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Research Article

Energy Exchange Using Spin-Boson Models with Infinite Loss

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Abstract

From experiment we know that energy is produced in the Fleischmann–Pons experiment, and that commensurate ^{4}He is observed, with about 24 MeV per helium atom. From the absence of neutrons in experiments producing excess heat, we know that the helium nuclei are born with less than 10 keV energy on average. This results in the key theoretical problem associated with the Fleischmann–Pons experiment: where does the energy go? In the lossy spin-boson model, a large energy quantum is converted into many small quanta. Here we present a new analysis of the lossy spin-boson model.

Keywords: Excess power, Fleischmann–Pons effect, Lossy-spin boson model, Phonon exchange

1. Introduction

Since the first days after the initial announcement of Fleischmann and Pons in 1989 [1,2], we have been drawn to the problem of how an energy quantum at the nuclear scale (MeV) is split up into a very large number of smaller quanta at the atomic scale (meV). The basic problem which received much attention in the early years was the absence of deuteron–deuteron fusion neutrons in amounts commensurate with the energy observed, which was taken as evidence that no commensurate deuteron–deuteron fusion reactions occurred [3]. We now recognize that this observation is symptomatic of a much larger version of the problem. Cells producing excess heat have in some cases been monitored for neutrons, gammas, and X-rays with the result that very little radiation of any kind is observed correlated with excess power [4]. Consequently, we are driven to the conclusion that there are essentially no energetic particles of any kind produced commensurate with the energy observed.

To date, only one product has been shown to be correlated with excess power in amounts commensurate with the energy: ^{4}He atoms in the gas, with about one helium atom observed for every 24 MeV of energy produced [5–7]. In light of the comments above, we became interested in the question of how much energy the ^{4}He atom is born with,

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in order to be consistent with experimental observations. For example, there have been proposals that two deuterons come together in the vicinity of a Pd nucleus, creating a ^{4}He nucleus which is born with nearly 24 MeV following a recoil with the Pd (as commented on in [4], and similar to the mechanism discussed in [8]). In another proposal, four deuterons come together to make two ^{4}He nuclei, each of which carries away about 24 MeV in kinetic energy [9]. Since energetic ^{4}He nuclei would subsequently collide with other nuclei [10], we should be able to determine from experiment on average how much energy that the ^{4}He nuclei are born with from observations of secondary radiation correlated with excess power.

We carried out a set of simple calculations for neutron, X-ray, and gamma-ray yields for a wide range of initial ^{4}He energies [4]. We found that the most sensitive diagnostic is secondary neutron production from deuteron–deuteron fusion following a hard collision of a ^{4}He nucleus with a deuteron. We then looked through a large number of papers to find experiments where excess power was observed under conditions where neutron emission was monitored. A modest number of such experiments have been reported, and in the majority of these experiments there is no correlation between neutron emission and excess power. From an analysis of the experimental papers, it was possible to estimate upper limits for the number of source neutrons produced per unit energy to be between 0.008 and 0.8 n/J, which correspond to initial ^{4}He energies between 6.2 and 20.2 keV [4].

We can conclude from this that theoretical proposals like those mentioned above, which result in energetic ^{4}He nuclei, are inconsistent with experiment. A natural response might be to suggest that perhaps the ^{4}He nuclei pushes off a lighter particle, resulting in less kinetic energy (a neutrino in [11], and an electron in [8]). In this case, the lighter particle will carry most of the energy. We can get a rough estimate of the minimum ^{4}He recoil energy for any light particle by estimating the recoil energy in the case of a photon. Assuming that momentum is conserved, the ^{4}He nucleus would have a momentum of roughly

$$P = \frac{\hbar\omega}{c} \tag{1}$$

and a corresponding recoil energy of

$$E = \frac{P^2}{2M} = \frac{(\hbar\omega)^2}{2Mc^2} = 76 \text{ keV}. \tag{2}$$

As a result, we are able to rule out all proposed mechanisms in which the end result is ^{4}He and another particle with energy and momentum conserved, as inconsistent with experiment (since the secondary neutron yield for 76 keV alphas should have been seen).

