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The amorphous state equivalent of crystallization: new glass types by first order transition from liquids, crystals, and biopolymers

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Dedicated to Professor J.M. Hönl on the occasion of his 75th birthday

Abstract

We review the normal state of glasses and explain some exceptional cases by referring to a mode of glass formation, which is distinct from the normal and involves a first order transition route. Important materials like amorphous water and silicon belong to the distinct class, which we expect will prove to have many members, and which we expect will occupy a position part way between quasi-crystals and ordinary glasses. There may also be many mesoscopic examples of this class of material, because the low energy tertiary structures obtained by the (first order) folding of specialized heteropolypeptides (proteins) satisfy many of the criteria that we utilize in defining the class. The mesoscopic examples have the advantage of undergoing the transition to the low energy state under conditions of relatively long-lived metastability so that the phenomenon can be studied at leisure. There is no obvious reason why the phenomenon should be confined to biomolecules. We discuss the relation of the new glass types to ordinary glasses, plastic crystals, folding proteins and quasi-crystals, within the energy landscape paradigm. The first order transition occurs in the lower levels of the landscape in all cases, implying that ‘funnels’ are the general rule. © 2000 Éditions scientifiques et médicales Elsevier SAS. All rights reserved.

Keywords: Amorphous state; First order transition; Quasi-crystals

1. Introduction

Glassy solids form an important class of solid state material one of whose members, vitreous silica, is indispensable to our present day communications technology while others, known generally as ‘plastics’, dominate the engineering materials scene [1,2]. Despite the importance of glassy solids, the description of how rigidity is acquired by a random assemblage of molecules during cooling, is incomplete. As a result the glass transition is widely regarded as a

major unsolved problem in condensed matter physics.

In general, glassy solids are formed by the continuous cooling of liquids, though it is well known [1,3] that they may be formed by a number of other routes and that the dependence of the material properties on the route used in its formation is, in general, a minor one. A glass may be characterized as a monolithic non-crystalline solid, though often some restrictive additional qualifications, such as the ability to exhibit the glass transition, are attached. It has been our own predisposition to attach such restrictions because of the respectability that such an abil-

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ity to be cycled into and out of ergodic (internally equilibrated) states, implies [4,5]. However, it has always been awkward that certain undisputed glasses, like the silica analog BeF_2 , do not exhibit any glass transition [6,7], and even the archetypal glass, silica itself, manifests it very weakly compared with other systems [8,9]. For the purposes of dealing with the special cases of this paper, we will drop the glass transition requirement and simply require that the material, we call glass, can be formed directly from a condensed phase by non-chemical processes. This latter stipulation will avoid including amorphous materials that are formed as crystals too small to give diffraction peaks.

In most cases known, glasses form by arrest of whatever process is causing the vitrification, with an entropy that exceeds the entropy of the crystal of the same substance at the same temperature by some 30% or more. While this difference can change markedly with temperature as temperature is reduced down to 0 K [10,11], it does not disappear in the cases that have been fully characterized. In addition, the glassy state seems to have some other universal features. There is a ‘ubiquitous’ excess in the density of states at low frequency [12–14] over the value expected from harmonic vibrations, and there is a ‘ubiquitous’ violation of the Debye T^3 law [13–16], caused by the existence of very low energy states called TLS (two level systems) that can be reached by heavy atom tunneling processes [17]. Recently, however, two exceptions to these generalizations have been recognized [18–21] and the exceptions involve two of the most important ‘amorphous’ states known. This forces attention to the problem.

One of these exceptions is amorphous water, which is arguably the most common form of water in the universe [21], and the other is amorphous silicon [22], familiar to most as the black window on solar cell powered calculators. We will discuss these ubiquity violations and then correlate their existence with other exceptional behaviors of the liquid states of these two substances, water and silicon relative to other liquids which are normal glassformers.

The boson peak is most commonly detected as a low frequency resonance vibrational mode or group of modes revealed at $20\text{--}50\text{ cm}^{-1}$ after the Raman spectrum has been corrected for population effects by the Bose factor. This feature is very strong in scattering spectra (light and neutron) of the glasses

formed from so-called ‘strong’ liquids, and less so in those formed from ‘fragile’ liquids. Its interpretation is very controversial but generally seems to be intimately related to the disorder characterizing the glassy state. Intensity in the density of vibrational states in this region has been shown to build up as initially crystalline systems are systematically disordered beyond the ranges characteristic of the crystalline state [23]. It is, therefore, somewhat astonishing to find that the low frequency density of states of vitreous water, formed in bulk by the pressure-induced amorphization route [24], accurately follows the ω^2 dependence characteristic of the crystalline state, i.e. lacks any Boson peak. This is also true for the low density amorphous ice obtained by annealing the high density form at 100 K for many hours, in the manner first reported by Mishima [25].

It has, furthermore, been reported [19] that the low density vitreous water has a negligible density of TLS and that the density of TLS is also very small in the high density polyamorph. It was striking that a similar observation was reported, a little earlier, for the case of almost pure amorphous silicon containing about 1% hydrogen [20]. The dissipation of acoustic energy at low frequencies characteristic of glasses with Debye T^3 law violations (due to the TLS) is quite pronounced in hydrogenated amorphous silicon but drops off to background levels as the H content is reduced from 6 to 1% [20].

Thus, we find that there are two chemically quite distinct tetrahedral amorphous solids, both usually formed by unconventional routes, that exhibit solid state properties which are in striking distinction from those of the canonical glass. We now examine the relation of these systems in their *liquid* states to the liquid states of other glassformers, to see if there are any features there with which we may correlate their abnormal, and distinct, glassy state behavior. Based on the results of this examination, we will then project the existence of a rather broad range of substances, which will fall into the same new pattern and form a distinct subclass of glassy systems, which it will be interesting to characterize in more detail. It is expected that they will all show exceptionally strong annealing behavior after formation by rapid quenching, generally strong thermal and pressure history dependence, exceptionally low excess entropy, and generally crystal-like character while remaining strictly aperiodic. With the exception of

their single molecule counterparts discussed below (the proteins) they may also, rather unfortunately, lack stability in the temperature range above (or even approaching) their glass transition temperatures. Fortunately, this does not preclude their providing technologically important materials (e.g. a-Si solar cells).

