Perspective: Excess-entropy scaling

Jeppe C. Dyre

Glass and Time, IMFUFA, Department of Science and Environment, Roskilde University, P.O. Box 260, DK-4000 Roskilde, Denmark

(Received 6 September 2018; accepted 31 October 2018; published online 5 December 2018)

This article gives an overview of excess-entropy scaling, the 1977 discovery by Rosenfeld that entropy determines properties of liquids like viscosity, diffusion constant, and heat conductivity. We give examples from computer simulations confirming this intriguing connection between dynamics and thermodynamics, counterexamples, and experimental validations. Recent uses in application-related contexts are reviewed, and theories proposed for the origin of excess-entropy scaling are briefly summarized. It is shown that if two thermodynamic state points of a liquid have the same microscopic dynamics, they must have the same excess entropy. In this case, the potential-energy function exhibits a symmetry termed hidden scale invariance, stating that the ordering of the potential energies of configurations is maintained if these are scaled uniformly to a different density. This property leads to the isomorph theory, which provides a general framework for excess-entropy scaling and illuminates, in particular, why this does not apply rigorously and universally. It remains an open question whether all aspects of excess-entropy scaling and related regularities reflect hidden scale invariance in one form or other. Published by AIP Publishing. https://doi.org/10.1063/1.5055064

I. INTRODUCTION

Entropy is one of the most fascinating concepts of the physical sciences. It is deeply connected to another fundamental concept—time—because the direction of time is that of increasing entropy. In fact, the second law of thermodynamics is the only fundamental law of physics that is not time reversible. In 1977, a brief paper by Rosenfeld appeared in Physical Review A entitled “Relation between the transport coefficients and the internal entropy of simple systems,” which proposed a relation between entropy and time with no obvious connection to irreversibility: a liquid’s equilibrium dynamical properties are controlled by its entropy. Despite the intriguing nature of this claim, what became known as excess-entropy scaling for many years attracted little interest from the scientific community. Thus until his death in 2002, Rosenfeld’s seminal paper had been cited less than twenty times, half of which were autocitations.

Rosenfeld reported computer simulation results for simple model liquids of point particles like the well-known Lennard-Jones (LJ) system. His argument for excess-entropy scaling was based on the quasiuniversality of simple liquids traditionally explained by reference to the hard-sphere system based van der Waals picture of liquids. A possible explanation of the initial lack of interest in excess-entropy scaling is that after about 1980, the consensus in the liquid-state community was that simple liquids are well understood in terms of the hard-sphere reference system. As an illustration, consider the random barrier model of a particle jumping on a lattice with identical energies and randomly varying nearest-neighbor jump probabilities. This model has a trivial thermodynamics, in fact zero specific heat, but a highly complex and spatially heterogeneous dynamics as reflected, e.g., in the particle mean-square displacement as a function of time.

Since the onset of the new millennium, there has been steadily growing interest in excess-entropy scaling, which has turned out to apply more generally than originally thought, e.g., also for mixtures, molecular liquids, confined systems, etc. While this has highlighted the importance of the excess entropy, at the same time it has been realized that excess-entropy scaling has exceptions and thus cannot be a general, rigorous consequence of statistical mechanics. In regard to the non-rigorous nature of excess-entropy scaling, Hoover already in 1986 described the situation as follows: “this scaling relationship is, like van der Waal’s equation of state, a semiquantitative model, rather than a theory.” Rosenfeld confirmed this characterization in 1999, adding “like any corresponding-states relationship that links non-scaling force laws, excess-entropy scaling can only be approximate.”

There is no universal link between thermodynamics and dynamics because thermodynamics reflects the equilibrium probabilities of states, whereas dynamics reflects the rate of transitions between the states. As an illustration, consider the random barrier model of a particle jumping on a lattice with identical energies and randomly varying nearest-neighbor jump probabilities. This model has a trivial thermodynamics, in fact zero specific heat, but a highly complex and spatially heterogeneous dynamics as reflected, e.g., in the particle mean-square displacement as a function of time.

This article provides a brief review of excess-entropy scaling aimed at non-experts. After summarizing some necessary preliminaries in Sec. II, Sec. III presents examples of excess-entropy scaling from computer simulations. Selected experimental data are shown in Sec. IV. Section V gives examples of how excess-entropy scaling has been applied recently, e.g., for deriving viscosity models for industrial purposes.

a)dyre@ruc.dk
Section VI summarizes theoretical explanations of excess-entropy scaling proposed over the years, beginning with Rosenfeld's own hard-sphere-model based justification. Section VII shows that if excess-entropy scaling is a consequence of the property that state points with the same excess entropy have the same microscopic dynamics, the system in question must conform to “hidden scale invariance.” This symmetry, which states that a uniform scaling of configurations maintains the ordering of their potential energies, compare Eq. (19) below, applies to a good approximation for many liquids and solids but is never exact for realistic systems. Hidden scale invariance implies that the thermodynamic phase diagram becomes effectively one-dimensional in regard to structure and dynamics. The “isomorph theory” for the dynamics of systems that obey hidden scale invariance to a good approximation is a semi-rigorous theoretical framework, which quantifies the consequences of hidden scale invariance. Isomorph theory does not apply universally, only for systems with strong virial potential-energy correlations. On the other hand, it covers not just bulk, single-component, simple liquids, but also mixtures, molecular systems, solids, confined systems, out-of-equilibrium situations, etc. In this way, the isomorph theory takes excess-entropy scaling beyond Rosenfeld’s original focus (Sec. VIII). Section IX gives a brief outlook.

II. PRELIMINARIES

Consider a system in thermodynamic equilibrium at temperature $T$ with (number) density $\rho \equiv N/V$, in which $N$ is the number of atoms or molecules (“particles”) and $V$ is the volume. The system may consist of identical particles or mixtures of two or more different types of particles. If $S(\rho, T)$ is the system’s entropy, the excess entropy $S_{\text{ex}}$ is defined by subtracting from $S$ the entropy of an ideal gas at the same temperature and density, $S_\text{id}$,

$$S_{\text{ex}}(\rho, T) \equiv S(\rho, T) - S_\text{id}(\rho, T).$$

(1)

According to statistical mechanics, entropy is the logarithm of the phase-space volume of all microscopic states consistent with the given macroscopic thermodynamic condition; thus entropy quantifies our ignorance about the system’s microscopic state.\textsuperscript{10,12} Because the molecules of an ideal gas are “all over the place” with equal probability, the ideal-gas state corresponds to maximum ignorance or, as often stated, is maximally disordered. Consequently, $S_{\text{ex}} \leq 0$ always applies. Note that $S_{\text{ex}}$ increases with temperature just like the full entropy $S$ does; in fact $S_{\text{ex}} \to 0$ as temperature goes to infinity at fixed density because the system approaches to an ideal gas.

