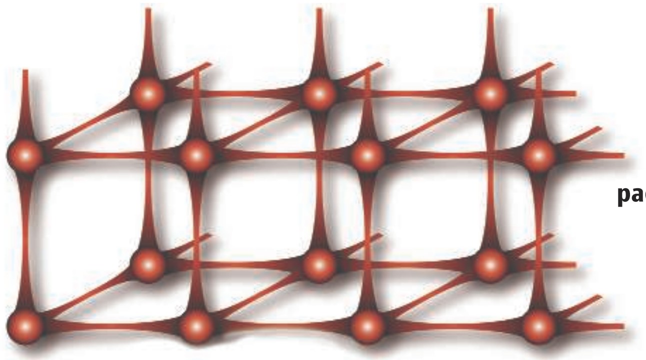
page 1211

INTRODUCTION

Quantum Wonderland

LIKE ALICE AND HER WONDERLAND, PHYSICISTS ALSO HAVE ACCESS TO TWO worlds: the classical and the quantum. Although both worlds are inhabited by the same two species, bosons and fermions, their behavior in either world can be remarkably different. The macroscopic or classical world is filled with the familiar and modeled with classical laws. Lowering the temperature sufficiently to enter the quantum world reveals that these species can interact in cooperative ways, giving rise to exotic phases of matter—quantum matter—not seen in the classical world. Here, things get interesting (and weird); solids, liquids, and electrons can flow without dissipation; exotic phases can emerge; fluctuations can be critical; and objects can be entangled and be in multiple places at once.

Many experimentalists and theorists have been exploring this quantum regime for some time now, studying how individual particles and ensembles of particles behave, in attempts to unravel the underlying physics producing these exotic properties and phases. Some others are heading straight to applications. The six Perspectives in this special section provide a taste of some of the topics that occupy the world of quantum matter.

With atoms trapped in a lattice of optical microtraps, Bloch (p. 1202) discusses how the ability to manipulate the magnitude and sign of the interaction between the atoms can provide a model system in which to explore the formation of the exotic phases seen in quantum gases, liquids, solids, and electronic and magnetic systems. Leggett (p. 1203) sets out the theoretical basics of such quantum systems, explaining how their behavior depends on which family of statistics (Bose-Einstein or Fermi-Dirac) the atoms belong to. Choosing the example of quantum criticality in fermionic systems, Zaanen (p. 1205) points out that the fermions and their statistical family are troublemakers. Trying to explain the complexity emerging from what are simple constituents, he tells us that our present mathematical toolbox is incapable of describing how these exotic electronic phases emerge and that new mathematical tools need to be developed. Another recent example of an observation in need of an explanation is the supersolid effect found in helium-4, where a solid crystal seems to move like a superfluid. Chan (p. 1207) presents the latest on this new phase and argues that imperfections in the crystal appear to be necessary for the effect to be seen. Communication is a vital technology in the classical world, and Walmsley (p. 1211) describes how developments made in the quantum world are carrying over to applications through the use of quantum optics in areas such as secure communication and cryptography. Lloyd (p. 1209) expands on the topic of communication and information, describing how quantum information can be considered as matter, as concrete as any of the matter we are familiar with in our classical world, and how theoretical ideas in quantum error correction will lend themselves to the realization of an operational quantum computer.

Outside the special section, Adrian Cho's story in News Focus (p. 1180) describes research in Fermi condensates, gases composed of fermionic atoms, which may help researchers model materials as diverse as high-temperature superconductors and the interiors of neutron stars.

So, armed with an Alice-like curiosity, let's take a short walk in this quantum landscape.

– IAN OSBORNE AND ROBERT COONTZ

Quantum Matter

CONTENTS

Perspectives

See also related News story page 1180

Science

PERSPECTIVE

Quantum Gases

Immanuel Bloch

Ultracold quantum gases are proving to be a powerful model system for strongly interacting electronic many-body systems. This Perspective explores how such atomic ensembles can help to unravel some of the outstanding open questions in the field.

When matter is cooled down close to zero temperature, particles can interact in a cooperative way and form novel states of matter with striking properties—superconductors, superfluids, or fractional quantum Hall liquids. Similar phenomena can now be observed in a dilute gas of atoms, five to six orders of magnitude less dense than the air surrounding us. Here, degenerate bosonic and fermionic quantum gases trapped in magnetic or optical traps are generated at temperatures in the nanokelvin regime (*1*). Whereas initial research concentrated on weakly interacting quantum states [for example, on elucidating the coherent matter wave features of Bose-Einstein condensates (BECs) and their superfluid properties], research has now turned toward strongly interacting bosonic and fermionic systems (*2*, *3*). In these systems, the interactions between the particles dominate over their kinetic energy, making them difficult to tackle theoretically but also opening the path to novel ground states with collective properties of the many-body system. This has given rise to the hope of using the highly controllable quantum gases as model systems for condensed-matter physics, along the lines of a quantum simulator, as originally suggested by Feynman (*4*).