2. Exceptions to the normal glassformer viscosity pattern

The viscosity data for a range of liquid systems, including water and silicon, are summarized in Fig. 1 [26], which is the T_g -scaled Arrhenius plot used in the classification of liquids according to the ‘strong/fragile’ pattern of behavior [27].¹ If the data for water and the weakened silica analog, BeF_2 , were dropped from Fig. 1 we would have the normal viscous liquid pattern for canonical glassformers. This is the pattern on the basis of which many correlations of other relaxation properties, such as

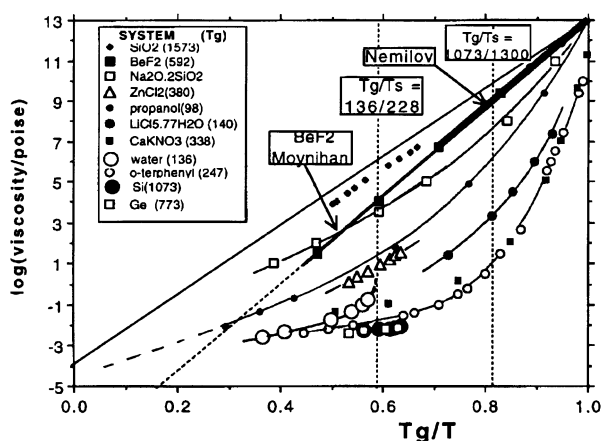


Fig. 1. The T_g -scaled Arrhenius plot representation of the viscosities of a selection of liquids, highlighting the exceptional behavior of two ‘tetrahedral’ liquids, BeF_2 and water. The temperature 228 K which power law fits of the viscosity and other data for water suggest as a singular temperature [25], is indicated by a vertical dashed line. The second vertical line corresponds to the temperature assigned to the first order liquid–liquid transition in silicon, see text. The assignment of scaling temperatures for Si and Ge is discussed later in an appendix.

¹ Note that the first use of the T_g -scaled Arrhenius plot for liquid viscosities was made by [28]. The corresponding T_g -scaled excess entropy plot [123] seems, on the other hand, to be a new development.

non-exponentiality of relaxation functions [30], non-linearity of annealing functions [31,32], and prominence of short time dynamic features [33], have been based. Now, however, we are interested in the exceptions. Liquid silicon and liquid germanium would simply appear to be exceptionally fragile liquids, which, like many molecular liquids included in the original data presentation [27] (see [3] for a reproduction) cannot be studied much below their melting points because of easy crystallization. However, we will see that they are very different.

Note the manner in which the data for pure BeF_2 , [33–35] in Fig. 1 follow the Arrhenius law with an unphysical pre-exponent. Water deviates from the norm in a different manner. Starting from high temperatures where the behavior appears fragile, the data for water suddenly commence a rapid increase shortly below the freezing point, and tend to diverge according to a power law [29] at a temperature 228 K that is far above the glass transition temperature of 136 K. This value of T_g is obtained both by direct measurements on the vapor-deposited and hyperquenched forms [37], and by extrapolation of solution data of which a great many exist [38–42].

The data for liquid silicon, deserve additional comment. The authors of these difficult viscosity measurements [43,44]² noted that, near the lower end of the temperature range, deviations occur from the Arrhenius behavior which is characteristic of most other liquids of low viscosity. The deviations are in the same sense as those characterizing supercooled water in Fig. 1, though they cannot be seen on the scale of Fig. 1. It is unfortunate that the supercooled range could not be explored in the laboratory study of silicon because computer simulations studies to be discussed further below show that striking behavior, even more pronounced than that of water, would be found.

We should note that the manner in which the Si and Ge data are placed on the diagram depends on the manner in which we assess a scaling temperature T_g for Si. This will be dealt with in Appendix A dealing with diffusivity behavior. By our assignment, the viscosities of liquid Si are seen to be lower than

² These authors report anomalous increases of viscosity near T_m . Molecular dynamic studies [27–29] show anomalous decreases in the corresponding diffusivity with decreasing temperature but only at lower temperatures near the temperature of the liquid-to-liquid transition anticipated in Section 2.

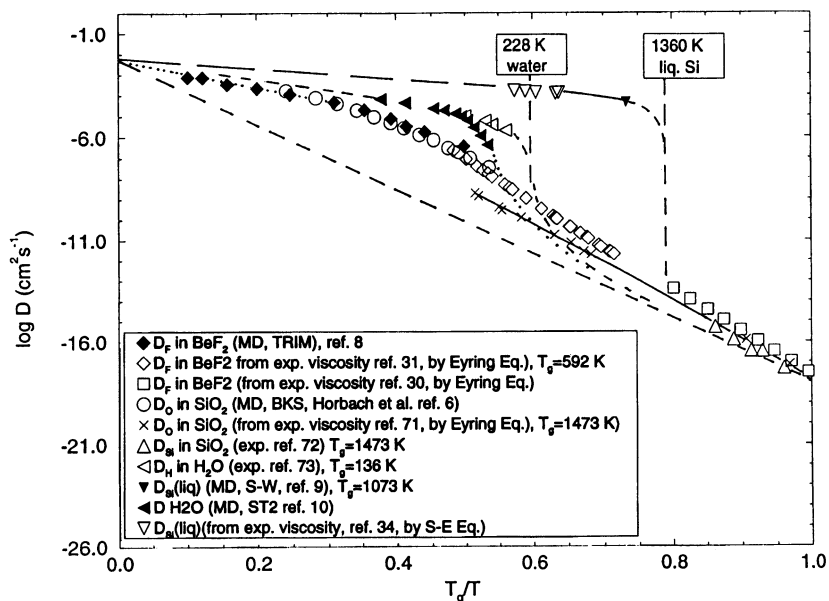


Fig. 2. A new pattern for tetrahedral liquids, based on the data of Fig. 1 cast in diffusivity form. These data are augmented by diffusivity data at high temperatures from computer simulation studies. Note the progressive deviation from the pattern of Fig. 1 as the constraints on the network topology tighten. This happens as the bridging species (O,F,H) diameter decreases and finally (Si) is absent altogether. In this case there is a first order ‘freezing’ transition to a highly viscous, still amorphous, state.

those of any of the other liquid phases at the same T_g/T , inviting the designation ‘superfragile’.