1.1. Connection with theory

This argument is powerful, and it allows us to make progress on theory in several ways. If we know a certain class of models is inconsistent with experiment, then we can use our time more productively looking at models that are not inconsistent with experiment. In order to develop models that can be consistent with experiment, we need to understand new kinds of reaction mechanisms that work differently from the simple Rutherford billiard ball model, in which local conservation of momentum and energy dictates that product particles must be energetic. People have thought about this problem for several years, and we are in a position to summarize some of the different lines of thought that have resulted.

In light of the comments above, there seem to be three basic possibilities: (a) we can imagine a scenario in which the energy produced is somehow spread among a large number of particles initially, as a new kind of many-particle effect [12,13]; (b) perhaps the reaction initially creates a small number of energetic particles, which then slow down (in a new way that does not lead to secondary radiation) and exchange energy with a large number of particles; or (c) the reaction energy is transferred initially in some way to a phonon mode, or some other low energy degree of freedom. Each of these possibilities individually would constitute a revolution in nuclear physics and condensed matter, if real.

At present, there is no known physical system that works as (a), in which direct coupling of a large reaction energy occurs between one system and a very large number of particles. The number of particles under discussion in the case of deuterons would be 24 MeV divided by the associated upper limit from experiment combined with a yield calculation, which in this case is under 1000 eV [14]. Hence, the first possibility imagines the direct coupling of 24 MeV to more than 24,000 deuterons. It may be that such energy redistribution might be possible in a tightly coupled system with much stronger particle–particle interactions, but it is very hard to see how this scenario might work in the PdD environment. While this kind of mechanism is discussed by Kim in [12,13], a calculation that demonstrates that such an effect can occur at a finite rate has to our knowledge not yet been given. We might also consider energy transfer to the Pd subsystem since it is coupled much more strongly; but even so, the Pd–Pd interactions are very weak compared to what would be required for direct distribution of 24 MeV.

The possibility that new physics might be involved in the slowing down of an energetic particle as in (b) can be contemplated. Working against this kind of approach is the problem of obtaining consistency with observations of secondary radiation, specifically secondary neutrons. The basic problem is that Rutherford scattering between a charged energetic particle (such as an alpha particle) and a deuteron is very efficient at producing energetic deuterons, which then have a high probability of subsequently producing neutrons from deuteron–deuteron fusion reactions. To be consistent with the observed lack of secondary neutron emission for MeV alphas, the secondary neutron yield would have to be reduced by 4–5 orders of magnitude [4]. If the initial alpha energy were 10 MeV or more, then the predicted yield of neutrons from deuteron break up reactions is nearly 10 orders of magnitude higher than observed. Experiments that have focused on the slowing down of charged particles over the past seventy years have unfortunately not provided evidence for the kind of large enhancement that would be required.

We have focused over the past decade or more on models in which the energy is transferred to low energy degrees of freedom as in (c). Working against the approach is the fact that there are no previous observations of similar effects in the nuclear physics literature. However, a somewhat similar effect is known in other areas. If we consider the basic problem to be one of converting a large energy quantum associated with a two-level system (representing nuclear energy levels) into a large number of low-energy oscillator quanta (representing a phonon mode), then there is the beginning of a precedent in the NMR and atomic physics literature [15]. In the multiphoton regime of the spin-boson model [16], coherent energy exchange occurs between two-level systems with a large transition energy and an oscillator with a low characteristic energy [17,18]. This basic effect has been observed experimentally [19]. The problem that we face is that this anomalous coherent energy exchange effect is very weak in the spin-boson model; coherent energy exchange is possible only when the large quantum is divided into a few tens of low energy quanta. In order to split a very large MeV quantum into a much larger number of oscillator quanta, we require a model something like the spin-boson model, but this new model must have a much stronger version of the effect to be relevant to the Fleischmann–Pons excess heat observations.

In the sections that follow, we will consider the lossy spin-boson model in the limit of infinite loss [20–22]. The spin-boson model augmented with loss is capable of a much stronger anomalous energy conversion effect than the conventional spin-boson model. The finite loss version of the model in our view is relevant to experiment, but it is complicated; if we make the loss infinitely fast, then the resulting model simplifies considerably, and allows us to quantify it in a useful way.