Not just the entropy but also the Helmholtz and Gibbs free energies may be written as sums of an ideal-gas term and an excess term.\textsuperscript{5,13} The system energy itself, $E$, is a sum of the kinetic energy and the potential energy $U$; the latter is zero for an ideal gas, so $U$ is the excess energy. The general relation $T = (\partial E/\partial S)_p$ has the following analog referring to the configurational degrees of freedom:\textsuperscript{12,14,15}

$$T = \left( \frac{\partial U}{\partial S_{\text{ex}}} \right)_p.$$  

(2)

Just as entropy and free energy are sums of an ideal-gas term and a term reflecting the loss of configurational degrees of freedom due to interactions, the same is the case for their derivatives. The pressure $p$, for instance, is given\textsuperscript{13} as

$$pV = Nk_BT + W,$$  

(3)

in which $W$ is the virial, an extensive quantity of dimension energy that is zero for an ideal gas, so $W/V$ is the excess pressure. The virial, which can be both positive and negative, is related to the potential energy\textsuperscript{5,13} by

$$W = \left( \frac{\partial U}{\partial \ln \rho} \right)_{S_{\text{ex}}}.$$  

(4)

Excess-entropy scaling uses the so-called macroscopically reduced units. In contrast to traditional unit systems, reduced units vary with the thermodynamic state point in question. The density defines the length unit $l_0$, the temperature defines the energy unit $e_0$, and the density and thermal velocity define the time unit $t_0$. If $m$ is the average particle mass, the length, energy, and time units are given\textsuperscript{1,8,14} by

$$l_0 = \rho^{-1/3}, \quad e_0 = k_BT, \quad t_0 = \rho^{-1/3} \sqrt{m/k_BT}.$$

(5)

It may seem impractical to employ a unit system that depends on the state point. On the other hand, these units do not require knowledge of the system’s Hamiltonian [the fact that the average particle mass $m$ appears in Eq. (5) is immaterial since this is just a constant].

Quantities made dimensionless by scaling with the above units are referred to as reduced, and this is henceforth indicated by a tilde. Consider, for instance, two quantities that Rosenfeld discussed in his original paper:\textsuperscript{1} the diffusion constant $D$ and the viscosity $\eta$.\textsuperscript{16} Since $D$ has dimension length squared over time, which is formally written as $[D] = l^2/t$, one defines $\tilde{D} \equiv D/(l_0^2/t_0)$, i.e.,

$$\tilde{D} \equiv \rho^{1/3} \sqrt{m/k_BT} D.$$  

(6)

Viscosity is shear stress over shear rate. The former has dimension energy over volume, and the latter has dimension inverse time, so $[\eta] = e_0 t_0/m$ since $[e] = m^2/h^2$ (compare the kinetic energy expression $mv^2/2$). Thus

$$\tilde{\eta} \equiv (\rho^{-2/3} \sqrt{mk_BT}) \eta.$$  

(7)

This dimensionless viscosity is found already in Andrade’s theory of viscosity from the 1930s.\textsuperscript{17}

We can now define: A liquid obeys excess-entropy scaling if its reduced dynamic properties at different temperatures and pressures are determined exclusively by $S_{\text{ex}}$. In other words, the lines of constant $S_{\text{ex}}$ in the thermodynamic phase diagram are lines of invariant reduced dynamics. This is the ideal situation; in practice, a system may obey excess-entropy scaling for some variables and not for others. For instance, one could imagine that all standard reduced transport coefficients—each
FIG. 1. Figure from Rosenfeld’s seminal 1977 excess-entropy paper showing simulation data for the reduced diffusion constant (left) and the reduced viscosity (right) as a function of the negative excess entropy per particle. Data are shown for the standard Lennard-Jones (LJ) liquid and the purely repulsive “soft-sphere” inverse-power-law pair potential \( v(r) \propto r^{-12} \). Rosenfeld discovered that the two systems give almost the same results for the dependence on excess entropy and, moreover, that the reduced diffusion constant and viscosity—in the figure denoted by \( D^* \) and \( \eta^* \)—are both approximately exponential functions of the excess entropy [Eq. (8)]. Reproduced with permission from Rosenfeld, Phys. Rev. A 15, 2545–2549 (1977). Copyright 1977 American Physical Society.

III. SIMULATION

Much of the evidence for excess-entropy scaling comes from computer simulations. We show in this section examples of this and discuss, at the end, how lines of constant excess entropy in the thermodynamic phase diagram are identified.

Figure 2 shows data for liquids of point particles, with (a) giving the reduced diffusion constant of the iconic Kob-Andersen binary Lennard-Jones liquid at different temperatures and densities, a standard model in the study of viscous liquids and the glass transition.\(^{18,21}\) The data collapse nicely, demonstrating excess-entropy scaling; Rosenfeld scaling Eq. (8) does not apply, although it is a good approximation to the high-temperature data. Figure 2(b) shows data inverse-power-law pair-potential system. The reduced diffusion constant is shown on the left, the reduced viscosity on the right (denoted by \( D^* \) and \( \eta^* \), respectively). By the Stokes-Einstein relation these two quantities are inversely related. The systems follow very similar trends. The same was reported for the one-component plasma, the system of same-charge particles interacting via Coulomb forces in a neutralizing background of opposite charge, as well as for the hard-sphere system.\(^1\) The diffusion data fall roughly on a line, indicating an exponential dependence on the excess entropy per particle,

\[ D \propto e^{\alpha S_{\text{ex}}/(Nk_B)}. \]  

Here \( \alpha \approx 0.8 \) is a numerical constant. A similar expression applies for the viscosity data in Fig. 1. Such an exponential dependence on \( S_{\text{ex}} \) of a reduced transport coefficient is nowadays referred to as “Rosenfeld scaling.” As we shall see, many systems conform to excess-entropy scaling, but not all of them obey Rosenfeld scaling.
for two metals modeled by many-body potentials. Rosenfeld scaling applies here, but note that the dynamic range is much smaller than in Fig. 2(a). In Fig. 2(b), the excess entropy is approximated by the two-particle entropy $S_2$ (see below).\textsuperscript{22}

Figure 3 shows simulation data for molecular models with (a) giving results for united-atom models of four hydrocarbons and (b) showing data for the symmetric dumbbell model consisting of two LJ spheres connected by a rigid bond. In both cases, excess-entropy scaling works well. Figure 3(b) gives data for two different inverse relaxation times. The fast process is the reorientation dynamics, the slow one is the decay of periodic density fluctuations at the wave vector corresponding to the maximum of the static structure factor. Despite a difference of more than one decade, both relaxation processes conform to excess entropy scaling. Rosenfeld scaling applies for all state points studied in (a), whereas in panel (b)—like in Fig. 2(a)—Rosenfeld scaling only applies at high temperatures, corresponding to low magnitudes of $S_{ex}$. (c) shows data for the incoherent intermediate scattering function as a function of time for the asymmetric dumbbell model probed along a curve of constant $S_{ex}$. The figure demonstrates that for this model not only the average relaxation time, but the entire dynamic signal is invariant along configurational adiabats. In experiments, such an invariance is referred to as “isochronal superposition;”\textsuperscript{26} in practice, this is tested by checking whether, independent of temperature and pressure, the average relaxation time determines the entire relaxation curve\textsuperscript{27} (compare Fig. 9 below).

A molecular model with internal degrees of freedom is considered in Fig. 4 presenting data for the flexible LJ-chain model, which consists of LJ particles connected by fixed-length, freely rotating bonds. Figure 4(a) shows the reduced viscosity for chain lengths 2, 4, 8, and 16, while (b) shows the reduced thermal conductivity. Excess-entropy scaling applies in both cases. The viscosity is chain-length dependent, but the thermal conductivity is not. Rosenfeld scaling applies for the viscosity at low temperatures, indicated by dashed lines in (a). Figure 4(c) shows data for three different intermediate scattering functions of the LJ-chain model, demonstrating invariant dynamics along a line of constant excess entropy, i.e., isochronal superposition. For comparison, (d) shows the same quantities along an isotherm for less than half the density variation. In all cases, Fig. 4 shows results for chain lengths that are too short to exhibit