To bring some order into these otherwise perplexing twists in the viscosity behavior, we turn to the results of computer simulations on these substances and focus on the corresponding diffusivities. The simulations permit us to explore ranges of temperature, which are both to the low and to the high ends of existing experimental data. These provide much clarification and allow a new and important pattern to emerge. Data are presented in Fig. 2.

We use diffusivities in Fig. 2 for several reasons, as follows. Firstly, diffusivity is the most fundamental transport process, occurring as it does in the absence of any driving force, and involving only a single particle correlation function. Secondly, and because of the first, it is the process most easily and most accurately extracted from molecular dynamics computer simulations on which the key observations in Fig. 2 depend. Thirdly, use of the diffusivity criterion for fixing T_g (see below) permits us to arrive at reasonable estimates of T_g for substances for which it has not yet been (or cannot be) measured, in particular for amorphous silicon, a-Si.

To convert the (high) viscosity data of Fig. 2 to (low) diffusivity data for Fig. 6, we use the Eyring equation [45] which is based on a jump migration model and is known to apply well for oxide glasses and liquids [46]. It is appropriate in a domain where MD studies show the diffusivity process to be episodic in character. The Eyring equation is

$$D_i = \frac{k_B}{\eta \lambda_i} \quad (1)$$

where λ_i is the jump distance, taken to be the diameter of the oxide (or fluoride) anion, 0.28 and 0.27 nm, respectively. For very low viscosities, molecular dynamics studies show diffusion is continuous. In this case the more familiar Stokes–Einstein equation, which predicts lower diffusivities (by ca. one order of magnitude) for a given viscosity than does the Eyring equation, is probably more appropriate. We use it to convert the high temperature viscosity data for Si [43,44] to diffusivities for comparison with the results of simulations [47–49] and for plotting in Fig. 6.

As in the case of liquid silicon, the data for SiO_2 and BeF_2 at very high diffusivity values have been obtained by means of molecular (ion) dynamics com-

puter simulations [50–53]. The data for SiO₂ are already in the literature [52,53]. In the case of water, on the other hand, some data for rather *low* diffusivities are taken from MD simulations, in this case from the study of Paschek and Geiger [54]. These data are possible because the rapid crystallization of laboratory water does not happen in ST2 water. Whether or not this is because of the rigid molecule form of ST2 is still uncertain.

Fig. 2 contains a lot of information, and many details and justifications are needed. These are given in the Appendix A, parts (a) and (b), which provide background to, as well as explanation of Fig. 2 data and plotting.

The important features of Fig. 2 are the following. Firstly we see that, for BeF₂ shortly above the highest experimental viscosity point of [35,36], the data enter a regime of rapid curvature that allows D to recover the normal high temperature limit. As might be expected from Eq. (2), this is a domain of high heat capacity, as is shown elsewhere [26]. Secondly, the pattern of Fig. 2 implies the existence of systematic effects, which we relate to the size and character of the bridging entities in the tetrahedral network. As the polarizability or size of the bridging unit (O, F, and H in SiO₂, BeF₂, and OH₂, respectively) decreases, the cooperativity of bond-breaking increases, and the departure from the normal excitation pattern of Fig. 1 becomes more pronounced. The relatively mild temperature dependence of D at high temperatures persists closer to the glass transition temperature (where $\langle \ell^2 \rangle / D$ reaches 200 s [55]³), and the final decrease towards $D = 10\text{--}18 \text{ cm}^2 \text{ s}^{-1}$ at T_g becomes more precipitous. Finally, in the case where there is no bridge at all, the case of liquid silicon, the overconstraint and the consequent cooperativity of bond breaking maximizes, and the extreme of a first order transition to the low excitation state is realized [47–49,56]. Elsewhere [57] this cooperativity has been correlated with the narrowness of the inter-tetrahedral bond angle distribution in the

amorphous phase. The existence of the first order transition in a-Si was first suggested on the basis of thermodynamic data both by Spaepen and Turnbull [58] and by Bagley and Chen [59], and later verified by the flash heating experiments of Thompson et al. [60]. It occurs only ca. 20% above the glass transition temperature as we have assigned it (see (a) below). The temperature of the liquid–liquid transition is about 1345 K though there is some uncertainty due to the interference of crystallization [61].

3. Modeling the liquid–liquid phase transition

Elsewhere [62,63] the behavior seen in Fig. 2 has been modeled by applying the Adam–Gibbs equation [64] to the results of a simple excitations model for the thermodynamic properties [65–67]. This model, in zeroth order, has recently [68,69] been shown to provide a good account of the excess entropy of a wide variety of glassforming liquids, and thus to provide a resolution of the Kauzmann paradox for glassformers. Use of this entropy temperature relation in the Adam–Gibbs equation, Eq. (2) below, then accounts at least qualitatively for the diffusivity/viscosity behavior of the normal glassformers of Fig. 1. The Adam–Gibbs equation for the diffusivity D is

$$D = D_o \exp \frac{C}{TS_c} \quad (2)$$

where D_o and C are constants and S_c is the configurational entropy. In experimental tests of the theory, S_c is usually taken as the excess entropy of the liquid over crystal, but this may differ considerably from the configurational entropy defined by Gibbs because of the presence of a system dependent vibrational component in the entropy difference [10,11].

The origin and importance of this component, and others, is under active discussion [68,69] as it may well be the principal determinant of the fragility of the liquid. It may also (see below) be the factor that, by determining the value of ΔS° in the excitations model, determines how close to T_g the temperature of the phase transition lies when the excitations of the quasilattice are cooperative. Since most successful tests of the Adam–Gibbs equation have utilized the excess entropy for S_c , there would seem to be a problem, and some have suggested that Eq. (2) should be regarded as an empirical relation. On the

³ In Zener's treatment of anelastic relaxation in solids [55] relaxation is held to occur via relaxation modes of timescale $\tau = \langle \ell^2 \rangle / D$, D being the diffusion coefficient and ℓ the wavelength of the mode. The wavelength of the dominant mode will be $2\pi/q_o$ where q_o is the q value of the main peak in the structure factor, hence will be of the order of the intermolecular spacing. Taking this as 0.3 nm, we find that, for $\langle \tau \rangle = 200$ s, D should be $2 \times 10^{-22} \text{ m}^2 \text{ s}^{-1}$ (the enthalpy relaxation time at the 10 K min^{-1} onset T_g is usually about 200 s according to [32]).

other hand, tests of the equation based on simulation studies where the quantity used in Eq. (2) is strictly the configurational entropy assessed from the inherent structures [70,71] have also proved to be successful. Thus the configurational entropy and the associated vibrational entropy evidently scale, as excitations models would require in the lower temperature regime.