2. Spin-boson model

In the basic spin-boson model, a two-level system is coupled to an oscillator with linear coupling (the two-level system can make a transition up or down, coupled with an increase or decrease of one oscillator quantum). This model is very popular in the atomic physics literature (several hundred papers have been published on it), mostly because it provides a test problem for new mathematical techniques. The following Hamiltonian describes the model [16]

$$\hat{H} = \Delta E \frac{\hat{S}_z}{\hbar} + \hbar\omega_0 \hat{a}^\dagger \hat{a} + U \left(\hat{a}^\dagger + \hat{a}\right) \frac{2\hat{S}_x}{\hbar}. \tag{3}$$

The first term describes the two-level system energy with transition energy $\triangle E$, the second describes the oscillator with a characteristic frequency of ω_0, and the last term accounts for linear coupling between the two systems.

In the multiphoton regime (where the characteristic energy of the oscillator is much less than the transition energy of the two-level system), the states that are coupled in the Hamiltonian are off of resonance, since the energies of the two individual systems are highly mismatched. Our intuition might suggest that no energy exchange is possible under such conditions, since the energy levels of the two systems are not matched. Nevertheless, many analyses carried out over the past 50 years have shown that coherent energy exchange can occur. The reason for this is that there can be a weak indirect coupling between distant states that are resonant with each other, resulting in a finite rate for coherent energy exchange as long as the two-level system energy (including the coupling) is matched to an odd number of oscillator quanta [17,18].

2.1. Nuclear levels and phonon mode

We might think about this model in relation to the Fleischmann–Pons excess heat effect. In this case, the two-level system might represent nuclear energy levels; the ground state of the two-level system might stand in for the ground state of a nucleus, and the excited state of the two-level system might stand in for an excited nuclear state. The oscillator in the two-level system might represent one phonon mode.

It is possible to use a spin-boson model to investigate what happens when these two systems are coupled. For the model to be relevant there needs to be some coupling so that a transition between the two nuclear states can occur if a phonon is exchanged. In general we would expect the nuclear states not to care whether the phonon modes (which the nucleus is part of) are excited or not. There are, however, some special cases where we would expect coupling to occur [23]. A phonon mode is going to care about the mass of a nucleus; hence if a nucleus is made up of a superposition of states with different masses, then we might expect that phonon exchange might occur if the mode is highly excited in association with transitions between these different states.. Alternatively, if there is a D_2 molecule embedded in the lattice, a phonon mode is going to notice if it makes a transition to a ^{4}He atom. In both of these cases, it would be plausible to use a spin-boson model to begin analyzing what happens if phonon exchange becomes possible.

Intuitively, we would not expect much to happen in either case. When off-resonant states mix, we can usually use second-order perturbation theory to compute the level shifts, and how much mixing occurs. The spin-boson model in this application would predict a shift of the nuclear levels sufficiently small as to be unmeasurable. The probability that the nuclear transition energy could be converted into vibrational energy in the spin-boson model for this example can be evaluated directly, and the result is that no coherent energy exchange can take place. We conclude that within the context of the spin-boson model nothing interesting happens.

3. Lossy Spin-Boson Model

We have argued for several years that when the spin-boson model is augmented with loss that things change qualitatively [20–22]. In this case, we are interested in the Hamiltonian

$$\hat{H} = \Delta E \frac{\hat{S}_z}{\hbar} + \hbar\omega_0 \hat{a}^\dagger \hat{a} + U \left(\hat{a}^\dagger + \hat{a}\right) \frac{2\hat{S}_x}{\hbar} - i\frac{\hbar\hat{\Gamma}(E)}{2}. \tag{4}$$

The first term in this case represents many two-level systems, the second term accounts for the oscillator as before, the third term describes linear coupling with the oscillator (which mediates transitions between all of the different two-level systems on equal footing), and the last term accounts for loss.

For coherent energy exchange in the multiphoton regime, we have shown explicitly that severe destructive interference occurs in the basic spin-boson model in the limit of weak coupling, resulting in a very slow associated rate. This destructive interference is lifted in the lossy version of the model, especially under conditions where the loss is very strong [21]. As such, the lossy version of the spin-boson model is very efficient at exchanging energy coherently under conditions where a large two-level system quantum is split up into many small oscillator quanta; it constitutes a new kind of model in this regard that we have not encountered previously in physics.