**FIG. 3.** Numerical data demonstrating excess-entropy scaling for molecular liquids. (a) Reduced diffusion constant at different densities of united-atom models of $n$-octane, 2,5-dimethylhexane, 2,2-dimethylhexane, and 3-methyl-3-ethylpentane, all conforming to Rosenfeld scaling Eq. (8). Reproduced with permission from Chopra et al., J. Phys. Chem. B \textbf{114}, 16487–16493 (2010). Copyright 2010 American Chemical Society.\textsuperscript{23} (b) Inverse relaxation times of the coherent intermediate scattering function (“collective dynamic”) and of the individual molecular reorientation dynamics for the symmetric dumbbell model consisting of two Lennard-Jones spheres connected by a rigid bond.\textsuperscript{24} Although the two dynamics are quite different, they both conform to excess-entropy scaling. Reproduced with permission from Chopra et al., J. Chem. Phys. \textbf{133}, 104506 (2010). Copyright 2010 AIP Publishing LLC. (c) Incoherent intermediate scattering function for the asymmetric dumbbell model of two different LJ particles, evaluated along a curve of constant excess entropy. The data refer to highly viscous liquid state points; similar data along an isotherm with a smaller density variation exhibit more than three decade’s change of the average relaxation time.\textsuperscript{25} Reproduced with permission from Ingebrigtsen et al., J. Phys. Chem. B \textbf{116}, 1018–1034 (2012). Copyright 2012 American Chemical Society.
entanglement. For longer chains, the picture is more complex; here excess-entropy scaling applies only if a renormalization of the diffusion constant is allowed for.\(^{31}\) Figure 5 compares data for liquids confined to small volumes to those of the corresponding bulk liquids, with (a) giving data for the hard-sphere diffusion constant. The same
excess-entropy dependence applies for confined and bulk liquids. This behavior is also observed in (b), which shows the relaxation-time data for the asymmetric dumbbell model in both bulk and confinement. The fact that data referring to quite different boundary conditions collapse as a function of the excess entropy emphasizes the fundamental role of this thermodynamic variable.

Excess-entropy scaling does not always apply.\textsuperscript{34,36–39} Figure 6 reports data for (a) the Hertzian model, (b) a standard water model, (c) the Stillinger-Weber potential, and (d) the Fermi-Jagla model. Liquids with anomalies like water or silica that have, e.g., a diffusion constant which increases instead of decreases upon isothermal compression or which expand upon freezing, usually disobey excess-entropy scaling in the regions of the phase diagram where the anomalies appear.\textsuperscript{35,40–42}

How is the excess entropy determined in simulations? Chopra \textit{et al.}\textsuperscript{24} employed a two-step process to determine the Helmholtz free energy $F$ by first determining $F$'s density dependence from a grand-canonical transition-matrix Monte Carlo simulation at a high temperature, subsequently using a canonical-ensemble simulation to determine the free energy’s temperature variation at constant density. From $F$, the excess entropy is easily determined. Agarwal \textit{et al.} calculated $S_{ex}$ using thermodynamic integration with an ideal-gas reference state.\textsuperscript{19} Vasisht \textit{et al.} used thermodynamic integration by first finding the excess entropy at zero pressure at a reference temperature and from this calculating $S_{ex}$ for $p = 0$ at other temperatures, finally moving to non-zero pressure.\textsuperscript{36}

The excess entropy may be evaluated analytically if the equation of state is known.\textsuperscript{28} Rosenfeld, for instance, for the hard-sphere reference system used the Carnahan-Starling equation of state to calculate $S_{ex}$ as a function of the hard-sphere packing fraction.\textsuperscript{1}

A systematic expansion of $S_{ex}$ exists in terms of two-particle, three-particle contributions, etc.,\textsuperscript{46,47}

$$S_{ex} = S_2 + S_3 + S_4 + \ldots$$  \hspace{1cm} (9)

The two-particle contribution is calculated from the radial distribution function $g(r)$ as follows:\cite{46,47}

$$S_2/N = -2\pi \rho \int_0^\infty (g(r) \ln g(r) - g(r) + 1)r^2 dr. \quad (10)$$

In many cases, this gives the dominant contribution to $S_{ex}$.\cite{47} This is fortunate because $g(r)$ is a standard outcome of simulations. It is not generally the case, however, that $S_{ex} \approx S_2$.\cite{48–53} Some examples where this identity does not work well are given in Fig. 7.

If one is primarily interested in whether or not excess-entropy scaling applies, more than in how different reduced quantities depend on $S_{ex}$, the lines of constant $S_{ex}$ in the thermodynamic phase diagram are conveniently identified without knowledge of $S_{ex}$ by means of the following statistical-mechanical fluctuation identity:\cite{14}

$$\frac{\partial \ln T}{\partial \ln \rho}_{S_{ex}} = \frac{\langle \Delta U/\Delta W \rangle}{\langle (\Delta U)^2 \rangle}. \quad (11)$$

Here, the angular brackets denote $NVT$ canonical thermal averages at the state point in question and $\Delta$ refers to the given quantity minus its state-point average. If, for instance, the right-hand side is 3, a 1% increase of density is to be accompanied by a 3% increase of temperature in order to keep $S_{ex}$ constant. Recalculating the right-hand side at the new state point, etc., allows for step-by-step mapping out a line of constant excess entropy in the phase diagram.\cite{14}

IV. EXPERIMENT

Figure 8 shows experimental data for the viscosity and diffusion constant measured at various pressures and temperatures. The first three figures report the reduced viscosity of argon, carbon dioxide, and methane, as a function of the excess entropy (“reduced residual entropy”). These data are from experiments by Abramson measuring the viscosity over a large range of temperatures and pressures up to 6-8 GPa, using a rolling-ball technique. (d) and (e) give the diffusion constant of alkanes and a colloidal monolayer, respectively, as a function of the excess entropy.

Identifying $S_{ex}$ in experiments is not straightforward. We refer to the references of Fig. 8 for details of how this is done. Briefly, the determination of $S_{ex}$ at a given state point is often based on the use of an equation of state, e.g., expressed via the Helmholtz free energy as a function of temperature and density, fitted to the thermodynamic data.\cite{35,60,61} Alternatively, if one assumes the system in question obeys excess-entropy scaling for all dynamic processes, the dynamics measured on a certain time scale may be used to identify the lines of constant $S_{ex}$. This leads to the above-mentioned principle of isochronal superposition,\cite{26,27} according to which two state points with the same average reduced relaxation time have the same relaxation spectra. As shown in Sec. VII, if two state points have the same microscopic dynamics (except for a uniform scaling of space and time), they must have the same $S_{ex}$. Figures 3(c) and 4(c) show instances of isochronal superposition in computer simulations. Experimental examples are given in Fig. 9 in which (a) shows how dielectric spectra at different temperatures and pressures superpose if the loss-peak frequency is the same. An example of isochronal superposition over fourteen orders of magnitude was recently given by Hansen et al.\cite{59} For two van der Waals liquids, it was shown that state points with different temperatures and pressures, but the same relaxation time on the second time scale (probed by dielectric spectroscopy) have the same dynamics on the picosecond time scale (probed by neutron scattering). Isochronal data for one of these liquids are shown in Fig. 9(c), which may be compared to the isothermal data on the same liquid shown in Fig. 9(b). Interestingly, isochronal superposition does not apply for the hydrogen-bonded liquid studied with the same techniques.\cite{59}

V. APPLICATION

Of relevance to industrial applications, excess entropy has recently been identified as the crucial quantity for
rationalizing, e.g., how the bulk or confined-geometry viscosity and diffusion constant vary throughout the thermodynamic phase diagram. Excess-entropy-scaling based models have been applied to the transport properties of electrolytes and silica melts, methane and hydrogen absorption in metal-organic frameworks, the viscosity
and thermal conductivity of refrigerants and other liquids,72,73 the viscosity of the Earth’s iron-nickel liquid core at the relevant extreme pressures,74 separation of carbon isotopes in methane with nanoporous materials,75,76 etc. As stated by Novak, models based on the excess entropy provide a “practical approach to determining viscosity in process engineering, product engineering, oil and gas reservoir engineering, pipelines, and fracking applications.”77 A basic-research oriented application of excess-entropy scaling was recently reported in a paper on understanding water’s structure and dynamics in protein hydration layers.78

Figure 10 gives examples of how experimental data in application-relevant contexts are rationalized, either by the use of excess entropy itself or by models based on it, usually with one or more fitting parameters. Panel (a) shows data for methane diffusion in zeolites, and (b) shows data for bulk n-alkane viscosity. Panel (c) shows ketone data versus excess entropy calculated from the so-called perturbed-chain statistical associating fluid theory (PCP-SAFT),81 while (d) shows isopropyl-benzene’s reduced thermal conductivity versus excess entropy.