Cooperativity is taken into account by extension of the bond lattice model to first order in the mixing free energy of bonds and broken bonds. In the zeroth order model, it is assumed that the bond-breaking energy is independent of how many bonds are already broken and it has been surprising how closely the excitation equations based on this seemingly naïve assumption, actually describe the majority of glassformers [65–67]. However, the cases of Fig. 2 are clearly exceptional and here we must assume that the mixing is less simple.

Let us assume, for simplicity, that the free energy of mixing of bonds and broken bonds behaves like that of the different molecular species in a non-ideal binary solution. Then the empirical observation [72,73] that the partial molar free energy of a component in such solutions varies as the square of its mole fraction will also describe first order deviations from ideal mixing of the bond lattice excitations. In this case, the excitable bond lattice problem maps onto the familiar regular solution model of binary solu-

tions. The cooperativity increases as the strength of the first order mixing parameter W increases. The state of excitation at temperature T is then obtained, as for the zeroth order model, by finding the extremum in the derivative of the free energy of the two-state system with respect to the fraction X_B of excited states. In the cooperative case this means solving the equation

$$\frac{\partial G}{\partial X_B} = \Delta H^\circ - T\Delta S^\circ + RT \ln \frac{X_B}{1 - X_B} + W(1 - 2X_B) = 0 \quad (3)$$

in which the only term distinguishing the non-ideal case from the ideal case is that containing the temperature independent interaction parameter W [65–69]. As W increases the excitation profile [70,71] (which is the plot of configurational energy vs. temperature) changes systematically from convex to concave towards the T axis until it generates a critical point at a temperature $W/2R$ (R the gas constant). For larger W , a first order transition between liquid states in different states of excitation occurs, the strength of the transition increasing with increasing W . As is obvious from Eq. (3), all curves must pass through the same point at $T = \Delta H^\circ / \Delta S^\circ$. Since the bond breaking energy parameter ΔH° is the principal determinant of T_g , the larger the value of ΔS° the closer to T_g the phase transition will occur.

This treatment yields a phase transition that is well understood in the binary solution context, and which has also been invoked for a number of conceptually distinct models of single component systems with phase transitions [74–79], all of which use the same algebraic representations. It is the lowest order member of a series of non-ideal binary solution treatments, which give successively more quantitative descriptions of the observed phase coexistence in such systems. Since the single component systems studied to date cannot be observed in detail because of their unsuccessful competition with the dominant crystallization process, refinements of the model are not yet in order. Elsewhere [53], we have discussed how it may be possible to observe the phase transition in compressed BeF_2 under difficult but thermodynamically stable conditions, in which it may be possible to obtain quantitative data to justify model refinements. We will not pursue that question here. Rather, we want to stress that the regular solution model is not the only way of parameterizing cooperativity in a single component excitation processes

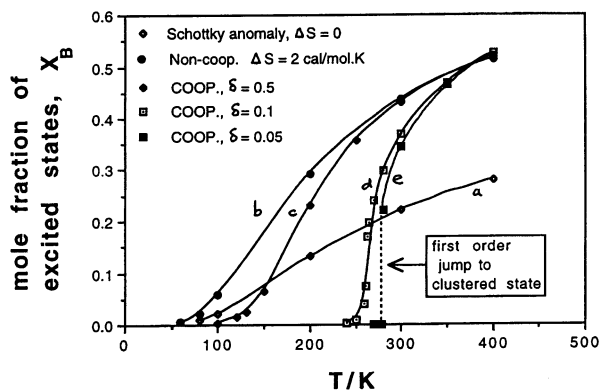


Fig. 3. Excitation profiles for the bond lattice model with fixed ΔH° and (curve a) no cooperativity and no elementary entropy of excitation (i.e. Schottky anomaly), (curve b) no cooperativity and a positive entropy of excitation, (curve c, d, e) and with the same entropy of excitation and three different degrees of cooperativity according to Eq. (4) (curves a, b, and c). For case (c), the de-excitation changes from fragile-but-continuous to first order discontinuous. (From ref. 60, by permission.)

with a single additional parameter. Fig. 3 shows the excitation profile obtained from a treatment [80] in which it is supposed that the excitation free energy is a function of a parameter δ according to the relation

$$\Delta G = \frac{\Delta G_o}{X + \delta} \quad (4)$$

The interesting feature of this model is that, when the first order transition from the high temperature phase occurs, the system drops into a low temperature phase in which the defect population (i.e. broken bond, or excitations, fraction) is extremely low (Fig. 3). In this respect, it is more like a crystallization than a liquid–liquid phase transition. If this parameterization should be appropriate to the first order transition in liquid silicon then it would be reasonable to think of the new phase as rather distinct thermodynamically from the regular glassy solid. It is notable in this connection that the entropy of low density amorphous water has been assessed by three different groups [81–83] using different experiments, and is found to be only marginally different from that of the stable crystal at the same temperature.

The thrust of this discussion has been to indicate that the glasses produced by systems which configurationally de-excite in a highly cooperative manner, particularly by outright first order phase transition, might be in thermodynamic states more akin to crystals than to glasses. The relative uniqueness of their configurations may in turn be the source of their exceptional and crystal-like densities of states and low temperature dynamics. Of course, it could be argued that these features obtain because the systems are indeed nanocrystalline, but this is certainly not the view of a large body of amorphous Si researchers. In this spirit, we will, therefore, consider the wider implications of our line of thought.

4. The range of systems with properties like a-Si and a-H₂O

There are many cases of substances with silicon-like interactions and water-like behavior [84] waiting to be studied in detail. From high pressure studies, it can already be told that the first-order transition between liquid states seen for liquid silicon must also occur for GaSb and InSb under the appropriate conditions. Indeed, it was for such substances that

the ‘two-liquids’ model of Rappoport [74,75] was first applied with specific focus on the predicted phase transition [76]. The analysis of Ponyatovsky and Barkalov [76], in particular, used the known band-gap energies and density differences between crystalline analogs of these substances to predict the approximate temperature/pressure relations for such a phase transitions.

