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Thermodynamic versus statistical nonequivalence of ensembles for the mean-field Blume-Emery-Griffiths model

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We illustrate a novel characterization of nonequivalent statistical mechanical ensembles using the mean-field Blume-Emery-Griffiths (BEG) model as a test model. The novel characterization takes effect at the level of the microcanonical and canonical equilibrium distributions of states. For this reason it may be viewed as a *statistical* characterization of nonequivalent ensembles which extends and complements the common *thermodynamic* characterization of nonequivalent ensembles based on nonconcave anomalies of the microcanonical entropy. By computing numerically both the microcanonical and canonical sets of equilibrium distributions of states of the BEG model, we show that for values of the mean energy where the microcanonical entropy is nonconcave, the microcanonical distributions of states are nowhere realized in the canonical ensemble. Moreover, we show that for values of the mean energy where the microcanonical entropy is strictly concave, the equilibrium microcanonical distributions of states can be put in one-to-one correspondence with equivalent canonical equilibrium distributions of states. Our numerical computations illustrate general results relating thermodynamic and statistical equivalence and nonequivalence of ensembles proved by Ellis, Haven, and Turkington [*J. Stat. Phys.* **101**, 999 (2000)].

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I. INTRODUCTION

The microcanonical and canonical ensembles are the two main probability distributions with respect to which the equilibrium properties of statistical mechanical models are calculated. As is well known, the microcanonical ensemble is a statistical mechanical expression of the conservation of the energy for a closed or isolated system whereas the canonical ensemble models a system in thermal interaction with a heat reservoir having a constant temperature [1, 2]. Although the two ensembles model two different physical situations, it is widely assumed that the ensembles give equivalent results in the thermodynamic limit; i.e., in the limit in which the volume of the system tends to infinity. The typical argument used to motivate this equivalence is given in the classic text of Landau and Lifshitz [3]. Although “the canonical distribution is ‘spread’ over a certain range of energies,” “the width of this range . . . is negligible for a macroscopic body.” The conclusion is that, in the thermodynamic limit, the canonical ensemble can be considered to be an ensemble of fixed mean energy or, in other words, a microcanonical ensemble. The texts [1, 4, 5, 6, 7] contain similar arguments.

In order to substantiate the argument of Landau and Lifshitz, it must be proved that the energy per particle converges to a constant in the infinite-volume limit of the canonical ensemble. This convergence can be proved to hold for noninteracting systems such as the perfect gas and for a variety of weakly interacting systems. For general systems, however, neither is this convergence valid nor is the conclusion true concerning ensemble equivalence which this convergence is intended to motivate. In fact, in the past three and a half decades, numerous sta-

tistical models have been discovered having microcanonical equilibrium properties that cannot be accounted for within the framework of the canonical ensemble. This lack of correspondence between the microcanonical and canonical ensembles has profound consequences for it implies that one is forbidden to substitute the mean-energy variable for the temperature variable, and vice versa, when parameterizing the equilibrium properties of systems. In such cases of nonequivalence, the questions of determining which of the two ensembles is the more fundamental and which ensemble is the one realized in the laboratory are of fundamental interest.

Until now, the phenomenon of nonequivalent ensembles has been identified and analyzed almost exclusively by determining regions of the mean energy where the microcanonical entropy function is anomalously nonconcave or by determining regions of the mean energy where the heat capacity, calculated microcanonically, is negative. This thermodynamic approach to the problem of nonequivalent ensembles has been propounded by a number of people, including Lynden-Bell and Wood [8], who in the 1960’s were among the first to observe negative heat capacities in certain gravitational many-body systems (see also [9, 10, 11, 12, 13, 14]). More recently, nonconcave anomalies in the microcanonical entropy as well as negative heat capacities have been observed in models of fluid turbulence [15, 16, 17, 18] and models of plasmas [19], in addition to long-range and mean-field spin models, including the mean-field XY model [20] and the mean-field Blume-Emery-Griffiths (BEG) model [21, 22].

The existence of such nonconcave anomalies invalidates yet another tacit principle of statistical mechanics which states that the one should always be able to express the microcanonical entropy, the basic thermodynamic func-

tion for the microcanonical ensemble, as the Legendre-Fenchel transform of the free energy, the basic thermodynamic function for the canonical ensemble. Indeed, if the microcanonical entropy is to be expressed as the Legendre-Fenchel transform of the canonical free energy, then the former function must necessarily be concave on its domain of definition. Hence, if the microcanonical entropy has nonconcave regions, then expressing it as a Legendre-Fenchel transform is impossible. When this occurs, we say that there is thermodynamic nonequivalence of ensembles.

For the BEG model, the existence of nonconcave anomalies in the microcanonical entropy was shown by Barré, Mukamel and Ruffo via a Landau expansion of this quantity [21]. The aim of the present paper is to extend their study of nonequivalent ensembles by showing how the thermodynamic nonequivalence of the microcanonical and canonical ensembles for this model reflects a deeper level of nonequivalence that takes place at the statistical level of its equilibrium distribution of states. We carry this out via numerical calculations both at the thermodynamic level and at the statistical level, illustrating, in particular, a striking statistical consequence of the nonconcavity of the microcanonical entropy. Namely, we demonstrate that if the microcanonical entropy is nonconcave at a mean-energy value u —i.e., if thermodynamic nonequivalence of ensembles holds—then nonequivalence of ensembles holds at the statistical level in the following sense: none of the equilibrium statistical distributions of states calculated in the microcanonical ensemble at that value of u can be realized in the canonical ensemble. Furthermore, we demonstrate that if the microcanonical entropy is strictly concave at u —i.e., if thermodynamic equivalence of ensembles holds—then all the statistical distributions of states found in the microcanonical ensemble at that value of u can be put in one-to-one correspondence with equivalent distributions of the canonical ensemble.

Our numerical findings illustrate clearly and directly a number of results on equivalence and nonequivalence of ensembles that are valid for a wide class of statistical mechanical systems. These results, derived recently by Ellis, Haven and Turkington, establish different levels of correspondence between microcanonical equilibrium macrostates and canonical equilibrium macrostates in terms of concavity properties of the microcanonical entropy. These results, as well as an overview of applications to turbulence, can be found in their comprehensive paper [23]. A specific application to a model of geophysical fluid turbulence is treated in their paper [18].

The organization of the present paper is as follows. In Section II we introduce the basic thermodynamic functions in the two ensembles, the microcanonical entropy and the canonical free energy. We also motivate the definitions of the sets of equilibrium statistical distributions of states for the two ensembles in the thermodynamic limit. This motivation is based in part on the theory of large deviations, a branch of probability that studies the

exponential decay of probabilities and is perfectly suited for analyzing asymptotic properties of the two ensembles. In Section III we discuss the equivalence and nonequivalence of ensembles first at the thermodynamic level and then at the statistical level. In Section IV, we finally present our numerical calculations that illustrate thermodynamic and statistical equivalence and nonequivalence of ensembles for the BEG model. Because our goal is to emphasize the physical ideas, we have omitted almost all mathematical details and have occasionally compromised perfect mathematical accuracy when it benefits the exposition. For complete mathematical details concerning equivalence and nonequivalence of ensembles and the theory of large deviations, the reader is referred to [23, 24, 25].