VI. THEORY

Since excess-entropy scaling is only approximate and does not apply universally (Fig. 6), it cannot be rigorously derived for systems in general. Arguments for excess-entropy scaling that have been given over the years are summarized below.

First, however, we note that there is one particular case for which excess-entropy scaling is exact. This happens if the potential-energy function is Euler homogeneous, i.e., obeys the following requirement, in which $\lambda$ is a uniform scaling parameter

$$U(\lambda R) = \lambda^{-\alpha} U(R).$$

In this case, the reduced-unit microscopic dynamics is exactly invariant along the lines in the thermodynamic phase diagram defined by $\rho^{2/3}/T = \text{Const.}$82,83 Physically, this reflects the
fact that a change of density can be compensated entirely by a change of temperature. Since the dynamics is invariant along certain lines except for a uniform scaling of space and time, this is the case also for the excess entropy that basically measures the available reduced volume in configuration space. Accordingly, the lines of constant excess entropy are lines of invariant reduced dynamics, which implies that, e.g., $\tilde{D}$ and $\tilde{\eta}$ are unique functions of the excess entropy whenever the potential-energy function is Euler homogeneous. An often studied example is that of particles interacting via purely repulsive inverse-power-law pair potentials.\textsuperscript{41,82,85–87}

**A. Rosenfeld-Hoover-Dzugutov hard-sphere argument**

The first justifications of excess-entropy scaling came from three authors giving closely related arguments based on the hard-sphere (HS) reference system. In Rosenfeld’s original publication,\textsuperscript{1} the reasoning was basically as follows. If a simple liquid is well represented by a HS system, its excess entropy is equal to that of the corresponding HS liquid, which is determined by the packing fraction $\phi$ (the occupied fraction of space). If the excess entropy per particle is denoted by $s_{\text{ex}} \equiv S_{\text{ex}}/N$, this implies that $s_{\text{ex}} = s_{\text{ex}}(\phi)$. Assuming the HS representation applies also for dynamic properties, one has $\tilde{D} = \tilde{D}(\phi)$, which by elimination implies $\tilde{D} = \tilde{D}(s_{\text{ex}})$. The new idea of Ref. 1 was that one can avoid the issue of determining the effective HS packing fraction of the system at the state point in question by referring instead to a thermodynamic property that is in a one-to-one relation with $\phi$, namely, the excess entropy.\textsuperscript{88}

At the time, Rosenfeld did not thoroughly justify the use of macroscopically reduced units [Eq. (5)], but in 1999 he stated that these are suggested “by elementary kinetic theory for a
dense medium of particles with thermal velocities but with a mean free path between collisions which is of the order of the average interparticle distance.”

An important finding in Rosenfeld’s 1977 paper was that the function $D(s_{ex})$ is quasiuniversal. In particular, virtually the same function applies for systems of purely repulsive pair forces and for the Lennard-Jones system, and this finding is highly nontrivial. From simulation data, Rosenfeld found that Eq. (8) applies; similarly he found for the viscosity that $\eta \propto \exp(-\beta s_{ex}/k_B)$, in which $\beta$ like $\alpha$ of Eq. (8) is a numerical constant of order unity (Fig. 1). This exponential Rosenfeld scaling, which is now known not to work in all cases of excess-entropy scaling (compare the above figures), was derived in 1986 by Hoover using an effective Einstein model for the HS system’s “cage rattling” particles,8 arguing as follows: If $\omega_0$ is the effective Einstein vibration frequency, i.e., the inverse of the time between collisions, any diffusion constant is estimated by $D \sim l_0^2 \omega_0$, in which $l_0 = \rho^{-1/3}$ is the average interparticle distance. This applies to the heat diffusion constant, the particle diffusion constant, and the transverse momentum diffusion constant, which is the so-called kinematic viscosity. According to the Navier-Stokes equation, the latter is $\eta/(m \rho)$, i.e., one has $\eta/(m \rho) = l_0^2 \omega_0$. In terms of the vibrational mean-square displacement ($x^2$), the excess entropy per particle is the logarithm of the available space relative to the interparticle distance, $s_{ex} = k_B \ln(\sqrt{x^2}/l_0)$.10,12 Writing the single-particle effective potential as $(1/2)m \omega_0^2 x^2$, equipartition implies $m \omega_0^2 = k_B T$. Eliminating ($x^2$) leads to $\exp(s_{ex}/k_B) = l_0^{-1} \omega_0^{-1} \sqrt{k_B T/m}$ or $\omega_0 = l_0^{-1} \sqrt{k_B T/m} \exp(-s_{ex}/k_B)$. Combining this with $\eta/(m \rho) = l_0^2 \omega_0$ leads to $\eta = m \rho l_0^2 \omega_0 = \rho^{3/2} \sqrt{mk_B T} \exp(-s_{ex}/k_B)$, which in reduced units [compare Eq. (7)] is simply $\tilde{\eta} = \exp(-s_{ex}/k_B)$.

Unaware of Rosenfeld’s little known paper, Dzugutov in 1996 rediscovered excess-entropy scaling.48 He did not use macroscopically reduced units, but instead scaled length by the effective hard-sphere radius and time by the inverse Einstein frequency. Nevertheless, Dzugutov’s physical ideas and arguments are quite similar to those of Rosenfeld and Hoover. Dzugutov justified the exponential excess-entropy dependence of the transport coefficients he found from simulations as follows: “The frequency of local structural relaxations, which defines the rate of the cage diffusion, is obviously proportional to the number of accessible configurations (per atom),” which by definition of the entropy is given by $\exp(s_{ex}/k_B)$. This argument, like the above, leads to Rosenfeld scaling Eq. (8) with $\alpha = 1$. Dzugutov argued, moreover, that the most important contribution to $s_{ex}$ is the two-particle entropy, making it possible to estimate diffusion constants from radial-distribution-function data.90

This extends excess-entropy scaling to the gas phase, although Rosenfeld scaling Eq. (8) does not apply in this part of phase space.

C. Kolmogorov-Sinai entropy connection

The dynamics of the hard-sphere system is determined by the rate at which trajectories in the multidimensional phase space diverge from one another in the course of time.92 This rate is quantified by the so-called Kolmogorov-Sinai (KS) entropy defined as the sum of all positive Lyapunov exponents.92–94 In 1998, an interesting connection to excess-entropy scaling was proposed by Dzugutov and co-workers in a paper demonstrating a linear relation between the KS entropy and the excess entropy for simple liquids.95 The existence of a universal relation between diffusion constant and KS entropy for simple liquids was confirmed in 2000 in simulations by Pang and co-workers,96 as well as in a theoretical study by Samanta et al. in 2004; the latter paper’s arguments apply also to mixtures.97,98

D. Mode-coupling theory

The main postulate of mode-coupling theory is that the structure determines the dynamics.99 In its simplest version, the pair distribution function $g(r)$ determines the incoherent intermediate scattering function. On the other hand, $g(r)$ also determines the two-particle contribution to the excess entropy $S_2$ [Eq. (10)], which as we have seen is often the dominant contribution to $S_{ex}$. Thus it makes good sense that mode-coupling theory predicts excess-entropy scaling in the $S_{ex} \approx S_2$ approximation.53,97,100,101

E. Generalized excess-entropy scaling

Truskett and co-workers in 2009 addressed the challenge that systems with soft pair potentials tend to disobey excess-entropy scaling. To remedy this, they proposed a modified scaling of the diffusion constant that by construction ensures excess-entropy scaling in the low-density (gas) limit. This is done by scaling $D$ with the product of $D_p$ (evaluated at a reference density) and $B(1 + d \ln B/d \ln T)$, in which $B$ is the second virial coefficient.102 In the dilute limit, rigorous theory implies that the diffusion constant scaled in this way varies as $1/s_{ex}$.102 For binary hard-sphere mixtures, Widom-Rowlinson mixtures, the Gaussian core model,103 and Hertzian particle fluids, this alternative form of excess-entropy scaling collapses diffusion-constant data much better than regular excess-entropy scaling, also in the non-dilute phase.104 The novel scaling is not consistent with standard Rosenfeld scaling, however, so for any given system these two ways of scaling cannot both collapse dynamic quantities as a function of the excess entropy.

