Simulation of phase transitions in highly asymmetric fluid mixtures

Jiwen Liu, Nigel B. Wilding, and Erik Luijten

1Dept. of Materials Science and Engineering, University of Illinois at Urbana-Champaign, Urbana, Illinois 61801, U.S.A.
2Department of Physics, University of Bath, Bath BA2 4LP, U.K.

(Dated: April 22, 2018)

We present a novel method for the accurate numerical determination of the phase behavior of fluid mixtures having large particle size asymmetries. By incorporating the recently developed geometric cluster algorithm within a restricted Gibbs ensemble, we are able to probe directly the density and concentration fluctuations that drive phase transitions, but that are inaccessible to conventional simulation algorithms. We develop a finite-size scaling theory that relates these density fluctuations to those of the grand-canonical ensemble, thereby enabling accurate location of critical points and coexistence curves of multicomponent fluids. Several illustrative examples are presented.

PACS numbers: 05.10.Ln, 64.70.Fx, 64.70.Ja, 82.70.Dd

The vast majority of commercially relevant fluids are multicomponent mixtures. An understanding of the phase behavior of these systems is of paramount importance for applications, and also a matter of great fundamental interest. With the advent of powerful computers, various computational techniques have been devised to directly determine fluid phase behavior. However, these methods are all restricted to fluids in which the various components have similar sizes, whereas important phenomena occur in highly size-asymmetric multicomponent fluids such as colloidal dispersions, colloidal nanoparticle mixtures, and polymer solutions.

The computational bottleneck for existing simulation methods arises from the fact that they cannot simultaneously relax a fluid system on disparate length scales. Specifically, for large size ratios, the big particles become ‘jammed’ by the smaller ones. Recently, however, building on earlier work, a geometric cluster Monte Carlo algorithm (GCA) was proposed that facilitates rejection-free simulations of highly size-asymmetric mixtures via large-scale collective updates which move whole groups of particles in a single step. Although this method has been successfully applied to problems relating to colloidal stabilization, in which size asymmetries span several orders of magnitude, it is incapable of dealing with the density and concentration fluctuations associated with phase separation, since it inherently operates in a canonical ensemble in which the particle number and volume are fixed. This limitation cannot be overcome by incorporating cluster moves into a standard grand-canonical (GC) or constant-\(NpT\) ensemble, because this does not address the underlying shortcomings of these ensembles with respect to sampling of density fluctuations in asymmetric mixtures.

It is the purpose of this Letter to introduce a method that overcomes these problems. This is achieved by embedding cluster moves in a variant of the Gibbs ensemble, in such as way that they couple to the density fluctuations, resulting in efficient exploration of configuration space. To exploit this approach we present a finite-size scaling theory that permits the determination of the critical point and the phase boundary. As an illustration, we apply the method to study liquid-vapor coexistence in asymmetric binary mixtures, for which we show that the presence of even small quantities of small-particle additives can strongly affect the location of the critical point. Furthermore, depending on the nature of the interaction of the additive with the fluid particles, the critical temperature can either be enhanced or depressed.

To enable density fluctuations, we distribute a prescribed number of particles \(N_0\) over two boxes, and devise an operation that exchanges particles between these boxes to maintain chemical equilibrium. By adopting the symmetrical restricted Gibbs (RG) ensemble, in which the boxes have equal constant volumes \(V = L^d\), geometric cluster moves can be used for this exchange operation. The prescription for a cluster move closely follows the original algorithm, with the crucial difference that a geometric operation not only alters the position of a particle, but also transfers it from one box to the other. Specifically, a pivot is placed at a random position within the first box, as well as at the corresponding position within the second box. A particle \(i\) is picked at random from one of the boxes (denoted 1) and point-reflected with respect to the pivot from its original position \(r_i\) to \(r_i'\). However, instead of placing the particle at \(r_i'\), we place it at the corresponding point \(r_i''\) in the other box (denoted 2), subject to periodic boundary conditions. Thereafter, any particle \(j\) interacting with particle \(i\) around its original position in box 1 or its new position in box 2 will also be considered for point reflection around the pivot and subsequent transfer to the opposite box, with a probability \(p_{ij} = \max(1 - \exp(-\beta \Delta_{ij}), 0)\), where \(\Delta_{ij} = -V(|r_i - r_j|)\) if particles \(i\) and \(j\) originally reside in the same box and \(\Delta_{ij} = V(|r_i' - r_j|)\) if \(i\) and \(j\) originally reside in different boxes. \(V(r)\) denotes a general pair potential and \(\beta\) the inverse temperature \(1/k_B T\). Note that \(p_{ij}\) solely depends on the pair interaction between particles \(i\) and \(j\), rather than on the total energy change resulting from the displacement of particle \(j\).
cluster construction proceeds iteratively for all particles interacting with each particle \( j \). Upon completion of a cluster move, a new pivot is selected. The proof of detailed balance is analogous to that for the generalized GCA \([3, 4]\).