Here, we have been concerned with bringing this phenomenon, and the range of amorphous solids, which it makes possible by pressure cycling [25,76,79,85], into systematic relation with the better understood field of liquid glassforming systems.

All the substances of Fig. 2 have in common a temperature at which their densities have a maximum value (TMD). Glazov and co-workers long ago [84] showed that there are a number of members of the liquid semiconductor family, which had the same anomaly. Most, but not all, of these share the same inverse relation between energy and molar volume that characterizes Si, and indeed the related binary analogs GaAs, AlP, and even CdTe. It will indeed be interesting to carry out experiments and simulations on these materials to see the extent to which the pattern of Fig. 2 can be extended. A limiting case would be AlP, because, in first row compounds BN and elemental diamond, the liquid state would have three coordination as a preference or option. Boron, however, can be incorporated by compounding with As, according to Wentzcovitch et al. [86,87].

Ponyatovsky and co-workers have made many characterizations of the conditions in which the amorphous semiconducting states of these materials may be obtained by first order, or pseudo first order, transition from metastable *crystalline* states of the same composition [76,79,85]. While the crystal-to-glass transition is to be distinguished qualitatively from the liquid–liquid transition that we have discussed, it is certainly closely related. Indeed, a model which has much in common with those of [74–79], is that of Granato [88,89] in which the phase transition of the model is considered to be one between crystal and amorphous states, i.e. a regular melting transition. We feel confident that when the annealed states of the tetrahedral amorphs formed by pressure destabilization of the high pressure crystalline metallic states are studied, the class properties we have discussed earlier will be found in most if not all of these cases. As in the case of water, with its crystal-to-glass transition under pressure and its two

polyamorphic forms, the occurrence of polyamorphic forms with the implied liquid–liquid transition at higher temperatures, should be a common feature of these semiconducting glasses.

Beyond the tetrahedral amorphs of the silicon genre, lie examples of liquid–liquid transitions in insulating glasses, and even metallic glasses [90–92]. In each of the latter cases, the low temperature glassy phase has been observed to form as droplets from the parent liquid during rapid cooling. Each of the low temperature phases is prone to crystallization, and in the metallic case it is a quasi-crystal that is formed. The low temperature properties of the annealed, expanded, phases have yet to be determined, but the fact that these glasses can be obtained in bulk form directly from the liquid adds much interest to their characterization.

5. Mesoscopic examples. The folding of proteins

It would be interesting to increase the *molecular* members of our proposed class of systems whose distinction from ordinary glasses is that they de-excite via first order phase transitions. To date we have referred only to the case of water. To this end we inquire into the extent to which the phenomena we have discussed might occur in polymeric systems. a-Si after all can be regarded as a giant inorganic polymer.

In general, chain polymer systems either crystallize from the melt or vitrify by continuous relaxation time increase, more or less like regular glassformers. There are, however, an intensively studied group of chain polymers, of moderately high molecular weight, that undergo abrupt first order transitions from states of high disorder to states of much lower disorder. They distinguish themselves from the systems we have considered so far by doing this on an individual molecule-by-molecule basis. We refer, of course, to the class of heteropolypeptides developed by nature to perform **specialized** catalytic and structural functions in living matter. These protein systems are, on an individual level, of a complexity comparable with, indeed exceeding, that of a macroscopic liquid sample. The latter can, after all, be represented to high accuracy by a periodic box containing only a few hundred particles, far fewer than in a single biomolecule, which cannot be reduced to any simpler unit.

These mesoscopic single molecule glassformers undergo their first order phase transitions in a reversible manner under the appropriate conditions, reaching the same low energy structural state (or, better, the same megabasin of connected states [93,94]) every time. For each molecule there is a particular temperature at which the folding and unfolding rates are the same, and the population of the two states are, therefore, equal [95,96]. This is the denaturation temperature, which is not sharply defined as in a macroscopic melting process because of the fluctuations characteristic of small systems [97]. Nature has of course designed these systems to be kinetically protected from passage into lower energy states, usually of tendril or amyloid plaque form, and the occasional breaching of this kinetic barrier has serious or fatal consequences. Examples are ‘folding diseases’ like ‘mad cow disease’ (BSE), and Parkinson’s disease [98]. The refolding to a state, which then easily forms tendrils at body temperatures appears to be heterogeneously nucleated in the BSE case [99], the mysterious ‘protein X’ [100] serving as the heterogeneous nucleus [93]. However, recent studies of the effect of prolonged heating of a number of proteins near but below the denaturation temperature, show the tendril formation to be of rather general occurrence [101]. Under these circumstances, the tendril formation does not require prior refolding of the individual protein molecules. On the other hand, a nucleation event to initiate the aggregated tendril state is presumably required and this step is probably homogeneously nucleated. Reorganizations into tendril and amyloid forms of proteins are in many ways the equivalent of the crystallization processes that bedevil the study of the liquid–liquid transitions in supercooled water and silicon.

The difference in thermodynamic properties of folded and unfolded states of proteins may be rather similar to the difference in properties of high and low temperature phases of a molecular system with liquid–liquid phase transition (or close approach to one as in the case of water). This is illustrated in Fig. 4, which is reproduced from [93]. Fig. 4 contains data from [102,103] on the heat capacity of a globular protein before and after the unfolding induced by the increase in temperature. The protein, legumin, had been dehydrated, hence did not refold on cooling. Then the increased heat capacity of the unfolded state is seen to be retained until the temperature of the glass transition is reached at ca. 130°C, where-