II. THERMODYNAMIC ENSEMBLES AND LARGE DEVIATIONS PRINCIPLES

A. Microcanonical ensemble

The main quantity characterizing the thermodynamic properties of statistical mechanical systems in the microcanonical ensemble is the entropy function. This quantity is defined in terms of the probability measure of all microstates of a system having the same value of the mean energy [1, 2]. To be precise, suppose that the system in question is composed of n particles, and denote the microstates of that system by the joint state $x^n = (x_1, x_2, \dots, x_n)$, where x_i represents the state of the i^{th} particle taking values in some state space \mathcal{X} . The set of all microstates is the n -fold product \mathcal{X}^n . Moreover, let us denote by $U(x^n)$ the energy or Hamiltonian of the system as a function of the microstates x^n and by $u(x^n)$ the mean energy or energy per particle, defined as

$$u(x^n) = \frac{U(x^n)}{n}. \quad (1)$$

In terms of this notation the microcanonical entropy function $s(u)$ is defined as

$$s(u) = \lim_{n \rightarrow \infty} \frac{1}{n} \ln P\{u(x^n) \in du\}, \quad (2)$$

where

$$P\{u(x^n) \in du\} = \int_{\{x^n: u(x^n) \in du\}} P(dx^n) \quad (3)$$

is the probability measure of all microstates x^n lying in the infinitesimal mean-energy ball du centered at u [23]. The probability measure $P(dx^n)$ in (3) is the *a priori measure* on \mathcal{X}^n , which is taken to be the uniform measure in accordance with Boltzmann's equiprobability hypothesis [1, 2]. In order for this probability measure to be well defined, the configuration space \mathcal{X}^n is assumed to be bounded. As we will soon explain, conditioning P on the

set of configurations $\{x^n : u(x^n) \in du\}$ defines the microcanonical ensemble. We call this set the *microcanonical conditioning set*.

We find it convenient to re-express the definition of $s(u)$ in (2) via the formula

$$P\{u(x^n) \in du\} \asymp e^{ns(u)}. \quad (4)$$

This is done in order to emphasize the facts that $P\{u(x^n) \in du\}$ has, to a first degree of approximation, the form of an exponential that decays with n and that the exponential decay rate is the microcanonical entropy. The heuristic sign “ \asymp ” is used here instead of the approximation sign “ \approx ” in order to stress that the dominant term describing the asymptotic behavior of $P\{u(x^n) \in du\}$, as $n \rightarrow \infty$, is the exponential function $e^{ns(u)}$.

In the theory of large deviations, the exponential asymptotic property of thermodynamic probability measures is referred to as a *large deviation principle* (LDP). Detailed explanations of this theory are available in a number of references including [24, 25, 26, 27, 28]. For our purpose it is important to note that an LDP can be formulated not only for the energy per particle, but also for more general macroscopic variables such as the total spin per particle or the vector of occupation numbers. As we point out in the next paragraph, the LDP is the basic tool for deriving the well known variational maximum-entropy principle characterizing the equilibrium macrostates of the associated macroscopic variables in the microcanonical ensemble.

The macroscopic variables that we consider are quantities $L(x^n)$ that take values in a space of macrostates \mathcal{L} , and that are assumed to satisfy an LDP expressed by

$$P\{L(x^n) \in dL\} = \int_{\{x^n : L(x^n) \in dL\}} P(dx^n) \quad (5)$$

The physical importance of equilibrium macrostates in statistical mechanics stems from the fact that any macrostate not in \mathcal{E}^u has an exponentially small probability of being observed given that the energy per particle of the system is fixed at the value u . This can be seen by introducing the microcanonical probability measure P^u , which is defined by conditioning the a priori measure P on the microcanonical conditioning set [23]; in symbols, $P^u(dx^n) = P\{dx^n | u(x^n) \in du\}$. Thus for $L \in \mathcal{L}$ we have

$$\begin{aligned} P^u\{L(x^n) \in dL\} &= P\{L(x^n) \in dL | u(x^n) \in du\} \\ &= \frac{P\{\{x^n : L(x^n) \in dL\} \cap \{x^n : u(x^n) \in du\}\}}{P\{u(x^n) \in du\}}. \end{aligned} \quad (8)$$

As shown in [23], the LDP's (4) and (5) and the asymptotic relationship (6) yield the LDP

$$P^u\{L(x^n) \in dL\} \asymp e^{-nI^u(L)}, \quad (9)$$

where

$$I^u(L) = \begin{cases} s(u) - s(L) & \text{if } u(L) = u \\ \infty & \text{otherwise.} \end{cases} \quad (10)$$

$$\asymp e^{ns(L)}.$$

In this formula L is an element of \mathcal{L} , and $P\{L(x^n) \in dL\}$ represents the probability measure of all microstates x^n such that $L(x^n)$ lies in the infinitesimal ball dL with center L . In addition, we assume that the energy per particle $u(x^n)$ can be written asymptotically as a function of $L(x^n)$ in the following sense. There exists a bounded, continuous function $u(L)$, called the *energy representation function*, with the property that $u(x^n) = u(L(x^n))$ for all microstates x^n or, more generally, that

$$|u(L(x^n)) - u(x^n)| \rightarrow 0 \quad (6)$$

uniformly over all microstates x^n as $n \rightarrow \infty$. In Section IV we give a concrete example, for the BEG model, of a macroscopic variable $L(x^n)$ for which assumptions (5) and (6) are valid. In the case of the BEG model, the energy representation will in fact be seen to be exact in the sense that $u(x^n) = u(L(x^n))$ for all x^n and n . However, the properties discussed in this paper are valid also for models for which the limit (6) holds.

Under assumptions (5) and (6) on the macroscopic variable $L(x^n)$, it is easily seen that the most probable macrostates L for configurations lying in the microcanonical conditioning set $\{x^n : u(x^n) \in du\}$ are those that maximize the entropy function $s(L)$ subject to the energy constraint $u(L) = u$ [23]. These constrained maximizers of $s(L)$ compose the set \mathcal{E}^u of *microcanonical equilibrium macrostates* associated with a given mean-energy value u ; in symbols,

$$\mathcal{E}^u = \{L : L \text{ maximizes } s(L) \text{ with } u(L) = u\}. \quad (7)$$

By the general theory of large deviations, $I^u(L)$ is non-negative for any macrostate $L \in \mathcal{L}$. Hence if $I^u(L) > 0$, then the microcanonical probability that $L(x^n)$ is near L goes to 0, as $n \rightarrow \infty$, at the exponential decay rate $I^u(L)$. This observation motivates our definition of the set \mathcal{E}^u of microcanonical equilibrium macrostates to be the set of macrostates L for which $I^u(L)$ attains its minimum of 0, and thus for which $s(L)$ is maximized under

the constraint $u(L) = u$. As a consequence of this definition, we have the following variational formula for the microcanonical entropy:

$$s(u) = \sup_{\{L:u(L)=u\}} s(L) = s(L^u), \quad (11)$$

where L^u is any macrostate contained in \mathcal{E}^u . [33]

B. Canonical ensemble

While the microcanonical ensemble is defined in terms of a fixed value of the mean energy u , the canonical ensemble is defined in terms of a fixed value of the inverse temperature β . In the canonical ensemble the relevant probability measure on \mathcal{X}^n is the Gibbs measure

$$P_\beta(dx^n) = \frac{1}{Z_n(\beta)} e^{-\beta n u(x^n)} P(dx^n). \quad (12)$$

As we did in the case of the microcanonical ensemble, we now state the LDP for $L(x^n)$ with respect to the canonical ensemble and then use this principle to define the set of canonical equilibrium macrostates. For any macrostate L and any microstates x^n satisfying $L(x^n) \in dL$, the continuity of the energy representation function implies that $u(x^n)$ is close to $u(L)$. Hence, we expect that

$$\begin{aligned} P_\beta\{L(x^n) \in dL\} &= \int_{\{x^n:L(x^n) \in dL\}} P_\beta(dx^n) \\ &= \frac{1}{Z_n(\beta)} \int_{\{x^n:L(x^n) \in dL\}} e^{-\beta n u(x^n)} P(dx^n) \end{aligned} \quad (15)$$

is close to

$$\frac{1}{Z_n(\beta)} e^{-\beta n u(L)} \int_{\{x^n:L(x^n) \in dL\}} P(dx^n) = \frac{1}{Z_n(\beta)} e^{-\beta n u(L)} P\{x^n : L(x^n) \in dL\}. \quad (16)$$