The exchange of particles between boxes leads to (complementary) density fluctuations around the average number density \( \rho_0 = N_0/(2V) \) in each box. The fluctuation spectrum of the number density in box 1, \( \rho_1 = N_1/V \), can be related to the grand-canonical probability distributions \( P \) of the number densities of both boxes via

\[
P^{RG}(\rho_1|\rho_0, V, T) \propto P(\rho_1|\mu, V, T)P(2\rho_0 - \rho_1|\mu, V, T),
\]

where we note that \( P^{RG}(\rho_1) \) is independent of the choice of chemical potential \( \mu \) on the right-hand side \([11]\).

Since the right-hand side of Eq. (1) is the product of a function and a shifted (by \( 2\rho_0 \)) and reflected form of this function, \( P^{RG}(\rho_1) \) is symmetric (even) with respect to its mean \( \bar{\rho}_1 = \rho_0 \). To facilitate comparison of the forms of \( P^{RG}(\rho_1) \) for various choices of \( \rho_0 \), it is expedient to consider distributions of zero mean, to which end we define \( x = \rho_1 - \rho_0 \) and write

\[
P^{RG}(x|\rho_0, T) \propto P(x + \rho_0|T)P(-x + \rho_0|T),
\]

where we have suppressed reference to the constant volume \( V \) and the (arbitrary) chemical potential \( \mu \).

The parameter space of the RG ensemble is spanned by \( \rho_0 \) and \( T \). In the vicinity of the critical point \((\rho_0^c, T^c)\) terminating a line of liquid-vapor coexistence of a pure fluid or fluid mixture, \( P^{RG} \) exhibits universal scaling behavior. Introducing reduced variables \( \tilde{\rho}_0 \equiv (\rho_0 - \rho_0^c)/\rho_0^c \) and \( t \equiv (T - T^c)/T^c \), we make the following ansatz for the finite-size scaling (FSS) properties of \( P^{RG}(x) \),

\[
P^{RG}_L(x|\rho_0, t) = a_0 L^{\beta/\nu} \tilde{P}^{RG}(a_0 L^{\beta/\nu} x, a_1 \rho_0 L^{\beta/\nu}, a_2 t L^{1/\nu}),
\]

where \( \tilde{P}^{RG} \) is a universal function which is symmetric in \( x \) for all values of \( t \) and \( \tilde{\rho}_0 \); \( \beta \) and \( \nu \) are critical exponents, and \( a_0, a_1, a_2 \) are nonuniversal metric factors. The arguments in \( t \) and \( \tilde{\rho}_0 \) control deviations from criticality. The temperature field has the form familiar from the FSS properties of magnets \([12]\) or fluids \([12]\), while that in \( \tilde{\rho}_0 \) is particular to the RG ensemble. As one can verify \([14]\) from an expansion of Eq. (2) with respect to \( \rho_0 \), together with the known \([13]\) symmetry properties of the derivatives of \( P(x) \), variations in the form of \( P^{RG}(x) \) are—to leading order—controlled by the value of \( \tilde{\rho}_0 \); all terms having odd powers of \( \tilde{\rho}_0 \) are antisymmetric in \( x \) and hence absent on symmetry grounds.