F. Diffusion in a rugged potential-energy landscape

The potential-energy landscape of a liquid is very complex.105 Representing complexity by randomness,106 liquid dynamics has been modeled as jumps in a random landscape,107–109 for instance with energies distributed according to a Gaussian.110–113 For this model, Bagchi and co-workers in
2014 and 2015 showed via the effective-medium approximation that at moderate disorder, the diffusion constant is given by the Rosenfeld relation Eq. (8) with \( \alpha = 1 \).

Excess-entropy scaling applies also in the “extreme disorder limit”\(^{111} \) of the Gaussian random landscape model, the limit in which barriers are much larger than \( k_B T \) and where diffusion is, consequently, controlled by percolation.\(^{111} \) To see this, consider a system in \( d \) dimensions with a Gaussian density of states \( n(E) \propto \exp(-E^2/2\sigma^2) \). At temperature \( T \), the energy probability distribution is given by \( P(E) \propto n(E) \exp(-E/k_B T) \propto \exp(-(E - \bar{E})^2/2\sigma^2) \), in which \( \Delta E = E - \bar{E} \) is the average energy. If the percolation energy is denoted by \( E_c \), the effective barrier for escaping entirely from a state of low energy is given by \( \Delta E = E_c - E \).\(^{112} \) At low temperatures, the relevant energies move far into the negative tail of the Gaussian and one has \( \Delta E \equiv -E \). The diffusion constant is estimated by \( D \sim \langle l \rangle^2/\tau \), in which \( \langle l \rangle \) is the average distance between the most likely states [of energy close to \( \bar{E}(T) \)] and the relaxation time \( \tau \) is given by \( \tau = \tau_0 \exp(\Delta E/k_B T) \sim \tau_0 \exp(-E/k_B T) \), where \( \tau_0 \) is a prefactor. We estimate \( \langle l \rangle \) from \( \langle l \rangle^{-\alpha} \sim n(\bar{E}(T)) \) leading to \( \langle l \rangle \sim \exp(\bar{E}(T)^2/(2d\sigma^2)) \). After evaluating a Gaussian integral, one finds \( \langle l \rangle \sim \exp(3\sigma^2/(2k_BT)^2) \). In terms of \( \bar{E}(T) \), the excess entropy is given by \( S_{ex}/k_B \sim \ln(n(\bar{E}(T))) = -\bar{E}(T)^2/2\sigma^2 = -\sigma^2/(2k_BT)^2 \). Combining these equations leads to \( \bar{D} \sim \exp(\alpha S_{ex}/k_B) \), i.e., Rosenfeld scaling with \( \alpha = 32/d \).

### G. Other theoretical connections

This section discusses briefly some interesting studies in which entropy is related to dynamics in various ways without aiming at excess-entropy scaling.

#### 1. The Adam-Gibbs configurational entropy theory for the glass transition

In the study of glass-forming liquids, going back in time more than half a century there have been proposals of a link between the relaxation time and the “configurational entropy,” defined as the entropy relative to the perfect crystalline state, just as \( S_{ex} \) is defined relative to the gas state.\(^{116} - 118 \) In fact, ever since the studies of Simon\(^{119} \) and Kauzmann\(^{120} \) in 1931 and 1948, respectively, it has been noted that the rapid increase in structural relaxation time of liquids cooled toward their glass transition correlates with a significant drop in entropy. Martinez and Angell in 2001 reviewed the parallelism between the dramatic temperature dependence of the viscosity/relaxation time of supercooled liquids and the configurational entropy’s temperature dependence.\(^{121} \) Wolynes and co-workers have developed the so-called random first-order transition theory (RFOT) of the glass transition, in which the configurational entropy plays a crucial role for the dynamics.\(^{122} \) More recently, a generalized entropy theory of glass formation has been proposed,\(^{123} - 124 \) which combines the Adam-Gibbs model\(^{118} \) with the so-called lattice cluster theory.

#### 2. Relating to the geometry of the potential-energy-landscape

Chakraborty and Chakravarty have shown numerically for the Lennard-Jones system that the diffusion constant depends linearly on a number of properties characterizing the potential-energy landscape, e.g., the fraction of negative curvature directions, and mean, maximum, and minimum eigenvalues of the Hessian matrix.\(^{125} \) Assuming Rosenfeld scaling Eq. (8), this implies that the logarithm of the excess entropy exhibits the same linear dependencies, which provides interesting connections between the excess entropy and landscape properties.

### 3. Coarse-graining with the relative entropy

Scott Shell has shown that the optimal approximation to a given system by one from a specified class of simpler systems\(^{126} \) is obtained by minimizing the so-called relative entropy, an information-theoretic quantity also known as the Kullback-Leibler divergence, which measures the information lost upon coarse graining.\(^{127} \) Shell’s “relative-entropy coarse-graining method” was designed for reducing the number of degrees of freedom as in proper coarse-graining,\(^{128} \) but it has also been used to approximate the LJ system by a system with inverse power-law pair interactions.\(^{128} \) In the latter case, since the relative entropy is the difference between the two systems’ excess entropies, assuming Rosenfeld scaling Eq. (8) one concludes that the smaller the relative entropy is, the better does the coarse-graining approximation work for dynamic properties.

### 4. String-theory-based lower bound on viscosity

A surprising connection between entropy and viscosity appeared in a paper in 2005 by Kovtun et al., which showed that for a large class of strongly interacting quantum field theories whose dual description involves black holes in anti-de Sitter space, the ratio of viscosity over entropy density is \( \eta/4\pi k_B \).\(^{129} \) The paper conjectured that this number provides a lower bound to the entropy density \( \eta/(s\rho) \), where \( s \) is the entropy per particle (the excess entropy plus the ideal-gas entropy), i.e., that

\[
\frac{\eta}{s\rho} \geq \frac{\hbar}{4\pi k_B} \approx 6 \cdot 10^{-13} \text{Ks.} \tag{13}
\]

This lower-bound prediction, which obviously excludes superfluids for which \( \eta = 0 \), has been applied to strongly interacting systems like the quark-glue plasma created in a heavy ion collision\(^{129} \) and to the electron fluid of cuprate strange metals.\(^{130} \) Equation (13) is obeyed by helium and other rare gasses, as well as by water, nitrogen, ammonia, and molten alkali metals.\(^{130} - 133 \) For these systems, \( \eta/(s\rho) \) reaches its minimum close to the critical point, where it is 10-100 times larger than the string-theory based lower bound Eq. (13).\(^{133} \)