Substituting into this formula the large deviation estimate (5) for $P\{L(x^n) \in dL\}$ and the limit (14) relating $Z_n(\beta)$ and $\varphi(\beta)$, we obtain the LDP

$$P_\beta\{L(x^n) \in dL\} \asymp e^{-n I_\beta(L)}, \quad (17)$$

where

$$I_\beta(L) = \beta u(L) - s(L) - \varphi(\beta). \quad (18)$$

The function $I_\beta(L)$ is nonnegative for any macrostate L . As in the case of the microcanonical ensemble, for any macrostate L satisfying $I_\beta(L) > 0$ the LDP (17) shows that the corresponding canonical probability $P_\beta\{L(x^n) \in dL\}$ converges to 0 exponentially fast. As a result, such macrostates are not observed in the thermodynamic limit. The conclusion that we draw from the LDP (17) is that, with respect to the Gibbsian probability measure P_β , the most probable macrostates are those for which $I_\beta(L)$ attains its minimum of 0 or, equivalently, those for which the quantity $\beta u(L) - s(L)$ is minimized

In this formula, $Z_n(\beta)$ is the n -particle partition function defined by

$$Z_n(\beta) = \int_{\mathcal{X}^n} e^{-\beta n u(x^n)} P(dx^n), \quad (13)$$

in terms of which we define the *canonical free energy*

$$\varphi(\beta) = - \lim_{n \rightarrow \infty} \frac{1}{n} \ln Z_n(\beta). \quad (14)$$

This last quantity plays an analogous role in the asymptotic analysis of the canonical ensemble as the microcanonical entropy plays in the asymptotic analysis of the microcanonical ensemble.

for a fixed value of β . Accordingly, the set \mathcal{E}_β of *canonical equilibrium macrostates* associated with a fixed value of β is defined by

$$\mathcal{E}_\beta = \{L : \beta u(L) - s(L) \text{ is minimized}\}. \quad (19)$$

We end this section by motivating the existence of the limit (14) that defines the canonical free energy $\varphi(\beta)$. In the definition (13) of $Z_n(\beta)$, we first use (6) to replace the energy per particle $u(x^n)$ by $u(L(x^n))$. Since the macroscopic variable $L(x^n)$ takes values in \mathcal{L} , we can rewrite the resulting expression for $Z_n(\beta)$ not as an integral over the set \mathcal{X}^n of microstates, but as an integral over \mathcal{L} . Applying the LDP (5) for $P\{L(x^n) \in dL\}$, we obtain

$$\begin{aligned} Z_n(\beta) &= \int_{\mathcal{L}} e^{-\beta n u(L)} P\{L(x^n) \in dL\} \\ &\asymp \int_{\mathcal{L}} e^{-n[\beta u(L) - s(L)]} dL \end{aligned} \quad (20)$$

$$\asymp \exp\left(-n \inf_{L \in \mathcal{L}} \{\beta u(L) - s(L)\}\right).$$

The last step is a consequence of Laplace's method, which states that, as $n \rightarrow \infty$, the main exponential contribution to the integral comes from the largest value of the integrand or equivalently the smallest exponent [24, 26]. These calculations motivate the variational formula

$$\begin{aligned} \varphi(\beta) &= \inf_{L \in \mathcal{L}} \{\beta u(L) - s(L)\} \\ &= \beta u(L_\beta) - s(L_\beta), \end{aligned} \quad (21)$$

where L_β is any member of \mathcal{E}_β . We call this the *macrostate representation* of $\varphi(\beta)$. Formula (21) can be derived rigorously from the LDP (5) using Varadhan's theorem [23].

We denote by \mathcal{U} the set of mean energy values u for which $s(u) > -\infty$. Any $u \in \mathcal{U}$ is called an *admissible value* of the mean energy. Using a similar chain of arguments as in the preceding paragraph, one can also write

$$\begin{aligned} Z_n(\beta) &= \int_{\mathcal{U}} e^{-\beta n u} P\{u(x^n) \in du\} \\ &\asymp \int_{\mathcal{U}} e^{-n[\beta n - s(u)]} du \\ &\asymp \exp\left(-n \inf_{u \in \mathcal{U}} \{\beta u - s(u)\}\right). \end{aligned} \quad (22)$$

This asymptotic formula motivates the fundamental relationship

$$\varphi(\beta) = \inf_{u \in \mathcal{U}} \{\beta u - s(u)\}, \quad (23)$$

which expresses the thermodynamic free energy $\varphi(\beta)$ as the *Legendre-Fenchel transform* of the microcanonical entropy $s(u)$ [23]. We call this formula the *thermodynamic representation* of $\varphi(\beta)$.

The same formula for $\varphi(\beta)$ can also be derived by rewriting the infimum over L in (21) as an infimum over all mean-energy values u followed by a constrained infimum over all L satisfying $u(L) = u$. Using (11), we obtain

$$\begin{aligned} \varphi(\beta) &= \inf_{u \in \mathcal{U}} \inf_{\{L: u(L)=u\}} \{\beta u(L) - s(L)\} \\ &= \inf_{u \in \mathcal{U}} \left\{ \beta u - \inf_{\{L: u(L)=u\}} \{s(L)\} \right\} \\ &= \inf_{u \in \mathcal{U}} \{\beta u - s(u)\}. \end{aligned} \quad (24)$$

In the case where $s(u)$ is a strictly concave differentiable function of u and β is in the range of s' , this Legendre-Fenchel transform reduces to the usual differential form of the Legendre transform. This is given by

$$\varphi(\beta) = \beta u(\beta) - s(u(\beta)), \quad (25)$$

where $u(\beta)$ is the unique solution of the equation $\beta = s'(u)$.

III. EQUIVALENCE AND NONEQUIVALENCE OF ENSEMBLES

A. Thermodynamic level

The problem of the equivalence or nonequivalence of the microcanonical and canonical ensembles at the thermodynamic level is fundamentally related to the properties of the Legendre-Fenchel transform, and especially its invertibility or noninvertibility properties as a functional transform. Equation (23) expresses $\varphi(\beta)$ as the Legendre-Fenchel transform of $s(u)$. A basic question is whether one can invert this Legendre-Fenchel transform by applying the same transform to $\varphi(\beta)$ so as to obtain

$$s(u) = \inf_{\beta} \{\beta u - \varphi(\beta)\}. \quad (26)$$

We claim that such an inversion of the Legendre-Fenchel transform is valid if and only if $s(u)$ is concave on its domain of definition. The sufficiency can be seen by introducing the function

$$s^{**}(u) = \inf_{\beta} \{\beta u - \varphi(\beta)\}, \quad (27)$$

which is concave on its domain of definition and equals the minimal concave function majorizing $s(u)$ for all u [23, 24]. We call $s^{**}(u)$ the *concave hull* of $s(u)$ and depict it in Figure 1. It follows that if $s(u)$ is concave on its domain of definition, then s and s^{**} coincide. Replacing $s^{**}(u)$ by $s(u)$ in (27) expresses the microcanonical entropy $s(u)$ as the Legendre-Fenchel transform of the free energy $\varphi(\beta)$. Conversely, since any function written as the Legendre-Fenchel transform of another function is automatically concave on its domain of definition, it follows that if $s(u)$ is not concave on its domain of definition, then it cannot be written as the Legendre-Fenchel transform of the free energy.