To characterize the form of \( \tilde{P}^{RG}(x) \) as a function of \( \tilde{\rho}_0 \) and \( t \), it is useful to consider the behavior of the dimensionless fourth-order cumulant ratio \( Q \equiv \langle x^4 \rangle/\langle x^2 \rangle \) \([12]\), whose scaling properties follow from Eq. (3) as

\[
Q_L(\tilde{\rho}_0, t) = \tilde{Q}(q_1 \tilde{\rho}_0 L^{\beta/\nu}, q_2 t L^{1/\nu}),
\]

with \( \tilde{Q} \) a universal function and \( \tilde{Q}(0, 0) \equiv Q^* \). The value of \( Q^* = 0.711901 \) is known \( a \) priori by virtue of the result that \( P^{RG}(a_0 L^{\beta/\nu} x, 0, 0) \propto \langle P^*(L^{\beta/\nu} x) \rangle^2 \) \([11]\), with \( P^*(L^{\beta/\nu} x) \) the universal critical Ising magnetization distribution \([13]\). Measurements of \( Q_L(\tilde{\rho}_0, t) \) for a range of global densities \( \rho_0 \) provide a useful route to locating criticality. Specifically, consider the locus of points in \( \tilde{\rho}_0 - t \) space for which \( Q_L(\tilde{\rho}_0, t) = Q^* \), which we term the "iso-\( Q^* \) curve." Expanding Eq. (4) with respect to \( t < 0 \) and \( \rho_0 \), and recalling that only terms involving even powers of \( \rho_0 \) are nonzero, one has

\[
Q_L(\rho_0, t) = Q^* \left[ 1 + q_1 \tilde{\rho}_0^2 L^{2\beta/\nu} + q_2 t L^{1/\nu} + O(\rho_0^4, t^2) \right],
\]

from which it follows that, sufficiently close to the critical point, the iso-\( Q^* \) curve is a parabola in \( \tilde{\rho}_0 - t \) space,

\[
\sigma_0^2 = -(q_2/q_1) L^{(1-2\beta)/\nu}. \tag{6}
\]

The maximum of this parabola (at \( \rho_0 = t = 0 \)) coincides with the critical point and hence, by fitting to a few estimated points on the iso-\( Q^* \) curve, one can readily determine the critical parameters \( \rho_0^c \) and \( T^c \).

Turning now to the task of obtaining subcritical coexistence properties within the RG framework, we consider the peak positions of \( P^{RG}(\rho_1|\rho_0, t) \) on some (subcritical) isotherm. One finds that when \( \rho_0 \) equals the coexistence density \( \rho_d \equiv (\rho_g + \rho_l)/2 \), with \( \rho_g \) and \( \rho_l \) the gas and liquid densities, the peak positions of \( P^{RG}(\rho_1) \) coincide with the coexisting densities, which can thus be simply read off from a measured histogram of its form \([16]\). An effective method for locating the diameter \( \rho_d \) exploits the fact that the even moments of \( P^{RG}(x|\rho_0, t) \) are maximized when \( \rho_0 = \rho_d \). From the absence of odd powers of \( \rho_0 \) in the expansion of Eq. (4), it can then be shown \([14]\) that the variance \( \sigma^2(\rho_0|t) \) of \( P^{RG}(x|\rho_0, t) \) varies to leading order quadratically in \((\rho_0 - \rho_d)\), i.e.,

\[
\sigma^2(\rho_0|t) = \sigma^2(\rho_d, t) - b(\rho_0 - \rho_d)^2, \tag{7}
\]

with \( b \) a positive constant. By fitting to estimates of \( \sigma^2(\rho_d|t) \), this result facilitates a determination of \( \rho_d \) and thence the coexisting densities.

To test the scaling theory, we first perform simulations using the GCA in the restricted Gibbs ensemble for a pure Lennard-Jones (LJ) fluid. We employ a potential cutoff \( 2.5\sigma \) and reduced system sizes \( L^* = L/\sigma = 10, 20, 30 \). At fixed global density \( \rho_0 \), histogram reweighting can be used to determine a point on the iso-\( Q^* \) curve, i.e., the temperature at which \( Q_L(\rho_0, T) \) takes the value \( Q^* \). Repeating for a range of \( \rho_0 \) values allows the entire iso-\( Q^* \) line to be mapped. As shown in Fig. 11 the iso-\( Q^* \) curves for the various \( L^* \) are indeed parabolas \([14]\) that coincide at their maximum. A careful analysis \([14]\) of the position of this maximum, \( T^c = k_b T^c/\varepsilon = 1.1878 \) \( (2) \), \( \beta^* = \rho_0^c \sigma^3 = 0.3204 \) \( (5) \), reveals excellent agreement with an existing GC estimate.
The inset confirms the finite-size scaling predicted in Eq. (10) with remarkable accuracy. Furthermore, the critical density distribution $\bar{P}_{RG}(x)$ is indeed in quantitative agreement (not shown) with the square of the critical Ising magnetization distribution, as predicted [11].