3.1. Coupling between nearly degenerate states

There is a joke in which a mathematician, a physicist, and an engineer each sees a fire breaking out; the mathematician realizes that a solution exists, and walks away; the physicist reduces it to a previously solved problem, and also walks away; and the engineer rolls up his sleeves and puts the fire out. In the sense of this joke, as applied physicists, we would like to reduce the problem to one previously solved. So, our focus then is finding some well known previously solved problem that is like this one, which perhaps would allow us to have some intuition about how the new model works.

As it turns out, there is a classic physics problem, which is something like the new model, that allows us to understand the new model simply. In the lossy spin-boson model, we are interested in the coherent dynamics which govern transitions between nearly degenerate states. But almost the same mathematical problem occurs when we consider the coherent dynamics of an electron in the conduction band of a crystal semiconductor. We can use our intuition about this well known problem then in order to understand how the lossy spin-boson model works.

Consider for example an idealized set of atoms next to one another in one dimension, each with the same energy level, with the possibility of electron tunneling between neighbors. The evolution equation for the probability amplitudes $c_n(t)$ at each site can be written as

$$i\hbar \frac{\mathrm{d}}{\mathrm{d}t} c_n(t) = -V c_{n-1}(t) - V c_{n+1}(t). \tag{5}$$

This, by itself, is almost enough for us to recognize that the electron will evolve pretty much as a free particle with an effective mass. We might develop an energy band by working with eigenfunctions of the problem

$$c_n(t) = \mathrm{e}^{-iEt/\hbar} \mathrm{e}^{ikna} \tag{6}$$

and then substituting to obtain

$$E = -2V \cos(ka), \tag{7}$$

where a is the lattice constant for this idealized model. The conduction band in this case is locally parabolic around $k = 0$, so that the dispersion relation is almost the same as a particle in free space. We can approximate the discrete probability amplitudes with a continuous wavefunction

$$\begin{aligned} c_n(t) &\to \psi(na, t) \\ c_{n\pm 1}(t) &\to \psi((n \pm 1)a, t) \to \psi(na, t) \pm a \left(\frac{\partial \psi}{\partial x} \right)_{x=na} + \frac{a^2}{2} \left(\frac{\partial^2 \psi}{\partial x^2} \right)_{x=na} + \cdots \end{aligned} \tag{8}$$

which allows us to approximate the evolution equation as a Schrodinger equation

$$i\hbar \frac{\partial}{\partial t} \psi(x, t) = -2V \psi(x, t) - V a^2 \frac{\partial^2}{\partial x^2} \psi(x, t). \tag{9}$$

The effective mass in this model can be obtained from

$$\frac{\hbar^2}{2m^*} = a^2 V. \tag{10}$$

The continuum limit of this problem is that of a free particle in a constant potential. This allows us to use our intuition about free particle dynamics to understand the more complicated problem of coupling between a set of ordered degenerate states. Because we have so much intuition about this problem, we would have little problem extending our intuition if there was an additional potential that depended on the position n, or if the coupling coefficients varied slowly with n.

3.2. Lossy spin-boson model as a nearly degenerate state problem

To proceed, we would like to find a way to recast the lossy spin-boson model into a form in which it looks like a set of ordered nearly degenerate states with coupling between neighbors. In the absence of coupling, the states are not degenerate; we may write them as

$$E_{M,n} = M\Delta E + n\hbar\omega_0. \tag{11}$$

In general there are many states, and only a few of them might have roughly the same energy. For example, if the transition energy were exactly 11 times larger than the oscillator energy, then we would have the same energy if we increase the number of oscillator quanta n by 11 each time we decreased the number of excited states of the two-level systems by 1. But then this would leave us with all kinds of states that have different energies, and which are not degenerate.

We can accomplish this by using a second-order formulation. We separate out all of the states that are nearly degenerate, and focus on them; all of the other states we eliminate algebraically. For example, the eigenvalue equations for two sets of expansion coefficients are

$$\begin{aligned} E\mathbf{c} &= \mathbf{H}_a \cdot \mathbf{c} + \mathbf{V}_{ab} \cdot \mathbf{d}, \\ E\mathbf{d} &= \mathbf{H}_b \cdot \mathbf{d} + \mathbf{V}_{ba} \cdot \mathbf{c}, \end{aligned} \tag{12}$$

where $\mathbf{c}$ is a vector of coefficients of the degenerate states that we want to focus on, and $\mathbf{d}$ is a vector of coefficients for all of the other states. We can focus on the $\mathbf{c}$ vector by eliminating the $\mathbf{d}$ vector to obtain

$$E\mathbf{c} = \mathbf{H}_a \cdot \mathbf{c} + \mathbf{V}_{ab} \cdot (E - \mathbf{H}_b)^{-1} \cdot \ \mathbf{V}_{ba} \cdot \mathbf{c}. \tag{13}$$

This problem now describes an ordered set of nearly degenerate states with coupling between neighbors. The indirect coupling between neighbors is included naturally in this formulation. All that remains is to understand the shifts and couplings of the resulting nearly degenerate states.