This discussion motivates the following definitions. We define the two ensembles to be *thermodynamically equivalent* at u if $s(u) = s^{**}(u)$; in this case we say that s is concave at u . In the opposite case—namely, if $s(u) \neq s^{**}(u)$ —we call the two ensembles *thermodynamically nonequivalent* at u . [34] If $s(u) \neq s^{**}(u)$, then s is said to be nonconcave at u . For later reference we say that s is strictly concave at u if $s(u) = s^{**}(u)$ and s^{**} is strictly concave at u in the sense that the graph of s^{**} is not flat around u .

The noninvertibility of the Legendre-Fenchel transform for nonconcave functions has no effect on how φ is to be calculated from s . The canonical free energy $\varphi(\beta)$ is always a concave function of the inverse temperature β ; regardless of the form of $s(u)$, $\varphi(\beta)$ can always be expressed via the fundamental Legendre-Fenchel relationship (23). As we have just seen, however, $s(u)$ can be expressed as the Legendre-Fenchel of $\varphi(\beta)$ if and only if $s(u)$ is concave on its domain of definition. In this sense, the microcanonical ensemble is more fundamental than the canonical ensemble.

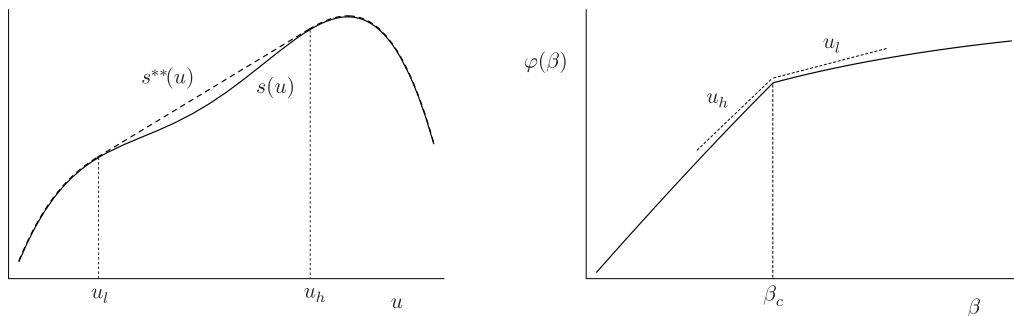


FIG. 1: (Left) Plot of a nonconcave entropy function $s(u)$ together with its concave envelope $s^{**}(u)$. The nonconcavity region equals the open interval (u_l, u_h) . (Right) The corresponding free energy function $\varphi(\beta)$ obtained by calculating the Legendre-Fenchel transform of $s(u)$. The region of nonconcavity of $s(u)$ is signalled by the appearance of a point β_c where $\varphi(\beta)$ is nondifferentiable. The quantity β_c equals the slope of the affine part of $s^{**}(u)$, and the left-hand and right-hand derivatives of φ at β_c equal u_h and u_l , respectively.

This apparent superiority of the microcanonical ensemble over the canonical ensemble does not prevent us from deriving a criterion based entirely on the canonical ensemble for verifying that the two ensembles are thermodynamically equivalent. Indeed, suppose that $\varphi(\beta)$ is differentiable for all β . Then the Gärtner-Ellis Theorem guarantees that, with respect to the a priori measure P , the energy per particle $u(x^n)$ satisfies the LDP with entropy function $s(u)$ given by the Legendre-Fenchel transform of $\varphi(\beta)$ [24, 25]. Because $\varphi(\beta)$ is assumed to be everywhere differentiable, the general theory of these transforms guarantees that $s(u)$ is strictly concave on its domain of definition. We conclude that if $\varphi(\beta)$ is everywhere differentiable, then thermodynamic equivalence of ensembles holds for all admissible values of the mean energy.[35] This can be expressed in more physical terms by saying that the absence of a first-order phase transition in the canonical ensemble implies that the ensembles are equivalent at the thermodynamic level. Unfortunately, the converse statement does not hold as the nondifferentiability of φ at some β corresponds to one of the following: either $s(u)$ is not concave over some range of mean-energy values or $s(u)$ is concave, but not strictly concave, over some range of mean-energy values (see Figure 1).

In the next subsection we examine the equivalence and nonequivalence of ensembles at a higher level; namely, that of equilibrium macrostates. Among other results we will see that the strict concavity of s at some mean-energy value u implies a strong form of equivalence that we call full equivalence.

B. Macrostate level

At the level of equilibrium macrostates, the natural questions to consider for characterizing the equivalence or nonequivalence of the microcanonical and canonical ensembles are the following. For every β and every L_β in the set \mathcal{E}_β of canonical equilibrium macrostates, does

there exist a value of u such that L_β lies in the set \mathcal{E}^u of microcanonical equilibrium macrostates? Conversely, for every u and every $L^u \in \mathcal{E}^u$, does there exist a value of β such that $L^u \in \mathcal{E}_\beta$? In trying to relate the macrostate level of equivalence and nonequivalence of ensembles with the thermodynamic level of equivalence and nonequivalence, we also ask the following. Are there thermodynamic conditions expressed in terms of properties of $s(u)$ or $\varphi(\beta)$ ensuring a correspondence or a lack of correspondence between the members of \mathcal{E}^u and those of \mathcal{E}_β ? In particular, does equivalence of ensembles at the thermodynamic level implies equivalence of ensembles at the level of equilibrium macrostates?

In [23] Ellis, Haven, and Turkington have provided precise answers to these questions expressed in terms of relationships between the solutions of the constrained minimization problem (7) that characterizes the set \mathcal{E}^u of microcanonical equilibrium macrostates and the solutions of the dual unconstrained minimization problem (19) that characterizes the set \mathcal{E}_β of canonical equilibrium macrostates. Their main results are summarized in items 1-4 below; they apply in general to any statistical mechanical model that have macroscopic variables $L(x^n)$ satisfying an LDP as in (5) with an entropy function $s(L)$ and that have an energy representation function $u(L)$ satisfying (6). In order to simplify the presentation, we consider only mean-energy values u lying in the interior of the domain of definition of s , and we assume that s is differentiable at all such u . The reader is referred to [23] for complete proofs of more general results that hold under weaker assumptions.

1. *Canonical equilibrium macrostates can always be realized microcanonically.* Let β be given. Then

$$\mathcal{E}_\beta = \bigcup_{u \in u(\mathcal{E}_\beta)} \mathcal{E}^u, \quad (28)$$

where $u(\mathcal{E}_\beta)$ denotes the set of mean-energy values u that can be written as $u(L)$ for some $L \in \mathcal{E}_\beta$. Formula (28) implies that for any macrostate $L_\beta \in \mathcal{E}_\beta$ there exists $u \in u(\mathcal{E}_\beta)$ such that $L_\beta \in \mathcal{E}^u$.