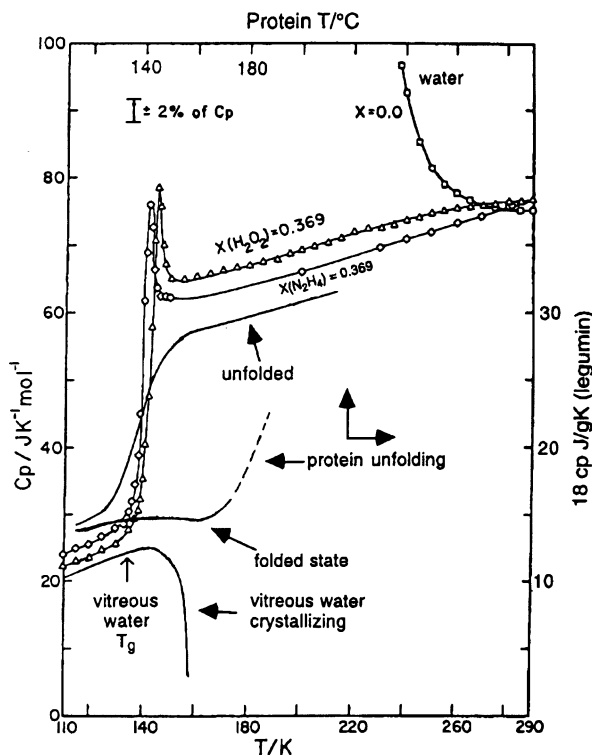


Fig. 4. Comparison of the thermodynamic character of the folded and unfolded protein legumin (top and right hand scales) with that of supercooled water with and without second components, H_2O_2 and N_2H_4 . These components hydrogen bond like water but destroy the cooperative de-excitation, hence preserving the aqueous system in the 'unfolded' state down to the glass transition temperature. The multiplying factor on the right hand scale is to change the specific heat capacity to values more comparable to the molar scale for the aqueous systems. (From ref. 93 by permission.)

upon it drops to the (vibrational) value of the folded protein. In the aqueous systems, the same behavior can be obtained by interfering, by solute additions, with the source of the cooperative deexcitation to the low temperature phase that occurs in pure water. Either hydrogen peroxide or hydrazine will achieve this end without changing the specific heat capacity much at all. Then instead of the upswing seen in pure water as $T < 270$ K (which is to water what the early part of the folding exotherm is to proteins), the heat capacity remains high but steady until the glass transition is reached at 140 K. On the scale of Fig. 4, the weak increase in heat capacity of amorphous water, that is taken as the signature of the glass transition at 136–150 K [37], is barely noticeable, just as a glass transition in the folded, but dehy-

drated protein cannot be detected in normal scanning. The anneal-and-scan method used by Mayer and colleagues [37] to 'bring out' the glass transition in amorphous solid water (which is not normally detectable [39]) has not yet been applied to dry proteins.

The thermodynamic properties of a mole of folded proteins should be dominated by the properties of the individual molecules. Thus if folding were to produce a ground configurational state for the individual molecule then the molar properties of the protein should mimic the molar properties of the ideal glass and the substance would conform to the third law of thermodynamics, to good approximation. Of course, the natural function of proteins requires that the folding produce a state in which much configurational freedom remains. However, this is a design feature of nature, which should not deflect attention from the fact that systems of this type, with single molecule phase transitions, are in fact very close in principle to the macroscopic systems with phase transitions that we have described in the earlier part of this article. Folding into much lower energy states could be induced by change of residue sequence [101], and such systems could be of considerable academic interest.

It is surely only a matter of time (well in advance of J.M. Honig's 100th) before inorganic equivalents of the proteins are developed and this whole phenomenology of folding (i.e. of single molecule phase transitions) can be shifted to higher temperatures. Then it will become a new and interesting part of solid state chemistry (and, presumably, also industrial catalysis technology).

6. Energy landscape considerations

In the search for understanding of the behavior of complex systems, it has been found helpful to invoke the higher dimensional ($3N + 1$ dimensions, N the number of particles) energy landscape [104–106]. Where a three-particle system can only have a single minimum energy configuration, complex systems can find of the order of e^N energetically distinct packings when the volume can freely be chosen. In the case of systems of the type we have considered here, there are certain volumes, intermediate between those of the two coexisting phases, in which the packing energies available are distinctly higher than those of

neighboring volumes, and in which those energies are connected in a continuous downhill path to those of the adjacent phases. This corresponds to the instability of structures of intermediate volume indicated by solutions of the equation of state that fall in between the two spinodal limits (superheating of the low temperature phase, and supercooling of the high temperature phase) [65]. In Fig. 5 [93] this circumstance has been represented by the region of high energy in the middle of the diagram, which is a cut across the $3N$ dimensional surface.

Fig. 5 has the configuration coordinate traversing regions of changing volume (it should be noted that it is simpler and more common to consider only the states available in a single volume, as is the situation for most of the MD computer simulations studies being reported at this time [50,51,106–108]). Fig. 5a suggests the sparse packing in energy of states in the low temperature phase, consistent with the evidence that these are ‘strong’ liquids, while the dense packing in the high temperature phase indicates the fra-

gility found in such cases (e.g. liquid silicon in Fig. 1).

The lower part of Fig. 5 shows one of the two common representations of the energy surface for protein folding [109–112]. Shakhnovitch [109,110] has argued that the folding pathway always includes, as a key element, the formation of a small knot of strongly interacting residues, which, once formed, direct the remainder of the folding process to an inexorable final state. This is the folding nucleus, a specially important grouping of residues that has the ability to promote the subsequent formation of a three-dimensionally stable structure from the chain molecule. The same group often occurs at the core of proteins which have quite different biological functions, and the same groups are strongly conserved in evolution [109,110,114]. Once the nucleus has arisen via fluctuations within the high temperature phase of the cooling protein, the remainder of the process of forming the low energy state follows quickly [109–112]. This is the same situation that prevails in the

