tverd. Tela (Leningrad) 30, 3048 (1988) [Sov. Phys. Solid State 30, 1753 (1988)].
5 V.M. Rozenbaum, V.M. Ogenko, and A.A. Chuiko, Usp. Fiz. Nauk 161, 79 (1991) [Sov. Phys. Usp. 34, 883 (1991)].
6 E. Artacho, F. Yndurain, B. Pajot, R. Ramirez, C.P. Herrero, L.I. Khirunenko, K.M. Itoh, and E.E. Haller, Phys. Rev. B 56, 3820 (1997).
7 M. Gienger, M. Glaser, and K. Lassmann, Solid State Commun. 86, 285 (1993).
8 W. Wegener, C. Bostoen, and G. Coddens, J. Phys.: Condens. Matter 2, 3177 (1990).
9 G.J. Kearley, H.G. Büttner, F. Fillaux, and M.F. Lautiè, Physica B 226B, 199 (1996).
10 P. Vorderwisch, S. Hautecler, G.J. Kearley, and F. Kubanek, Chem. Phys. 261, 157 (2000).
11 O. Rogalsky, P. Vorderwisch, A. Hüller, and S. Hautecler, J. Chem. Phys. 116, 1063 (2002).
12 A. Würg, Phys. Rev. Lett. 88, 063002 (2002).
13 W. Press, Single-Particle Rotations in Molecular Crystals, Springer Tracts in Modern Physics Vol. 92 (Springer, New York, 1981).
14 C. Gutt, B. Asmussen, W. Press, C. Merkl, H. Casalta, J. Greinert, G. Bohrmann, J.S. Tse, and A. Hüller, Europhys. Lett. 48, 269 (1999).
15 C. Gutt, W. Press, A. Hüller, J.S. Tse, and H. Casalta, J. Chem. Phys. 114, 4160 (2001).
16 S. Kettemann, P. Fulde, and P. Strehlow, Phys. Rev. Lett. 83, 4325 (1999).
17 The Mathieu equation was employed in Ref. 16 for calculating eigenenergies of single quantum rotor. The main difference between the Mathieu equation used in Ref. 16 and that in this paper is the physical origin of two potential minima expressed by the terms $\propto \cos(2\theta)$ in Eq. (9). The model of Ref. 16 does not take into account the interactions between rotors, namely, two potential minima for single rotors are presumed. Equation (9) was derived as a result of the dipole interaction, so the similarity is fortuitous, having no underlying physical origin.
18 C. Enss, M. Gaukler, S. Hunklinger, M. Tornow, R. Weis, and A. Würg, Phys. Rev. B 53, 12 094 (1996).
19 O. Terzidis and A. Würg, J. Phys.: Condens. Matter 8, 7303 (1996).
20 A. Würg, From Coherent Tunneling to Relaxation, Springer Tracts in Modern Physics Vol. 135 (Springer, Berlin, 1997).