2. *Full equivalence.* If s is strictly concave at u , then $\mathcal{E}^u = \mathcal{E}_\beta$ for $\beta = s'(u)$.
3. *Partial equivalence.* If s is concave at u but not strictly concave, then $\mathcal{E}^u \subsetneq \mathcal{E}_\beta$ for $\beta = s'(u)$; i.e., \mathcal{E}^u is a proper subset of \mathcal{E}_β which does not coincide with \mathcal{E}_β . According to items 2 and 3, thermodynamic equivalence of ensembles for some value of u implies either full or partial equivalence of ensembles at the level of equilibrium macrostates for that u .
4. *Nonequivalence.* If s is nonconcave at u , then $\mathcal{E}^u \cap \mathcal{E}_\beta = \emptyset$ for all β . In other words, if there is thermodynamic nonequivalence of ensembles for some value of u , then the microcanonical equilibrium macrostates corresponding to that u are nowhere realized within the canonical ensemble.

We spend the remainder of this section sketching part of the proof of full equivalence as stated in item 2 and proving nonequivalence as stated in item 4. The crucial insight needed to prove that $s(u) = s^{**}(u)$ implies $\mathcal{E}^u = \mathcal{E}_\beta$ with $\beta = s'(u)$ is provided by a basic result of convex analysis which states that

$$s(v) \leq s(u) + \beta(v - u) \quad (29)$$

for all v if and only if $s(u) = s^{**}(u)$ and $\beta = s'(u)$. Accordingly, if we suppose that $s(u) = s^{**}(u)$, then

$$\beta u - s(u) \leq \beta v - s(v) \quad (30)$$

for all v . Using the thermodynamic representation of $\varphi(\beta)$ in (23) and the macrostate representation of $\varphi(\beta)$ in (21), we see that

$$\begin{aligned} \beta u - s(u) &\leq \inf_{v \in \mathcal{U}} \{\beta v - s(v)\} \\ &= \varphi(\beta) \\ &= \inf_{L \in \mathcal{L}} \{\beta u(L) - s(L)\}. \end{aligned} \quad (31)$$

Now choose any $L^u \in \mathcal{E}^u$; by the definition of this set $u(L^u) = u$, and $s(L^u) = s(u)$. This allows us to write

$$\beta u(L^u) - s(L^u) \leq \inf_{L \in \mathcal{L}} \{\beta u(L) - s(L)\}. \quad (32)$$

We deduce that L^u minimizes $\beta u(L) - s(L)$ or equivalently that $L^u \in \mathcal{E}_\beta$. Since L^u is an arbitrary element of \mathcal{E}^u , it follows that $\mathcal{E}^u \subseteq \mathcal{E}_\beta$. One completes the proof that $\mathcal{E}^u = \mathcal{E}_\beta$ by showing that if s is strictly concave at u , then \mathcal{E}^u is not a proper subset of \mathcal{E}_β .

To prove the assertion on nonequivalence of ensembles given in item 4, we proceed in a similar fashion. The assumption that s is nonconcave at u implies that

$$\begin{aligned} s(u) < s^{**}(u) &= \inf_{\gamma} \{\gamma u - \varphi(\gamma)\} \\ &\leq \beta u - \varphi(\beta) \end{aligned} \quad (33)$$

for all β . This can be rewritten as

$$\beta u - s(u) > \varphi(\beta). \quad (34)$$

Now choose any $L^u \in \mathcal{E}^u$ and any β . Since $u(L^u) = u$ and $s(L^u) = s(u)$, it follows that

$$\beta u(L^u) - s(L^u) > \varphi(\beta) = \inf_{L \in \mathcal{L}} \{\beta u(L) - s(L)\}. \quad (35)$$

This shows that L^u is not a minimizer of $\beta u(L) - s(L)$ and thus that $L^u \notin \mathcal{E}_\beta$. Since L^u is an arbitrary element of \mathcal{E}^u and β is arbitrary, we conclude that $\mathcal{E}^u \cap \mathcal{E}_\beta = \emptyset$ for all β , as claimed.

IV. ILLUSTRATION OF THE RESULTS FOR THE BLUME-EMERY-GRIFFITHS MODEL

We now come to the main point of our study which is to illustrate, in the context of a simple spin model, the general results presented in the previous section about macrostate equivalence and nonequivalence of ensembles and the relationship with thermodynamic equivalence and nonequivalence. The model that we consider for this purpose is a slight variant of the mean-field Blume-Emery-Griffiths (BEG) model defined by the following Hamiltonian [21, 29]:

$$U(x^n) = \sum_{i=1}^n x_i^2 - \frac{K}{n} \left(\sum_{i=1}^n x_i \right)^2. \quad (36)$$

In this formula x_i represents a spin variable at site i taking values in the set $\mathcal{X} = \{-1, 0, +1\}$, and K is a positive real constant. The a priori measure for the model is defined by

$$P(x^n) = \frac{1}{3^n} \text{ for every } x^n \in \mathcal{X}^n. \quad (37)$$

The macroscopic variable that we use to investigate the equivalence and nonequivalence of the microcanonical and canonical ensembles is the *empirical vector* $L(x^n)$ defined as

$$L(x^n) = (L_{-1}(x^n), L_0(x^n), L_{+1}(x^n)), \quad (38)$$

where for $j = -1, 0, +1$

$$L_j(x^n) = \frac{1}{n} \sum_{i=1}^n \delta_{x_i, j}. \quad (39)$$

In physics, this quantity is often referred to as the *one-particle state distribution* or *statistical distribution*; its three components $L_{-1}(x^n)$, $L_0(x^n)$, and $L_{+1}(x^n)$ give the proportion of spins in the microstate x^n that take the respective values -1 , 0 , and $+1$. Accordingly, $L(x^n)$ takes values in the space \mathcal{L} of probability vectors $L = (L_{-1}, L_0, L_{+1})$ defined by the conditions $L_j \geq 0$ for $j = -1, 0, +1$, and $L_{-1} + L_0 + L_{+1} = 1$.

For $L \in \mathcal{L}$ we define the function

$$u(L) = L_{+1} + L_{-1} - K(L_{+1} - L_{-1})^2. \quad (40)$$

The choice of the empirical vector for studying equivalence and nonequivalence of ensembles is dictated by the fact that we can represent the energy per particle

$$u(x^n) = \frac{U(x^n)}{n} = \frac{1}{n} \sum_{i=1}^n x_i^2 - K \left(\frac{1}{n} \sum_{i=1}^n x_i \right)^2 \quad (41)$$

in the form

$$\begin{aligned} u(x^n) &= \sum_{j \in \mathcal{X}} L_j(x^n) j^2 - K \left(\sum_{j \in \mathcal{X}} L_j(x^n) j \right)^2 \\ &= L_{+1}(x^n) + L_{-1}(x^n) - K(L_{+1}(x^n) - L_{-1}(x^n))^2 \\ &= u(L(x^n)). \end{aligned} \quad (42)$$