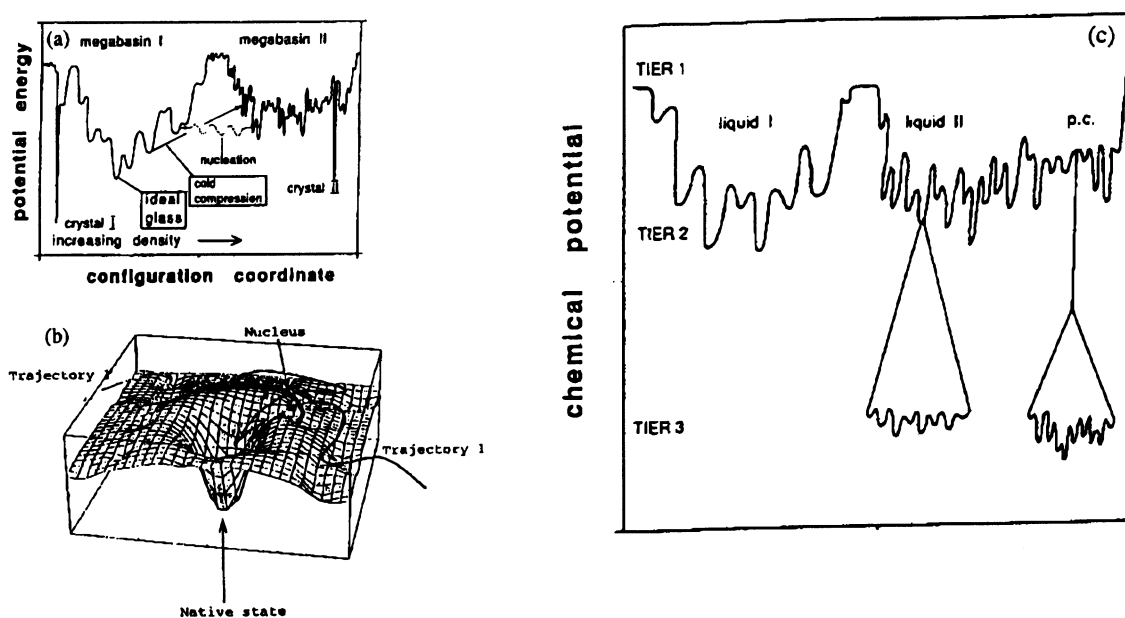


Fig. 5. Energy landscape depiction of the energetic relations between the different microstates of complex systems such as those considered in this article. Part a shows a cut through the $3N$ dimensional surface, in which the density of the system increases along the configuration coordinate, and energetically unfavored states interpose between the megabasins of configuration space belonging to the two alternative liquid phases. Part b shows the Shakhnovitch representation of the folding protein energy surface [86] in which the key feature, apart from the deep minimum of the folded state, is the presence of a narrow minimum energy pathway, corresponding structurally to passage through a specific nucleus. The deep minimum corresponds to an aperiodic structure in both protein and polyamorphic glass cases. (From refs. 69 and 86, by permission.) Part c shows the situation for a plastic crystal glassformer, in which the multiple packing states and slow relaxation occur within a crystalline megabasin. So far, plastic crystal glass formers are found to be intermediate and strong in character [89,90] rather than fragile like so many liquids and chain polymers. (From ref. 90 by permission.)

supercooling of water, and it is the reason that it is difficult to supercool below the temperature -40°C [115]. In nature, the need for an easy transition between the unfolded and folded states is obvious and the engineering of an efficient nucleus has, therefore, been a priority in evolution. According to the ‘folding funnel’ picture of protein folding [111,112], this process has been developed so efficiently that the nucleation barrier has virtually disappeared. The folding funnel notion could equally well be taken to describe the energetics of water or silicon as they pass from high temperature states to the states occupied just before the phase transition. In each of these cases, most of the bonds characteristic of the low temperature state are in place before the phase transition is reached, meaning that the systems are far down in energy on their landscapes before the final step is taken. It is in the final step, however, that the structures are locked into place, and the entropy of the remaining disorder is lost.

On the other hand, the folding funnel scenario for understanding how systems are guided towards first order transitions to their configurational ground states seems equally applicable to the crystallization of ordinary glassformers, and in this sense there is a paradox. We have seen that the position of T_c on the energy landscape of mixed LJ, is closer to the ground state in energy than it is to the top of the landscape [65,106]. There are a number of glassforming liquids for which the temperature of melting of the crystal is close to, or even below, the value of T_c [3]. These are also cases of crystals whose entropies of fusion are rather small. Therefore, the final crystallization would appear to be guided by a funnel energy landscape, though we know that these are some of the best glassformers available. Particularly in the case of the archetypal glassformer SiO_2 , the normal melting point would appear to be well below the value of T_c [50,51]. This is why the entropy of fusion of SiO_2 is so much less than the value R per mole of particles expected for the case where melting gives access to $\exp(N_{\text{av}})$ states per mole of particles (thought to be typical of the number of alternate packing minima per mole of simple particles). With so few minima available, the relaxation time is long (according to Eq. (2)) and so the crystallization rate becomes small. The paradox is that the same scenario utilized for explaining how a seemingly improbable event, the folding of a protein, will occur, can be used to explain how a seemingly probable

event (the crystallization) will not occur. The resolution must be, as emphasized in [112], that the folding temperature must not be too close to T_g or the folding will not occur. The conclusion of this line of thought, however, is that ‘folding funnels’ are the rule rather than the exception for systems with weak first order transitions.

Finally, we consider the case of glassforming rotator phases, often called plastic crystals, in which the orientational states can in some cases become slow to access, whereupon a very similar phenomenology to that of liquid glassformers is encountered. In the plastic crystal state [113,116], the system is initially isolated in a crystalline megabasin whose energy landscape is almost as complex as that of the typical liquid glassformer. Though these systems are center-of-mass ordered, they still exhibit the same ‘canonical glassy dynamics’ features as liquids. However, ‘strong’ character is much more common than for liquids, and unlike liquids, many molecular examples are included in the ‘strong’ class, e.g. thiophene [114], and neopentanol [118].

For a quasi-crystal, on the other hand, notwithstanding the lack of real space periodicity, the system exists within a crystal basin which has very little structure apart from some low lying minima due to defects (as in ordinary crystals).

7. Conclusions

We have considered an aspect of glassy state phenomenology that is not well known or understood at this time but which we think deserves some attention. This aspect involves systems which de-excite, on temperature decrease, in a discontinuous rather than in the continuous manner of normal glassformers. This may occur by happenstance of the interaction potential or, in the case of biomolecular systems, by design of nature. We have given evidence of a series relationship between several liquids in which the inter-particle potentials direct them towards tetrahedral coordination. Differences in the strength of constraints on the tetrahedrality cause them to de-excite with increasing sharpness down the series, culminating in a distinct phase transition between macroscopic states of different excitation level in the case of the most constrained system, Si. The sharpness of the transition in the case of silicon and water is related here to the uniqueness of the struc-

ture into which these constraints direct the system on cooling. This helps explain the crystal-like characteristics that they exhibit in the absence of three-dimensional periodic ordering. Our development suggests that there should be many other examples of such aperiodic solids that form by first order transition. We argue that they are all phenomenologically related to the mesoscopic structures formed by the folding of single molecule systems of the protein type. Finally, we envisage the development of intermediate cases that are single molecule systems with phase transitions but are inorganic in nature and which may well provide single molecule catalysts for future technology.

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Appendix A

Here we give some details and justifications for the data entries for water and liquid silicon in Fig. 2.