4. Results

We have carried out a variety of calculations using this general approach, and the results are interesting. There is a scaling law apparent within this model that we found from our first few calculations, which will help us in what follows to understand the results more simply. Our focus in what follows is on the special case where the loss is infinite for all basis states with an energy less than a cut-off energy we selected.

4.1. Scaling

There are a relatively small number of parameters that appear in the augmented spin-boson model in the limit of infinite loss, which will make it easier for us to understand. Perhaps the most important parameter for our discussion is the ratio of the transition energy to the oscillator energy, which will appear in our discussion in connection with the number of oscillator quanta Δn which is exchanged for a two-level system quantum. The strength of the coupling between the two systems is also important, and we have found that a useful measure of this can be developed through the dimensionless coupling parameter g defined according to

$$g = \frac{\max\left(\left\langle i|\hat{H}|f\right\rangle\right)}{\Delta E} \rightarrow \frac{(N+1)U\sqrt{n_0}}{2\Delta E}, \tag{14}$$

where N is the number of two-level systems. The number of two-level systems by itself constitutes an important additional parameter of the model.

One way to think about scaling in this model is to start with a problem in which the number of quanta Δn to be transferred is fixed, and then imagine that we would like to see what happens when we vary the number of two-level systems N or the coupling strength g. If we fix g and look at models with different N, we find that the energy offsets of the nearly degenerate states are very similar. As we increase the number of two-level systems, then the states become more and more nearly degenerate as the maximum offset remains constant. Also, the maximum indirect coupling between nearly degenerate states remains approximately fixed in this limit. Hence, if we compute results for one such model, the answer will be very similar to other models with different values of N.

We have found another concept that is very helpful in understanding these models. The details of the local coupling between the nearly degenerate states seem to depend only on the local coupling strength of the original problem when N is large. As a result, we can assume that all of the coupling matrix elements are the same locally (as a function of the number of excited two-level systems) in the initial Hamiltonian, to get approximate values for the shifts and indirect coupling matrix elements for the nearly degenerate states. We find that these approximate values are very close to the exact ones computed for the full problem. Since this seems to work, it allows us to focus on the level shifts and indirect coupling matrix elements as a function of g for this simpler uniform problem. Instead of having to analyze each version of the model separately, all we really need is to characterize the shift and indirect coupling matrix elements as a function of the dimensionless coupling strength g to understand most of what is important about the model.

4.2. Indirect coupling between neighboring states

By far the most important result from this model is the indirect coupling strength between neighboring nearly degenerate states. Results are shown in Figure 1 for the scaled indirect coupling matrix element as a function of the dimensionless coupling strength g.

Several features of these results are of interest to us here. To begin with, these indirect coupling coefficients are important for coherent energy exchange between the two-level systems and the oscillator, because the maximum rate at which two-level system transitions occur is proportional to this matrix element. When we scale the matrix element by the number of oscillator quanta exchanged Δn, the resulting scaled matrix element is proportional to the rate at which individual oscillator quanta are exchanged.

When g is small, the coupling is weak, and the matrix elements scale as

$$\frac{V_{\text{eff}}}{\Delta E} \rightarrow A(\Delta n)g^{\Delta n}. \tag{15}$$