### 5. Rationalizing the behavior of anomalous systems

It is well known that water exhibits a multitude of anomalies, e.g., by having a diffusion constant that increases instead of decreases upon isothermal compression. Such anomalies are common for tetrahedrally coordinated systems and have also been reported, e.g., for the Gaussian-core model of spherically symmetric interactions based on a Gaussian function.\(^{103} \) Interestingly, the anomalies correlate with entropy anomalies.\(^{35,51,134,135} \) Thus regions of anomalous diffusivity and density behavior appear as “nested domes” in the
thermodynamic phase diagram within a structurally anomalous envelope. Despite the fact that water does not obey excess-entropy scaling, a given property exhibits anomalous behavior when the strength of the excess entropy anomaly exceeds a property-specific threshold.\textsuperscript{51}

6. Excess-entropy scaling in out-of-equilibrium situations

The excess entropy is an equilibrium thermodynamic quantity. In order to extend excess-entropy scaling to out-of-equilibrium situations, Krekelberg et al. proposed using the two-body excess entropy $S_2$ of Eq. (10).\textsuperscript{136,137} The point is that $S_2$ is given by the radial distribution function $g(r)$, which is also defined in, e.g., a steady-state flow. This idea was applied recently by Ingebrigtsen and Tanaka to collapse relaxation time data for different systems with shear rates extending far into the nonlinear regime, in a version in which the anisotropic structure under shear was taken into account in the definition of $S_2$.\textsuperscript{138}

VII. HIDDEN SCALE INVARIANCE

Rosenfeld’s 1977 paper\textsuperscript{1} reported that state points with different densities and temperatures but the same $S_{\text{ex}}$ have the same reduced diffusion constant, viscosity, and thermal conductivity. The simplest explanation is that at state points with the same excess entropy, the system’s atoms or molecules move about another in the same way to a good approximation, as suggested by the observation of the isochronal superposition in Figs. 3(c) and 4(c). This section investigates the possibility that the equation of motion is approximately invariant along the configurational adiabats.

If one imagines filming a liquid’s particles, the same movie cannot be seen at state points with different densities. The best one can hope for is that the particle motions are identical at two state points with the same $S_{\text{ex}}$ except for a uniform scaling of space and time. To eliminate the density dependence, coordinates must be scaled by the density as in the units defined in Eq. (5), compare the definition of the reduced coordinate vector $\tilde{\mathbf{R}} \equiv \mathbf{R}/\rho_0 = \rho^{1/3} \mathbf{R}$. Likewise, time must be scaled to ensure invariant thermal velocities along the lines of constant $S_{\text{ex}}$, and this is precisely what is obtained by using the reduced time coordinate $\tilde{t}$ defined via Eq. (5).

Assuming a system of identical particles, if $\tilde{\mathbf{F}} \equiv l_0 \mathbf{F}/e_0 = \rho^{-1/3} \mathbf{F}/k_B T$ is the reduced 3N-dimensional vector of all particle forces, it is straightforward to show that Newton’s equation of motion in reduced coordinates is

$$\frac{d^2 \tilde{\mathbf{R}}}{d\tilde{t}^2} = \tilde{\mathbf{F}}.$$ \hfill (14)

No mass appears on the left-hand side because the particle mass has been absorbed into the reduced time, compare Eq. (5). Equation (14) refers to the case of a single-component system, but the arguments given below apply also for mixtures. For a mixture, the left-hand side of Eq. (14) is for each particle $(m_i / m) d^2 \tilde{\mathbf{R}}_i / d\tilde{t}^2 = \tilde{\mathbf{F}}_i$.

Equation (14) applies for any system in equilibrium at a specified state point, whether or not the system obeys excess-entropy scaling. In general, the reduced force vector $\tilde{\mathbf{F}}$ depends on the coordinate vector $\mathbf{R}$, not just on $\mathbf{R}$, and the reduced dynamics at different state points are consequently not identical. If, however, the reduced force depends only on the reduced coordinates, i.e., if $\tilde{\mathbf{F}} = \tilde{\mathbf{F}}(\mathbf{R})$, there is invariance because in this case the equation of motion has no reference to the density.

How to check whether a given system obeys this requirement to a good approximation? To answer this question it is convenient to define the following microscopic excess-entropy function $S_{\text{ex}}(\mathbf{R})$:

$$S_{\text{ex}}(\mathbf{R}) \equiv S_{\text{ex}}(\rho, U(\mathbf{R})).$$ \hfill (15)

On the right-hand side, $S_{\text{ex}}(\rho, U)$ is the excess entropy of the thermodynamic state point with density $\rho$ and average potential energy $U$. In other words, Eq. (15) defines the excess entropy of the configuration $\mathbf{R}$ as the excess entropy of the thermodynamic state point with average potential energy equal to $U(\mathbf{R})$ and density corresponding to $\mathbf{R}$. The thermodynamic excess entropy of a given state point is then the average of its microscopic excess entropies, i.e., $S_{\text{ex}} = \langle S_{\text{ex}}(\mathbf{R}) \rangle$, just as the thermodynamic pressure is the average of the microscopic pressures.

The above is completely general and so is the inverse of Eq. (15),

$$U(\mathbf{R}) = U(\rho, S_{\text{ex}}(\mathbf{R})).$$ \hfill (16)

Recalling the expression for temperature Eq. (2), Eq. (16) implies for the force vector $\mathbf{F} = -\nabla U(\mathbf{R}) = -T \nabla S_{\text{ex}}(\mathbf{R})$. The reduced force vector is given by $\tilde{\mathbf{F}} = \rho^{-1/3} \mathbf{F}/k_B T = -\nabla S_{\text{ex}}(\mathbf{R})/k_B$, in which $\nabla \equiv \rho^{-1/3}$ is the gradient operator with respect to the reduced coordinates. If $\tilde{\mathbf{F}} = -\nabla S_{\text{ex}}(\mathbf{R})/k_B$ depends only on the reduced coordinate vector $\tilde{\mathbf{R}}$, by integration one concludes that $S_{\text{ex}}(\mathbf{R})$ is a function of $\tilde{\mathbf{R}}$ plus a function of the system’s overall density $\rho$, denoted by $\psi(\rho)$. Letting $T$ go to infinity at fixed $\rho$ takes the system to the gas limit in which $S_{\text{ex}} \to 0$. Thus $\psi(\rho)$ is independent of $\rho$ and may be put to zero without loss of generality. In summary, if $\lambda$ is a uniform scaling factor, the “same-movie requirement” $\tilde{\mathbf{F}} = \tilde{\mathbf{F}}(\mathbf{R})$ implies

$$S_{\text{ex}}(\mathbf{R}) = S_{\text{ex}}(\lambda \mathbf{R}).$$ \hfill (17)

In particular, it follows that the only possible lines of invariant microscopic dynamics in the phase diagram are the configurational adiabats, the lines of constant $S_{\text{ex}}$. This is not surprising, given the fact that the excess entropy is the logarithm of the reduced volume traced out in the course of time.\textsuperscript{10,12}

Since $S_{\text{ex}}$ for a system with hidden scale invariance depends only on a configuration’s reduced coordinates $\tilde{\mathbf{R}}$, Eq. (16) becomes

$$U(\mathbf{R}) = U(\rho, S_{\text{ex}}(\mathbf{R})).$$ \hfill (18)

How does one test whether Eq. (17) or, equivalently Eq. (18), applies to a good approximation for a given system? Consider two arbitrary configurations corresponding to the same density, $\mathbf{R}_a$ and $\mathbf{R}_b$, and suppose that $U(\mathbf{R}_a) < U(\mathbf{R}_b)$. Since $(\partial U / \partial S_{\text{ex}})_\rho > T > 0$, one concludes from Eq. (18) that $S_{\text{ex}}(\mathbf{R}_a) < S_{\text{ex}}(\mathbf{R}_b)$. Equation (18) then implies that a uniform scaling of the two configurations to a different density preserves the relation between their potential energies. That is, if one configuration has lower potential energy than
another, this is the case also after a uniform scaling. Formally, this property is expressed by the following logical implication:

$$U(R_a) < U(R_b) \Rightarrow U(\lambda R_a) < U(\lambda R_b).$$  \hspace{1cm} (19)

This is referred to as the hidden-scale-invariance condition. It can be shown that Eq. (18) is mathematically equivalent to Eq. (19). Equation (19) applies rigorously only for an Euler-homogeneous potential-energy function [Eq. (12)] plus a constant. Hidden scale invariance is easily checked in a simulation. Figure 11 shows in (a) data for the potential energies of scaled Lennard-Jones-Gaussian model configurations for which Eq. (19) is not a good approximation and, for comparison, in (b) data for the exponentially repulsive pair potential where Eq. (19) works much better.