This display, which holds exactly for all x^n and n , verifies assumption (6) for the BEG model with the energy representation function $u(L)$ defined in (40). Moreover, one shows either by a combinatorial argument based on Stirling's approximation [1, 2, 24] or from Sanov's Theorem [24, 25] that with respect to the a priori probability measure P , $L(x^n)$ satisfies an LDP of the form (5) with the entropy function

$$s(L) = - \sum_{x \in \mathcal{X}} L_x \ln L_x - \ln 3. \quad (43)$$

These properties of $L(x^n)$ allow us to characterize the equilibrium macrostates with respect to each ensemble as solutions of an appropriate optimization problem. In order to simplify the notation, the components of probability vectors $L \in \mathcal{L}$ will be written as (L_-, L_0, L_+) instead of as (L_{-1}, L_0, L_{+1}) . We first consider the set \mathcal{E}^u of microcanonical equilibrium empirical vectors L^u associated with the mean energy u . According to the definition (7), the equilibrium macrostates L^u are characterized as maximizers of $s(L)$ over the space \mathcal{L} subject to the constraint $u(L^u) = u$. Solving this problem necessitates only the maximization of a function of one variable since the normalization constraint on the components of the empirical vector reduces the number of independent components of L^u to two, while the microcanonical energy constraint reduces this number by one more. The set \mathcal{E}_β of canonical equilibrium empirical vectors parameterized by the inverse temperature β is defined in (19). The elements L_β of this set are characterized as maximizers of the quantity $\beta u(L) - s(L)$ over \mathcal{L} . In this case, we are faced with an unconstrained two-dimensional maximization problem involving the two components L_+ and L_- . [36]

In Figure 2 we present a first set of solutions for \mathcal{E}^u and \mathcal{E}_β corresponding to the value $K = 1.1111$, together with a plot of the derivative of the microcanonical entropy function $s(u)$. Because neither of the two optimization

problems involved in the definitions of \mathcal{E}^u and \mathcal{E}_β could be solved analytically, we provide numerical results obtained using various routines available in the scientific software Mathematica. The top left plot of Figure 2 showing $s'(u)$ was obtained by calculating an empirical vector $L^u \in \mathcal{E}^u$, which satisfies $u(L^u) = u$ and $s(L^u) = s(u)$. The top right and the bottom left plots display, respectively, the canonical and microcanonical equilibrium components of the empirical vector as a function of the parameters β and u defining each of the two ensembles. In the top right plot, the solid curve can be taken to represent the spin +1 component of the equilibrium empirical vector L_β , while the dashed curve can be taken to represent the spin -1 component of the same equilibrium empirical vector. Since the BEG Hamiltonian satisfies the exchange symmetry $L_+ \leftrightarrow L_-$, the roles of the solid and dashed curves can also be reversed. For $\beta \leq \beta_c$, the solid curve represents the common value of $L_+ = L_-$. In all cases, the component L_0 of L_β is determined by $L_0 = 1 - L_+ - L_-$. The same explanation applies to the bottom left plot of L^u .

The first series of plots displayed in Figure 2 were designed to illustrate a case where $s(u)$ is concave and where, accordingly, we expect equivalence of ensembles. That the equivalence of ensembles holds in this case at the level of the empirical vector can be seen by noting that the solid and dashed curves representing the L_+ and L_- components of L_β in the top right plot can be put in one-to-one correspondence with the solid and dashed curves representing the same two components of L^u in the bottom left plot. The one-to-one correspondence is defined by the derivative of the microcanonical entropy $s(u)$: for a given u we have $L^u = L_{\beta(u)}$ with $\beta(u) = s'(u)$. Moreover, since the monotonic function $s'(u)$ can be inverted to yield a function $u(\beta)$ satisfying $s'(u(\beta)) = \beta$, we have $L_\beta = L^{u(\beta)}$ for all β . Thus, the equilibrium statistics of the BEG model in the microcanonical ensemble can be translated unambiguously into equivalent equilibrium statistics in the canonical ensemble and vice versa. In this case, the critical mean energy u_c at which the BEG model goes from a high-energy phase of zero magnetization $m(L) = L_+ - L_-$ to a low-energy phase of nonzero magnetization in the microcanonical ensemble can be calculated from the viewpoint of the canonical ensemble by finding the critical inverse temperature β_c that determines the onset of the same phase transition in the canonical ensemble. Since the two ensembles are equivalent, both the microcanonical and canonical phase transitions must be of the same order, which in this case is second-order.

In the second series of plots in Figure 3, a case of ensemble nonequivalence corresponding to the value $K = 1.0817$ is shown. Since in the top left plot $s'(u)$ is not monotonic, $s(u)$ is not concave. As in Figure 1, the open interval (u_l, u_h) of mean-energy values is the interval on which $s(u) \neq s^{**}(u)$; on this interval $s(u)$ is nonconcave and $s^{**}(u)$ is affine with slope β_c . By comparing the top right plot of L_β and the bottom left plot of L^u , we

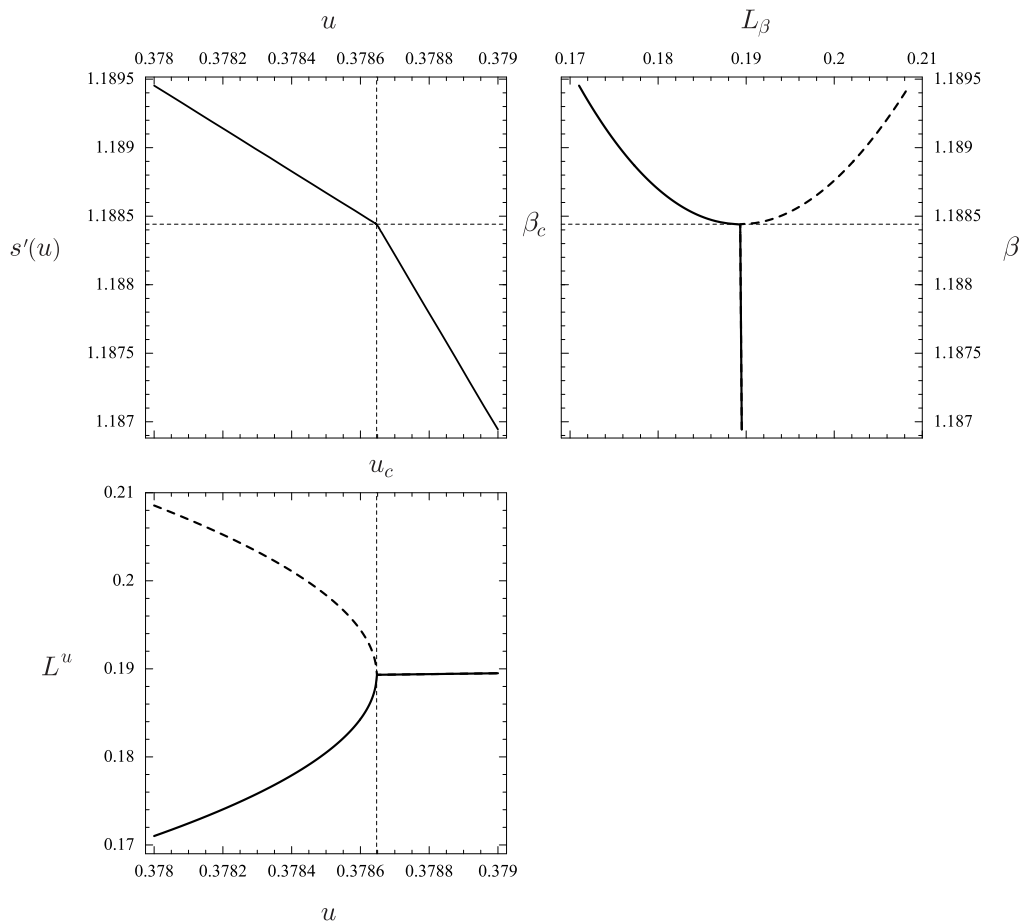


FIG. 2: Full equivalence of ensembles for the BEG model with $K = 1.1111$. (Top left) Derivative of the microcanonical entropy $s(u)$. (Top right) The components L_+ and L_- of the equilibrium empirical measure L_β in the canonical ensemble as functions of β . For $\beta > \beta_c$ the solid and dashed curves can be taken to represent L_+ and L_- , respectively, and vice versa. (Bottom left) The components L_+ and L_- of the equilibrium empirical measure L^u in the microcanonical ensemble as functions of u . For $u < u_c$ the solid and dashed curves can be taken to represent L_+ and L_- , respectively, and vice versa.