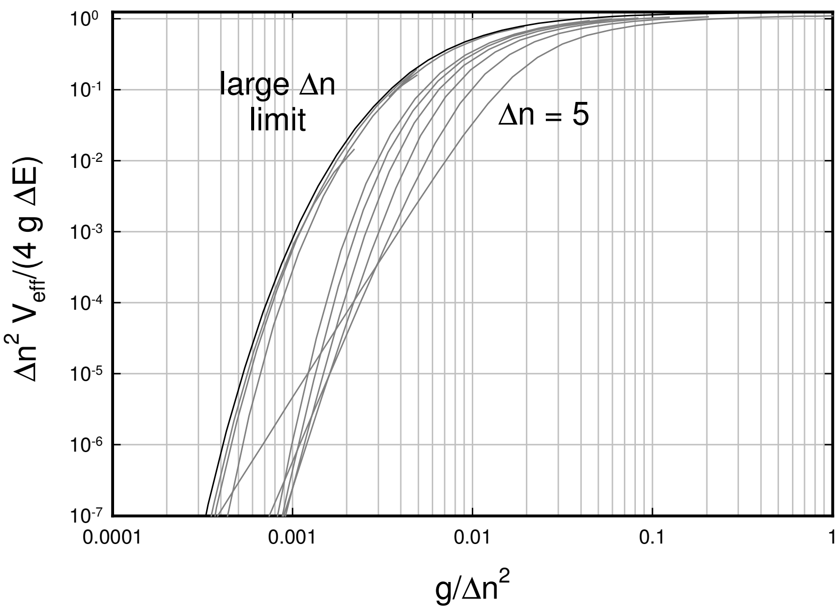

Figure 1. Scaled indirect coupling matrix element as a function of the scaled dimensionless coupling parameter g for different odd Δn, including 5,7,9,11,15,51,101,151 (grey); also shown is the large Δn limit (black) .

From perturbation theory we know that $A(3)$ is 9/8, in agreement with the numerical calculations. For larger values of Δn, an approximate fit is given by

$$A(\Delta n) = a\mathrm{e}^{b\Delta n} \tag{16}$$

with $a = 0.3139\Delta E$, and $b = 0.3441$. When g is somewhat above unity for modest Δn, the nearest neighbor indirect coupling matrix element is roughly

$$V_{\text{eff}} \to 4.8\frac{g\Delta E}{\Delta n^2} = 4.8\frac{g\hbar\omega_0}{\Delta n}. \tag{17}$$

We have found that in the limit of large Δn, it is possible to express the indirect coupling matrix element in terms of a universal function according to

$$\frac{\Delta n^2 V_{\text{eff}}}{4g\Delta E} = \Phi\left(\frac{g}{\Delta n^2}\right), \tag{18}$$

where the universal function Φ can be fit according to

$$\ln\Phi(y) = \frac{a_{-3}}{y^3} + \frac{a_{-2}}{y^2} + \frac{a_{-1}}{y} + a_0 \tag{19}$$

with the fitting parameters

$$\begin{aligned} a_{-3} &= -2.11317\times 10^{-10}, \qquad & a_{-2} &= 1.8517\times 10^{-6},\\ a_{-1} &= -0.00899174, & a_0 &= 0.208444. \end{aligned} \tag{20}$$

This result is plotted in Fig. 1 to compare with the finite Δn results.

In connection with coherent energy exchange in the Fleischmann–Pons experiment, this result begins to clarify some of the issues. In this case Δn is very large (on the order of 10^8–10^9). Consequently, we expect that when g is small, the indirect coupling matrix element will be essentially zero, consistent with no coherent energy exchange. We would not expect any coherent energy exchange below $g/\Delta n^2 = 10^{-4}$ for large Δn.

5. Discussion

In the Fleischmann–Pons experiment, energy is produced by some new physical process that results in ^{4}He as a reaction product. Nuclear reactions as understood in the nuclear physics literature lead to energetic reaction products as a consequence of Rutherford-style conservation of energy and momentum. Due to the absence of energetic reaction products observed in the Fleischmann–Pons experiment, we know that the ^{4}He is born with sufficiently low energy so as to be inconsistent with any reasonable extrapolation of Rutherford reaction kinetics. Based on this kind of consideration alone, we can conclude that our attention should be focused on new reaction schemes that split up the nuclear energy quantum into a very large number of small quanta. At present, there seem to be two basic ways that this might be done. One approach is being pursued by Kim [12,13], where the reaction energy is proposed to be converted to kinetic energy of a large number of deuterons in a way that remains to be clarified. The other approach is through a coherent energy exchange mechanism with a condensed matter degree of freedom, such as a phonon or plasmon mode.