Having established the validity of our methodology, we now exploit the strengths of the GCA to address a typical problem that is intractable for conventional simulations in the GC, $NpT$, or Gibbs ensemble. We consider a binary, strongly size-asymmetric mixture of LJ particles of size $\sigma$ and small particles ("additives") of diameter $\sigma_s = \sigma/10$. Depending on their interaction with the large particles, the presence of additives will affect the phase behavior and shift the location of the liquid-vapor critical point compared to the pure fluid. The additives mutually interact via a weakened LJ potential,

$$V_{ss}(r) = 4 \left( \frac{\sigma}{10} \right)^{12} \left( \frac{\sigma_s}{r} \right)^{6} \quad (r < 2.5 \sigma_s),$$

whereas a large and a small particle interact as hard spheres at a separation $\sigma_{ls} = (\sigma + \sigma_s)/2$.

As before, we perform simulations for a range of $\rho_0$, at fixed additive volume fraction $\phi_s = \frac{\pi}{6} \sigma_s^3 \rho_s = 0.005$, corresponding to $N_s = 10000$. Because the small particles are so numerous, and disperse relatively homogeneously, the insertion probability of a fluid (large) particle in a standard GC approach would be prohibitively small. By contrast, the present scheme renders it feasible to equilibrate the system and sample the density fluctuations. Figure 2 (diamonds) shows that the iso-$Q^*$ curve (plotted as a function of the total reduced density $\bar{\rho}_0 \equiv \rho (\sigma^3 + \rho_s \sigma_s^3)$) has a parabolic shape, as for the pure fluid. However, despite the relatively small volume fraction of additives, the maximum of this curve, i.e., the liquid-vapor critical point of the mixture, is markedly shifted. The increase in the critical temperature reflects an enhanced attraction between the large particles which stems from the entropic depletion interactions induced by the additives.

Our approach not only yields accurate estimates of critical points, but also entire coexistence curves. As described above, for each subcritical temperature, the variance of $P_{RG}(x)$ has a maximum at the coexistence curve diameter $\rho_d$ [Eq. (7)], as is confirmed in Fig. 4. The (total) densities of the coexisting liquid and vapor phases are determined from the peak positions of $P_{RG}$ [see Fig. 4(a) for an example], resulting in the phase diagram in Fig. 4(b). We emphasize that obtaining such a phase diagram in a reasonable timescale would not be feasible using even the most efficient traditional approach to fluid phase equilibria, namely GC simulation [17]. Our tests show that the GC relaxation time is too large to be reliably estimated. Nevertheless, a lower bound on

$T^c = 1.1876$ (3), $\rho^c = 0.3197$ (4) [17]. The inset confirms the finite-size scaling predicted in Eq. (10) with remarkable accuracy. Furthermore, the critical density distribution $\bar{P}_{RG}(x)$ is indeed in quantitative agreement (not shown) with the square of the critical Ising magnetization distribution, as predicted [11].

To highlight the subtle role of the nature of the interactions of the additives with the larger species, as well as their concentration, we study several systems in which there is a weak attraction between large and small particles. The interaction is again of the LJ form, Eq. (8), in which $\sigma_s$ is replaced with $\sigma_{ls}$. Already at $\phi_s = 10^{-3}$, the critical temperature is now noticeably depressed, as is evident from the shift of the iso-$Q^*$ maximum in Fig. 2 and at $\phi_s = 10^{-2}$, $T^c$ has decreased by almost 20%!