Because of continuity, Eq. (19) always applies to a good approximation for small density changes, so how does one judge to which degree hidden scale invariance applies? A quantitative measure for this is provided by the constant-density virial potential-energy Pearson correlation coefficient $R$ defined as follows:

$$R \equiv \frac{\langle \Delta U \Delta W \rangle}{\sqrt{\langle \Delta U^2 \rangle \langle \Delta W^2 \rangle}}. \hspace{1cm} (20)$$

Here $W$ is the microscopic virial, the thermal average of which (also denoted by $W$) obeys Eq. (3). In general, $-1 \leq R \leq 1$. If Eq. (18) applies rigorously, one has $R = 1$ (in exotic cases, $R = -1$). This is because at any given density, knowledge of $U(R)$ via Eq. (18) implies knowledge of $S_{ex}(\tilde{R})$, which determines the microscopic virial via the analog of Eq. (18), $W(R) = W(\rho, S_{ex}(\tilde{R}))$ (this expression follows from the definition $W(R) \equiv (\partial U(R)/\partial \ln \rho)_R^{13}$. Thus if $R$ is close to unity, the hidden-scale-invariance condition Eq. (19) applies to a good approximation and Newton’s equation of motion [Eq. (14)] is approximately invariant along the configurational adiabats.

FIG. 11. Potential energies of configurations generated by uniform scaling of equilibrium configurations. At each density, the average potential energy was subtracted, after which the data were normalized to unit variance. Equation (19) applies to a good approximation if the curves rarely cross one another. (a) Data for the Lennard-Jones-Gaussian model with twenty configurations selected from an equilibrium simulation (at the density marked by the red dashed line) scaled uniformly $\pm 20\%$ in density. There are several crossings, showing that Eq. (19) does not apply very well. Reproduced with permission from Schrøder and Dyre, J. Chem. Phys. 141, 204502 (2014). Copyright 2014 AIP Publishing LLC. (b) Similar data for the EXP exponentially repulsive pair-potential system.

FIG. 12. State-point jumps monitored in NVT simulations. Isomorphs were traced out numerically in the phase diagram using Eq. (11) in a step-by-step fashion, involving a density change of 1%. (a) Relaxation of the potential energy toward its equilibrium value after three jumps at $t = 0$ in the thermodynamic phase diagram for the viscous Lewis-Wahnström ortho-terphenyl model. The black curve is for a jump between two state points on the same isomorph. Here there is, after an initial instantaneous jump in potential energy (not visible), no subsequent relaxation toward equilibrium. The red and green curves are for non-isomorph jumps. Reproduced with permission from Ingebrigtsen et al., J. Phys. Chem. B 116, 1018–1034 (2012). Copyright 2012 American Chemical Society. (b) Same after jumps from three different state points of a face-centered cubic crystal of gold modeled via a realistic effective medium non-pair potential. Only the jump from a state point on the final state point’s isomorph (red curve) leads to instantaneous equilibration. Reproduced with permission from Friedeheim et al., e-print arXiv:1810.07255 (2018). Copyright 2018 Author(s).
The criterion $R > 0.9$ has been used for defining the class of “strongly correlating systems,” which to avoid confusion with strongly correlated quantum systems have been termed “Roskilde (R)-simple” systems. These are the systems that obey hidden scale invariance to a good approximation.

All systems have lines in the thermodynamic phase diagram of constant excess entropy. For R-simple systems, these lines are termed isomorphs because the reduced-unit structure and dynamics are invariant to a good approximation along isomorphs.

VIII. BEYOND EXCESS-ENTROPY SCALING

The simplest explanation of excess-entropy scaling is that the microscopic dynamics at state points with the same excess entropy is virtually the same. If this is the case, the system in question has hidden scale invariance [Eq. (19)]. This property, however, only applies approximately and only for certain systems, and even for these it applies only in part of the thermodynamic phase diagram. This explains why excess-entropy scaling is neither exact nor universal. The class of R-simple systems—those with hidden scale invariance—is believed to include most metals and van der Waals bonded systems, but exclude most systems with strong directional bonds like hydrogen-bonded or covalently bonded systems. Ionic and dipolar systems constitute in-between cases that may or may not exhibit hidden scale invariance, depending on how strong the Coulomb forces are. Whenever hidden scale invariance applies, however, it has a number of interesting consequences beyond the classical excess-entropy scaling of transport coefficients. The present section gives examples of this.

While Rosenfeld originally proposed excess-entropy scaling for single-component liquids of particles interacting via pair forces, hidden scale invariance applies equally well to...
mixtures, non-pair-force systems, confined systems, molecular systems, and solids. For molecules, Eq. (19) relates to a uniform scaling of the center-of-mass coordinates, while orientations and molecular sizes are kept fixed. An example of isomorph-invariant dynamics was given in Fig. 4(c) reporting simulations of a liquid of ten-bead flexible Lennard-Jones chain molecules.

Isomorph invariance of the dynamics implies that a uniform scaling of a configuration selected from an equilibrated simulation at one state point results in an equilibrium configuration corresponding to a state point of different density and temperature (with the same excess entropy). Such a scaling is referred to as an “isomorph jump” because the system jumps instantaneously from equilibrium to equilibrium along an isomorph, in effect creating kind of a wormhole in the thermodynamic phase diagram. Two examples of this are provided in Fig. 12 showing in (a) an isomorph jump for a united-atom model of ortho-terphenyl and in (b) for a gold crystal. The latter figure illustrates that hidden scale invariance is not limited to the liquid phase; in fact, although this has not yet been studied very much, isomorph-theory predictions appear to apply equally well for solids as for liquids.

A recent application of the isomorph theory was to the thermodynamics of melting. For R-simple systems, the melting line is an approximate isomorph. Via first-order Taylor expansions, this fact makes it possible to predict from simulations carried out at a single coexistence state point the variation along the melting line of quantities like pressure, density, heat of melting, viscosity, Lindemann ratio, etc. Examples are given for the LJ system in Fig. 13, in which the red curves are theoretical predictions based on simulations carried out at $T = 2$ and the black dots are numerical data.

The isomorph theory has been applied to rationalize results from computer simulations of single-component and binary LJ-type systems, simple molecular models, crystals, nano-confined liquids [compare Fig. 5(b)], non-linear shear flows, zero-temperature plastic flows of glasses, polymer-like flexible molecules [compare Fig. 4(c)], metals studied by ab initio density-functional-theory computer simulations, plasmas, physical aging, and for justifying a quasiuniversal viscosity equation for supercritical R-simple liquids.

Experimental confirmations of isomorph-theory predictions were presented for van der Waals bonded liquids in Refs. 59 and 178–180. As an example, the isomorph theory predicts how much the dielectric loss is suppressed by the application of pressure (Fig. 14). For hydrogen-bonded liquids, the theory does not work, though; this confirms earlier findings from simulations of water and methanol.