Many years ago we noticed that models with coupled two-level systems and an oscillator, augmented with loss, could dramatically improve the rate at which coherent energy exchange occurs in the multiphonon regime [24]. Over the years we have explored this effect using a number of different tools, including quantum flow calculations, periodic lattice models, and coupling between nearly degenerate states. In the case of the quantum flow calculations, we were able to demonstrate coherence under conditions where up to 10^4 phonons were exchanged by brute force numerical calculations. The periodic lattice calculations are useful for quantifying the threshold effect.

A few years ago, we made an effort to connect our approach with results in the conventional literature on coherent energy exchange in the spin-boson model [17,18]. In this case, we were able to develop a rotation that allowed the problem to be analyzed as a set of nearly degenerate states with dominant nearest neighbor coupling. By viewing the lossless version of the problem in this way, we were able to analyze coherent energy exchange in the multiphonon limit systematically for weak and moderate coupling.

We have been interested in recent years in extending this style of analysis to the lossy version of the problem. In the results discussed briefly in this paper, we have succeeded. Perhaps the most important feature of this kind of analysis is that we have identified scaling laws that make the system much simpler to understand. As a result, we are able to present results for the indirect coupling matrix elements systematically for models with different numbers of exchanged phonons. These indirect matrix elements determine nearly everything about how fast coherent energy exchange occurs in the models.

We have not emphasized the dramatic increase in the rate of coherent energy exchange with loss as compared to the lossless case, but this issue deserves some comment here. One way to see this is in the definition of the dimensionless coupling parameter g for the two problems: the definition of g appropriate for the lossy problem is larger essentially by the number of two-level systems than the definition in the lossless case. This means that in the lossless problem, the energy coupling effects act the same if one two-level system is coupled to the oscillator, or if one million two-level systems are coupled. In the lossy version of the problem, the presence of the additional two-level systems increases the dimensionless coupling strength in proportion to the number of them. Hence it is much easier for a physical system to act like a strongly coupled system if it behaves similarly to the spin-boson model augmented with loss.

In a sense, this captures what is special about the model. We know that only a strongly coupled system is going to be capable of splitting up a large quantum into a very large number of small quanta. The physical coupling between the atoms in PdD does not seem to be strong enough by at least six orders of magnitude to do the job. In the lossless spin-boson model, the energy exchange process is nearly the same as the coupling between one two-level system and the oscillator. Not only would we need the transition matrix element for single phonon exchange to be on the order of 24 MeV to get into the strong coupling limit, but once there we would at most be able to split a 24 MeV quantum into smaller quanta on the order of 1 MeV. However, in the spin-boson model augmented with loss, whether the coupling is strong (or not) is determined in part by the magnitude of the phonon exchange matrix element, and in part by the number of two-level systems involved (and their excitation). Hence, we can have a strongly coupled system in this case if the individual local interaction matrix element is much less than 24 MeV, as long as the effect of all of the individual transitions in the system add up in phase (which occurs in this kind of model if the phonon mode wavelength is long), and if there are a lot of them.

We have focused on the spin-boson model augmented with loss in this paper since the energy exchange part of the problem is clearly the most important. Our overall model for excess heat in the Fleischmann–Pons experiment involves two sets of two-level systems instead of one, so the associated model is more complicated. In the event that the excitation transfer step is the bottleneck, then this step limits the reaction rate for the overall model. Our calculation of this transition matrix element seems to be consistent with the transition rates needed to account quantitatively for the observed rates of excess power. The related computation of the energy exchange rate has been for us much more difficult, since it has become clear than pretty much all excited states which can be coupled to be consistent with the requirement that phonon exchange occurs.

Some of our colleagues have expressed displeasure with the model, arguing that any significant nuclear excitation in such a model would be lost immediately. A feature of the model worthy of comment in this regard is that the indirect coupling between states in this model occurs through off-resonant states that are in energy deficit. Although it remains to be demonstrated explicitly that this results in the stabilization of these states, the associated intuition is that the system is going to have trouble decaying if it does not have enough energy to decay.

Other colleagues have misunderstood energy exchange in this kind of model, talking about a very energetic phonon being emitted. Such language is unhelpful. Instead, the way this model works is that the coupled system undergoes a very large number of rapid sequential excitation and de-excitation steps, exchanging one low energy oscillator quantum at a time; this occurs in a coherent process similar to that already described in the spin-boson model, and observed in experiment.

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