see that the elements of \mathcal{E}^u cease to be related to elements of \mathcal{E}_β for all mean-energy values u in the interval (u_l, u_h) . In fact, for any u in this interval of thermodynamic nonequivalence of ensembles (shaded region) no L_β exists that can be put in correspondence with an equivalent equilibrium empirical vector contained in \mathcal{E}^u . This lack of correspondence agrees with the rigorous results of Ellis, Haven, and Turkington [23] reviewed in Section III. Thus, although the equilibrium macrostates L^u corresponding to $u \in (u_l, u_h)$ are characterized by a well defined value of the mean energy, it is impossible to assign a temperature to those macrostates from the viewpoint of the canonical ensemble. In other words, the canonical ensemble is blind to all mean-energy values u contained in the domain of nonconcavity of $s(u)$. By decreasing β continuously through the critical value β_c , the equilibrium value of the energy per particle associated with the empirical vectors in \mathcal{E}_β jumps discontinuously from u_l to u_h (canonical first-order phase transition). However, outside the range (u_l, u_h) we have equivalence of ensembles,

and a continuous variation of β induces a continuous variation of u .

We can go further in our analysis of the plots of Figure 3 by noting that the phase transition exhibited in the microcanonical ensemble is second-order (continuous) whereas it is first-order (discontinuous) in the canonical ensemble. This provides another clear evidence of the nonequivalence of the two ensembles. Again, because the canonical ensemble is blind to all mean-energy values located in the nonequivalence region, only a microcanonical analysis of the model can yield the critical mean energy u_c . As for the critical inverse temperature β_c , which signals the onset of the first-order transition in the canonical ensemble, its precise value can be found by calculating the slope of the affine part of $s^{**}(u)$ or, equivalently, by identifying the point of nondifferentiability of $\varphi(\beta)$ (see the caption of Figure 1).

A further characterization of β_c can also be given in terms of the three solutions of the equation $s'(u) = \beta_c$. In Figure 3, u_l is the smallest of these solutions and u_h

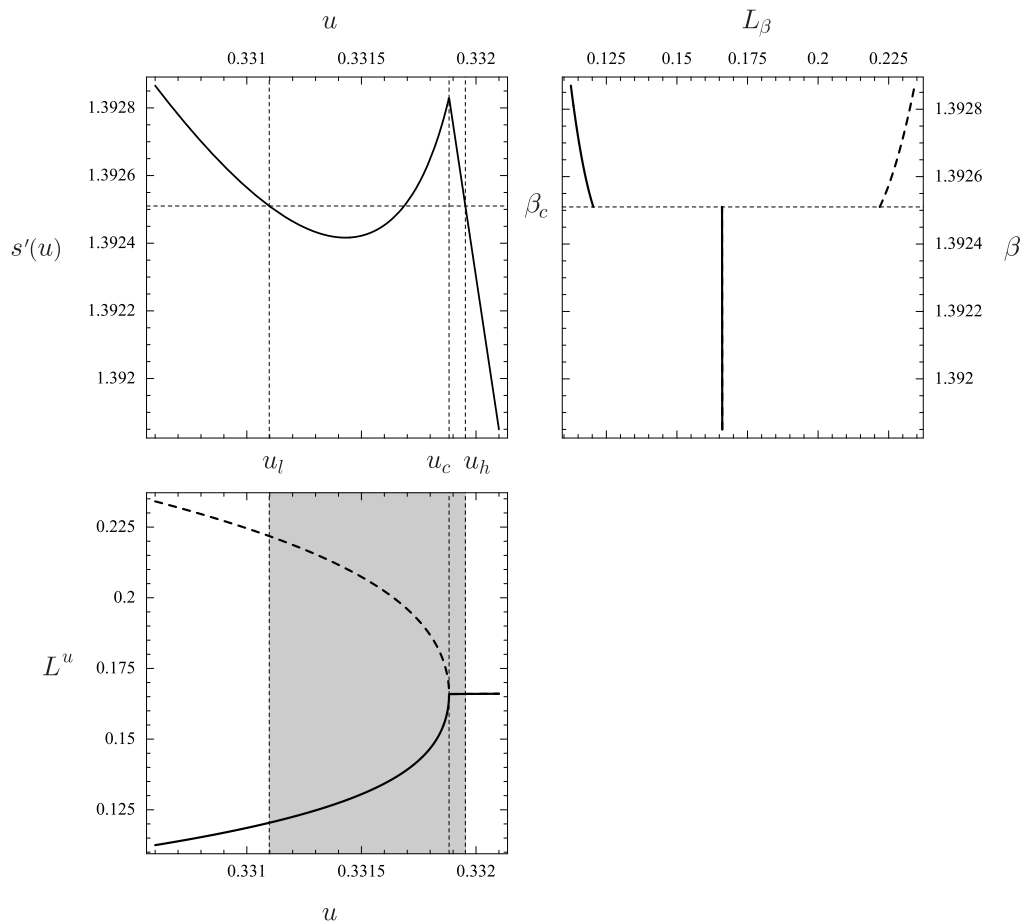


FIG. 3: Equivalence and nonequivalence regions for the BEG model with $K = 1.0817$. The solid and dashed curves are interpreted as in Figure 2. The shaded area in the bottom left plot corresponds to the region of nonequivalence of ensembles delimited by the open interval (u_l, u_h) . The ranges of the inverse temperature and the mean energy used to draw the plots were chosen so as to obtain a good view of the phase transitions.

the largest. We denote by u_m the intermediate solution of $s'(u) = \beta_c$. Because

$$s(u_l) = s^{**}(u_l), \quad s(u_h) = s^{**}(u_h), \quad (44)$$

and

$$s^{**}(u_h) - s^{**}(u_l) = \beta_c(u_h - u_l), \quad (45)$$

it follows that

$$\int_{u_l}^{u_h} [\beta_c - s'(u)] du = \beta_c(u_h - u_l) - [s(u_h) - s(u_l)] = 0. \quad (46)$$

Rewriting this integral in terms of u_m , we see that

$$\int_{u_l}^{u_m} [\beta_c - s'(u)] du = \int_{u_m}^{u_h} [s'(u) - \beta_c] du. \quad (47)$$

This equation expresses the equal-area property of β_c , first observed by Maxwell (see [6]).

To conclude this section, we present in Figure 4 a final series of plots of $s'(u)$, L_β , and L^u corresponding to

$K = 1.0805$, a slightly smaller value than the one considered in Figure 3. As in Figure 3, there also exists in Figure 4 an open interval (u_l, u_h) over which $s(u)$ is non-concave. For $u \in (u_l, u_h)$ we consequently have nonequivalence of ensembles, illustrated by the shaded region in the bottom left plot. As in Figure 3, the nonequivalence of ensembles is associated with a first-order phase transition in the canonical ensemble determined by β_c . The microcanonical phase transition in Figure 4 is also first-order due to the jump in $s'(u)$ as u increases through the critical value u_c . By contrast, the microcanonical transition is second-order in Figure 3.