On the theoretical side, the isomorph theory led to “NVU molecular dynamics” defined as motion at a constant velocity along geodesic curves on the high-dimensional constant-potential-energy hypersurface. Despite the fact that the potential and kinetic energies are both strictly conserved in NVU dynamics, this novel dynamics gives results that are for most quantities identical to those obtained by standard NVE or NVT Newtonian dynamics. The isomorph theory has also been applied to formulate a theoretical tool, the “isomorph filter,” according to which any universality valid theory for the relaxation time of glass-forming liquids must be based on isomorph-invariant quantities since the (reduced) relaxation time itself is an isomorph invariant. This makes it possible to rule out a number of well-known models of the glass transition.

According to the standard explanation of quasiuniversality, two systems have the same structure and dynamics to a good approximation if they correspond to hard-sphere systems with the same packing fraction. Rosenfeld realized that this translates into the two systems having the same excess entropy. The isomorph theory has recently made it possible to derive simple liquids’ quasiuniversality without reference to the hard-sphere system. The proof is based on using the exponentially repulsive EXP pair-potential system as a reference system. Different simple liquids with the same $S_{ex}$ have virtually the same structure and dynamics because they have virtually the same reduced-coordinate potential-energy function, and this function is conveniently identified with that of the exponentially repulsive pair-potential system. Systems that are quasiuniversal are those for which the pair potential may be approximated by a finite sum of exponentially repulsive pair potentials with prefactors that in reduced units are numerically much larger than unity. This provides a method for determining which pair-potential systems are quasiuniversal and which are not.

Maimbourgh and Kurchan recently showed that for systems with strong repulsive forces, the isomorph theory is exact in infinite dimensions. This suggests the intriguing scenario
that excess-entropy scaling is a finite-dimensional manifestation of a high-dimensional regularity and that, consequently, deviations from excess-entropy scaling should be quantified via an expansion in $1/d$ in which $d$ is the spatial dimension.\cite{171}

Readers that are further interested in the isomorph theory are referred to recent reviews of the theory.\cite{15,144,159,162}

**IX. CONCLUDING REMARKS**

Rosenfeld’s discovery of excess-entropy scaling more than forty years ago was based on those time’s primitive computer simulations. It provides an early example of a theoretical insight generated by simulation results, which preceded experimental confirmation. Rosenfeld’s explanation of excess-entropy scaling was based on the fact that simple liquids have virtually the same physics as the hard-sphere system. Thus he viewed excess-entropy scaling as a confirmation of quasiuniversality and, in particular, as a demonstration that quasiuniversality applies not just to the structure but also to the dynamics.\cite{1,158,180}

The hard-sphere system has a trivial potential-energy function, which is either zero or infinite, whereas the scaling properties of $U(R)$ [Eq. (19)] are central in isomorph theory. Despite this fundamental difference, the isomorph theory is fully compatible with Rosenfeld’s explanation of excess-entropy scaling based on simple liquids’ quasiuniversality. However, while Rosenfeld explained excess-entropy scaling from quasiuniversality, the isomorph theory in a sense does the opposite and explains quasiuniversality from the exponentially repulsive pair-potential system’s hidden scale invariance.\cite{15,185}

For the hard-sphere system, the packing fraction determines quantities like the reduced diffusion constant, viscosity, and thermal conductivity.\cite{1} While many suggestions exist for how to determine the effective HS-system packing fraction of a simple liquid at a given state point,\cite{191,192,193} Rosenfeld avoided this by referring instead to the system’s excess entropy, a purely thermodynamic quantity that for the hard-sphere system is in a one-to-one correspondence with the packing fraction.\cite{5,100} This is how excess-entropy scaling was arrived at; it also led to the proposal of a universal $S_{\text{ex}}$ dependence of the reduced transport coefficients.\cite{1} The exponential $S_{\text{ex}}$ dependence of Rosenfeld scaling Eq. (8) was not justified theoretically in Ref. 1 but reported as an approximate empirical fact; this is now known to apply only in some cases of excess-entropy scaling.

The basic message of hidden scale invariance is the irrelevance of the pair potential’s characteristic length scale.\cite{159} The Lennard-Jones potential-energy minimum is found at the pair distance $r = 2^{1/6}\sigma$, in which $\sigma$ is the characteristic length of the LJ pair potential. If the average nearest-neighbor distance is close to this value, the pressure is small. Hidden scale invariance reflects the non-trivial fact that the length $\sigma$ is irrelevant for the physics\cite{15,162} because the pressure is irrelevant for the reduced-unit structure and dynamics. By contrast, for the hard-sphere system the virial part of the pressure is determined by the packing fraction (equivalently: the excess entropy), and the pressure determines the excess entropy and the dynamical properties.\cite{200}

A quantitative measure of how well hidden scale invariance applies for a given system is provided by the virial potential-energy correlation coefficient $R$ of Eq. (20). Whenever $R$ is close to unity, excess-entropy scaling applies to a good approximation. If this is the case, the system obeys isochronal superposition, the property that two different state points of a given system with, e.g., the same dielectric loss-peak frequency have the same reduced relaxation functions.\cite{26,27,181}

Isomorph theory is not limited to explaining excess-entropy scaling for single-component pair-potential liquids; it applies also for solids, non-pair-potential systems, mixtures, molecules, confined systems, aging systems, etc. Note that $S_2$, as well as the higher-order configurational entropies [Eq. (9)], are all isomorph invariant because the reduced-unit structure is. This means that for tracing out the lines of invariant dynamics, i.e., the isomorphs, instead of keeping $S_{\text{ex}}$ constant, one might just as well keep $S_2$, $S_3$, or the so-called residual multiparticle excess entropy $S_{\text{RMPE}} \equiv S_{\text{ex}} - S_2$ constant.\cite{5,201}

Isomorph theory provides insights into why excess-entropy scaling may apply also for molecular systems quite different from the quasiuniversal simple point-particle liquids traditionally studied in liquid-state theory.\cite{3} The above “derivation” of isomorph theory from excess-entropy scaling showed that if the phase diagram has lines of the invariant structure and dynamics, these lines must be the configurational adiabats, i.e., excess-entropy scaling must apply.

It must be emphasized that not all observations in this intriguing field of research are explained by the isomorph theory. Rosenfeld scaling Eq. (8) cannot be derived from the hidden-scale-invariance condition Eq. (19), which is not surprising given that there are cases of excess-entropy scaling for which Rosenfeld scaling does not apply. Moreover, excess-entropy scaling or closely related regularities have been observed to apply also for a number of systems that do not have strong virial potential-energy correlations.\cite{35,51,103,134,135} Isomorph theory cannot explain this, at least not in its present version. There are examples of systems, e.g., with vibrational degrees of freedom modeled by harmonic springs, which do not have strong virial potential-energy correlations, but can be coarse-grained into a system that obeys isomorph theory by replacing harmonic molecular bonds by rigid bonds.\cite{174} More work is needed to clarify to which extent coarse-graining may lead to hidden scale invariance of the remaining degrees of freedom, and it remains an open question whether hidden scale invariance in one form or the other will eventually be able to provide a full explanation of excess-entropy scaling and related regularities.

**ACKNOWLEDGMENTS**

The author is indebted to Nick Bailey, Lorenzo Costigliola, Jack Douglas, Trond Ingebrigtsen, and Kristine Niss for suggestions improving the presentation. Heine Larsen provided invaluable technical assistance. This work was supported by the VILLUM Foundation’s Matter Grant (No. 16515).

\footnotesize
\begin{itemize}
\end{itemize}
E. Voyiatzis, F. Muller-Plathe, and M. C. Bohm, “Do transport properties scaling of transport properties in network-forming ionic melts (SiO$_2$).