The $P_{RG}$ curve at $\phi_s = 0.005$ [Fig. 4(a)] shows a pronounced depression of $T^c$, as is confirmed in Fig. 4(b). The (total) densities of the coexisting liquid and vapor phases are determined from the peak positions of $P_{RG}$ [see Fig. 4(a) for an example], resulting in the phase diagram in Fig. 4(b). We emphasize that obtaining such a phase diagram in a reasonable timescale would not be feasible using even the most efficient traditional approach to fluid phase equilibria, namely GC simulation [17]. Our tests show that the GC relaxation time is too large to be reliably estimated. Nevertheless, a lower bound on...
the GC relaxation time, relative to that of the pure LJ fluid, can be estimated via a comparison of the large-particle transfer (insertion/deletion) acceptance probability $p_{\text{acc}}$. For liquid-like densities of the large particles ($\rho_1 \approx 0.6$), we find that for $\phi_s = 0.005$, $p_{\text{acc}} \sim 10^{-4}$; while for $\phi_s = 0.01$, this falls to $p_{\text{acc}} \sim 10^{-6}$. These values are to be compared with $p_{\text{acc}} \sim 10^{-1}$ for the pure LJ fluid. One can therefore expect the GC relaxation time of the mixtures we have studied to be several orders of magnitude greater than for the pure LJ fluid. Since the algorithm presented here operates along fundamentally distinct lines—proposing large-scale collective updates of clusters of small and large particles, and accepting them with unit probability—it is not hampered by this problem. Consequently it allows the efficient calculation of phase diagrams, even under conditions for which the GC approach fails.

Summarizing, we have extended the rejection-free GCA to the study of phase transitions, by embedding it within a restricted Gibbs ensemble. The accurate location of critical points and coexistence curves within this ensemble requires a suitable FSS theory, which has been presented as well. By means of illustration, we have applied our method to a strongly size-asymmetric LJ mixture, which cannot be studied with existing direct methods, e.g., the GC ensemble. We find that the liquid-vapor phase behavior is highly sensitive to the concentration of small particles and the nature of their interaction with the large ones. Thus our method should prove useful in predicting the alterations to phase behavior which occur when small particles of various types are added to a fluid. Furthermore, by employing the method with a variant of the GCA suitable for electrostatic interactions [19], it becomes possible to study the effects of adding salt on the phase behavior of charged colloids. Finally, we note that while for concreteness we have developed the formalism for the case of phase transitions whose order parameter is the density, the structure of our theory holds also for consolute points, or indeed situations where the order parameter is a linear combination of density and concentration.

This material is based upon work supported by the National Science Foundation under Grant DMR-0346914 (to EL) and by EPSRC Grant GR/S59208/01 (to NBW).

[1] R. G. Larson, The Structure and Rheology of Complex Fluids (Oxford University Press, Oxford, U.K., 1999).
[2] D. Frenkel and B. Smit, Understanding Molecular Simulation, 2nd ed. (Academic, San Diego, 2002).
[3] V. Tohver et al., Proc. Natl. Acad. Sci. U.S.A. 98, 8950 (2001).
[4] C. Dress and W. Krauth, J. Phys. A 28, L597 (1995).
[5] J. Liu and E. Luijten, Phys. Rev. Lett. 92, 035504 (2004).
[6] J. Liu and E. Luijten, Phys. Rev. E 71, 066701 (2005).
[7] J. Liu and E. Luijten, Phys. Rev. Lett. 93, 247802 (2004).
[8] C. J. Martinez et al., Langmuir 21, 9978 (2005).
[9] A. Z. Panagiotopoulos, Mol. Phys. 61, 813 (1987).
[10] K. K. Mon and K. Binder, J. Chem. Phys. 96, 6989 (1992).
[11] A. D. Bruce, Phys. Rev. E 55, 2315 (1997).
[12] K. Binder, Z. Phys. B 43, 119 (1981).
[13] A. D. Bruce and N. B. Wilding, Phys. Rev. Lett. 68, 193 (1992).
[14] J. Liu, N. B. Wilding, and E. Luijten, to be published.
[15] M. M. Tsypin and H. W. J. Blöte, Phys. Rev. E 62, 73 (2000).
[16] This is confirmed by inserting $\rho_0 = (\rho_0 + \rho_1)/2$ into Eq. (1); clearly when $\rho_1 = \rho_0$, one has $2\rho_0 - \rho_0 = \rho_1$ and thus both $P(\rho_1|t)$ and $P(2\rho_0 - \rho_1|t)$ are maximized, implying that so too is $P^{RG}(\rho_1)$. The same holds for $\rho_1 = \rho_2$.
[17] N. B. Wilding, Phys. Rev. E 52, 602 (1995).
[18] S. Assakura and F. Oosawa, J. Chem. Phys. 22, 1255 (1954).
[19] S. A. Barr and E. Luijten, to be published.