V. CONCLUSION

We have provided in this paper a simple illustration of the fact that the nonequivalence of the microcanonical and canonical ensembles at the thermodynamic level entails a much more fundamental nonequivalence of these ensembles at the level of equilibrium macrostates. Fo-

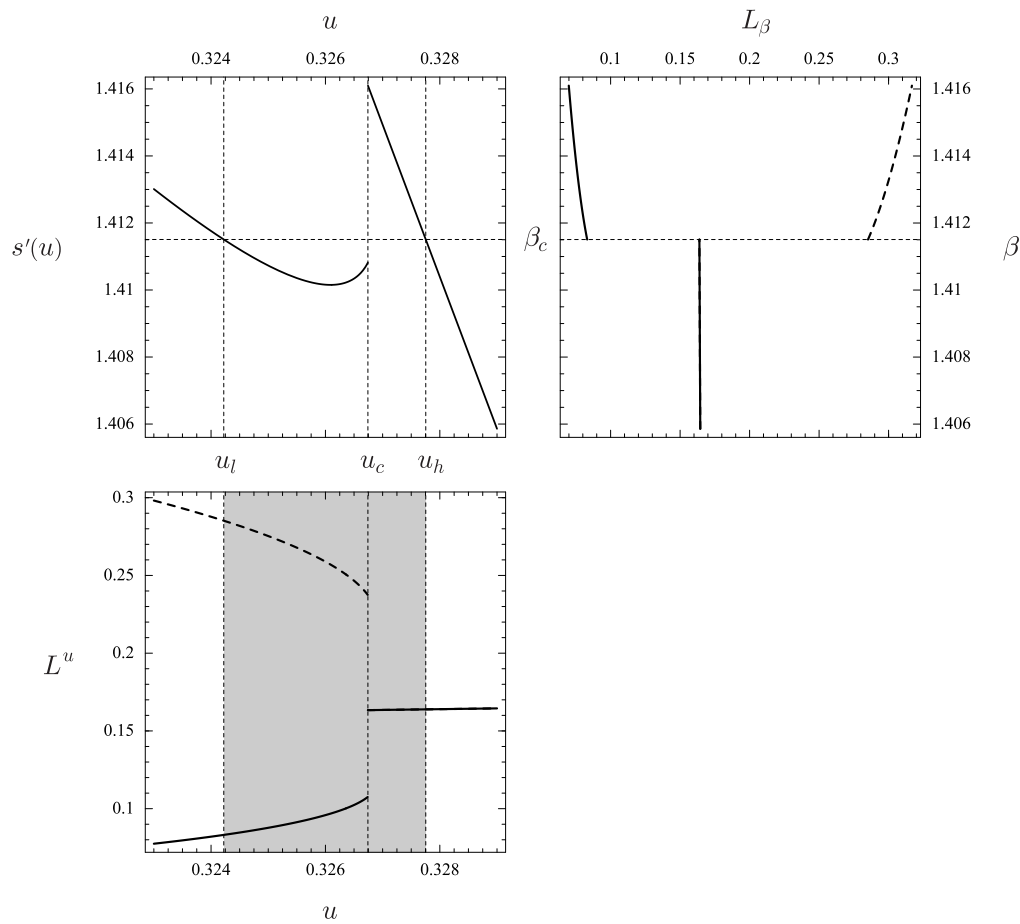


FIG. 4: Equivalence and nonequivalence regions for the BEG model with $K = 1.0805$. The solid and dashed curves are interpreted as in Figure 2. The shaded area in the bottom left plot corresponds to the region of nonequivalence of ensembles delimited by the open interval (u_l, u_h) .

cusing our attention on the mean-field Blume-Emery-Griffiths (BEG) model, we showed that if the microcanonical entropy s is strictly concave at a mean-energy value u , then the microcanonical equilibrium distributions of states characterizing the equilibrium configurations of the BEG model at the fixed value u are realized canonically for inverse temperature β given by $\beta = s'(u)$. We also showed that if the microcanonical entropy s is nonconcave at u , then the equilibrium distributions of states calculated microcanonically at the fixed value u are nowhere to be found in the canonical ensemble.

This latter case of macrostate nonequivalence is illustrated in Figures 3 and 4 for various values of the parameter K entering in the expression of the Hamiltonian of the BEG model. In each case there exists a continuum of *nonequivalent* equilibrium distributions of states that can be associated with mean-energy values u satisfying $s(u) < s^{**}(u)$, but cannot be associated with values of the inverse temperature β , the controllable parameter of the canonical ensemble. These results make it clear that the microcanonical ensemble is richer than the canonical ensemble, since the latter ensemble skips over the entire

range of mean-energy values for which the entropy is nonconcave. As we have seen, this implies the presence of a first-order phase transition in the canonical ensemble.

It should be remarked that although nonequivalent microcanonical equilibrium distributions cannot be associated with any value of the parameter β , one is not prevented from assigning to these distributions a *microcanonical* inverse temperature equal to the derivative of the microcanonical entropy. While this is a well-defined theoretical possibility, important questions are whether this procedure has any physical significance and whether such a microcanonical analog of inverse temperature can be measured.

We end this paper by discussing another point related to nonequivalent equilibrium macrostates corresponding to mean-energy values at which the microcanonical entropy is nonconcave. In the course of doing our numerical calculations leading to the determination of the sets \mathcal{E}^u and \mathcal{E}_β , we noticed that the equilibrium distributions of states L^u that exist microcanonically but not canonically are *metastable* macrostates of the canonical ensemble; i.e., they are local but not global minima of the quan-

tity $\beta u(L) - s(L)$. This was observed to occur when the mean energy u associated with these states satisfies $s''(u) < 0$. For u satisfying $s''(u) > 0$, we found instead that the macrostates L^u in the region of nonequivalence are saddle points of $\beta u(L) - s(L)$.

At this stage we cannot prove that this phenomenon holds in generality. However, our results lead us to conjecture that it is valid for a wide range of statistical mechanical models that have macroscopic variables satisfying an LDP as in (5) with an entropy function $s(L)$, and that have an energy representation function $u(L)$ satisfying (6). A number of theoretical and computational results found by Eyink and Spohn [30] and Antoni et al. [31], respectively, seem also to support such a conjecture. Besides, we know from the theory of Lagrange multipliers that all the microcanonical equilibrium macrostates are extremal points of the quantity $\beta u(L) - s(L)$, the global minimizers of which define the canonical equilibrium macrostates. In this context, what remains to be found is then just a way to determine the

precise nature of these extremal points based on the properties of $s(u)$. Work aimed at elucidating this problem is ongoing [32].

Acknowledgments

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- [33] Since $s(L)$ is nonpositive, $s(u)$ is also nonpositive. However, $s(u)$ can equal $-\infty$. Throughout this paper, we shall restrict ourselves to values of u for which $s(u) > -\infty$.
- [34] The occurrence of a negative heat capacity in the microcanonical ensemble provides a sufficient but not necessary criterion for characterizing the microcanonical and canonical ensembles as being thermodynamically nonequivalent. Indeed, the microcanonical heat capacity calculated at u can be positive even if $s(u) \neq s^{**}(u)$. On the other hand, a negative microcanonical heat capacity at u im-

plies necessarily that $s(u) \neq s^{**}(u)$.

- [35] The differentiability of $\varphi(\beta)$ for all β also implies that with respect to the canonical ensemble the energy per particle $u(x^n)$ converges to the constant value $\varphi'(\beta)$ in the limit $n \rightarrow \infty$ [24].
- [36] Another method for constructing \mathcal{E}_β can be based on the

determination of the canonical equilibrium value of the total spin per particle. The advantage of this alternate method is that the associated minimization problem is one-dimensional (R. S. Ellis, P. Otto, and H. Touchette, in preparation).