(a) The glass transition temperature for a-Si (and a-Ge).

In order to place the diffusion coefficient data for liquid silicon on Figs. 1 and 2 it is first necessary to have a glass transition temperature for amorphous silicon (a-Si). Since this material always crystallizes at temperatures around 800 K on reheating, the use of T_g of 1070 K in construction of Figs. 1 and 2 might seem arbitrary. The value used is based on the diffusivity data on crystalline silicon that are summarized by Frank [118], and reproduced in Fig. 6.

Fig. 6a shows that the self-diffusivity of Si in crystalline Si reaches the value characterizing the glass transition temperature, $D = 10\text{--}18 \text{ cm}^2 \text{ s}^{-1}$, at a temperature of 1170 K. It seems improbable that the amorphous state of Si could have a self-diffusivity smaller than that of the organized crystalline state (though some impurities diffuse faster in c-Si than in a-Si [118,119]). Thus we assume T_g lies below 1170 K, but not by too much. We assign $T_g = 1070 \text{ K}$. However, if we use the higher

value 1170 K ($T_g/T_{\text{li}} = 0.87$), the qualitative aspects of Fig. 4 are not affected. Quantitatively, the anomaly is sharpened since the diffusivity jump then occurs closer to T_g . Qualitative support for our choice of a lower value is provided by a comparison, shown in Fig. 5c, with the case of ice and amorphous water, for which T_g is known [37].

In the case of Ge, which has been reported to show many of the same anomalous liquid and glassy features that we have described for a-silicon, diffusivity data are also available over wide ranges for the crystal [118]. In this case, it is interesting to observe (from Fig. 5b) that the T_g value $D = 10\text{--}18 \text{ cm}^2 \text{ s}^{-1}$, falls at a temperature 540°C, that is much closer to that reported for the crystallization of amorphous Ge, 450°C [120] than in the case of Si. Since again the diffusivity for the amorphous phase would be as high or higher, this implies that liquid-like flow could probably be observed in a-Ge before it crystallizes. This is consistent with the reports of structural relaxation observed in a-Ge [121]. Some form of penetrometric viscosity measurement on a-Ge in the temperature range where the structural relaxation is observed needs to be performed on a priority basis to support the picture we have presented here.

(b) The diffusivity of supercooled water near 150K.

The manner in which we have drawn the diffusivity of H₂O below 220 K might also seem arbitrary in view of the absence of any data in that region, and particularly in view of the diffusivity data published recently for H₂O near 150 K [122], which show very different behavior. Our representation is based on the analysis of the glass transition width [123], supported by the very small heat capacity increase seen at the glass transition, both of which argue for extremely strong liquid behavior in this range. This is supported by the analysis of Starr et al. [63] of the behavior of the thermodynamic properties, entropy and heat capacity, of H₂O in the ‘forbidden’ region between 150 and 230 K needed for the measured values of these properties at 150 and 230 K to be correct. Applying the Adam–Gibbs equation to these data using parameters from above 230 K, Starr et al. [63] obtain data similar to, but steeper than, those shown in Fig. 6. We have added the requirement that $D = 10\text{--}18 \text{ cm}^2 \text{ s}^{-1}$ at $T_g = 136 \text{ K}$ as an additional constraint.

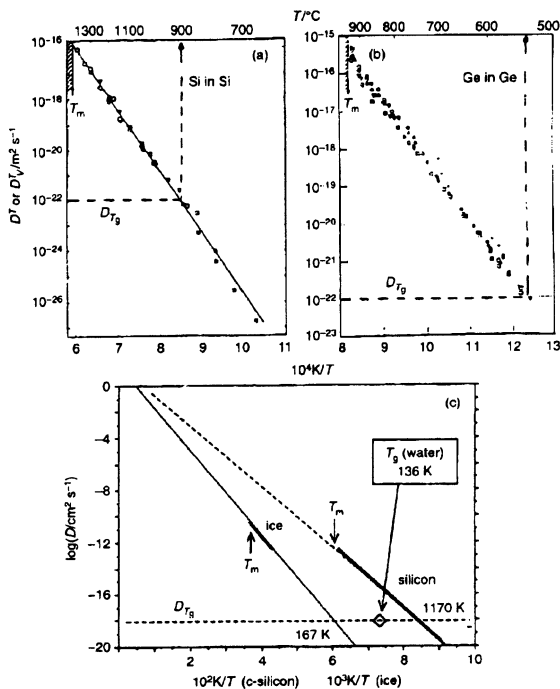


Fig. 6. Diffusivities in (a) crystalline Si and (b) crystalline Ge showing the temperatures at which the diffusivity reaches the value characteristic of the glass transition temperature in non-fragile glassforming systems, and indicating the maximum value considered feasible for the glass transition temperatures of the amorphous states of silicon and germanium. (c) The relation between diffusivities in ice and c-Si and the position of the glass transition in a-H₂O, 136 K. Diffusivities for ice are taken from ref. 98. The glass transition temperature is that usually attributed to water (see ref. 32).

Support for this diffusivity versus temperature form can be found in the data of Pashek and Geiger [54] from MD simulation of ST2 water. The ST2 potential is known to shift the important phenomena of water some 35 K to higher temperatures in the vicinity of the density maximum. We would accordingly expect to see a rapid drop in $D_{\text{H}_2\text{O}}$ at $230 + 35 \text{ K}$ (0.513 on the T_g -reduced scale of Fig. 2) if the proportionality holds to lower temperatures. Pashek and Geiger indeed found a rapid decrease, but at 0.55 on the Fig. 2 scale, which is only some 15 K above that of the laboratory water. It does not follow a power law as the experiments seem to. However, these authors also observe an extremum in the expansivity, which would imply a maximum in heat capacity (at 265 K, $T_g/T = 0.53$) below which Eq. (2) would predict a return to Arrhenius behavior such as we have drawn in Fig. 6. Thus the apparent

Arrhenius behavior reported by Pashek and Geiger at their lowest temperatures could be due to an inflection in the Arrhenius plot that is masked by computational noise, c.f. the inflection in the Arrhenius plot of [63]. Unfortunately, the data at lower temperatures needed to resolve this ambiguity will be very difficult to obtain because of the very long run times involved. As far as they go, however, we see these ST2 data as offering considerable support for the course of $D_{\text{H}_2\text{O}}$ that we have shown in Fig. 2.

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