Two prominent examples have dominated the research in this respect: (i) the transition from a superfluid to a Mott insulator of bosonic atoms trapped in an optical lattice potential (*5*–*7*) and (ii) the BEC–Bardeen-Cooper-Schrieffer (BCS) crossover of a two-component Fermi gas across a Feshbach resonance through which the magnitude and sign of the interactions between pairs of atoms can be tuned (*8*–*11*). In the first, a weakly interacting and superfluid gas of quantum degenerate bosons can be turned into an incompressible and insulating gas in a three-dimensional lattice of optical microtraps. The Mott insulator can be visualized as a many-body system in which strong repulsive interactions between the particles sort them into a perfectly ordered array and each lattice site is occupied by a single atom. In the second example, pairs of fermionic atoms can form bosonic composite particles when their interactions are tuned by Feshbach resonances. Such bosonic composites can themselves undergo Bose-Einstein condensation, thus fundamentally altering the properties of the many-body system. When a true two-body bound state exists between the particles, the composite bosonic particle is simply a molecule, albeit very large, whereas in the case of attractive interactions without a two-body bound state the composite pair can be seen to be related to a BCS-type Cooper pair, which can then undergo condensation. It is the possibility of changing almost all the underlying parameters dynamically and the ability to model the complex many-body quantum systems by first principles that have led to a surge in experimental and theoretical research.

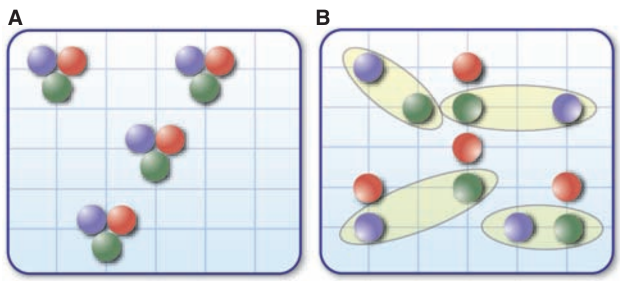

Fig. 1. Three-species fermionic atoms (red, green, and blue spheres) in an optical lattice can form two distinct phases when the interactions between the atoms are tuned. In the first case of strong attractive interactions between the atoms, they join as "trions" (**A**), whereas in the second case of weaker interactions, a color superfluid is formed (**B**), in which atoms pair up between only two species. The two phases have strong analogies to the baryonic phase (A) and the color superfluid phase (B) in quantum chromodynamics [see (*12*)].

What is next on the agenda? For fermionic systems with and without a lattice, researchers are trying to see whether they can pair up particles with very different mass ratios, such as lithium and potassium, or possibly even three different fermionic atomic species. This line of work is fueled by a theoretical prediction that such fermionic mixtures could show phases in which three fermions join to form a "trion" analogous to quarks forming baryonic matter or, alternatively, only two of the fermionic components pair to form a "color" superfluid (*12*) as in quantum chromodynamics (Fig. 1). For the case of two particles with highly different mass ratios, one hopes to observe exotic forms of superconductivity such as the Fulde-Ferrell-Larkin-Ovchinnikov superconducting phase (*13*, *14*), where particles condense into pairs with nonzero momentum. Early experiments have produced degenerate mixtures of two fermionic atomic species (*15*) and two fermionic species with an additional third bosonic component (*16*), and both are progressing quickly toward exploiting Feshbach resonances to control the interactions between the fermionic atoms.

For lattice-based systems, efforts are under way to explore the feasibility of using ultracold atoms as quantum simulators for strongly interacting many-body systems. For example, in the famous class of high-Tc superconductors, such as the CuO compounds, one observes that these form antiferromagnetically ordered ground states when undoped. Upon doping, and thereby changing the effective filling in the system, the antiferromagnetic order is destroyed and a superconducting phase with *d*-wave symmetry of the order parameter emerges (*17*) (Fig. 2). What exactly happens during the transition and how it can be described theoretically is currently a subject of heated debates and one of the fundamental unsolved problems in the field of condensed-matter physics. Cold-atom researchers are currently trying to determine whether they can help to resolve some of these issues (*18*). As a starting point, several groups are preparing to observe antiferromagnetically ordered states in two-component Fermi mixtures in an optical lattice. To achieve this, however, one needs to cool the many-body system to challenging temperatures T below the superexchange interaction energy J_{ex}, which characterizes the coupling strength between the spins of atoms on neighboring lattice sites. If the temperature is not low enough, thermal fluctuations would simply destroy the fragile magnetic order present in the ground state. Superexchange interactions form the basis of quantum magnetism in strongly correlated electronic media and can be described as an effective spin-spin interaction between the neighboring particles on a lattice (*19*). They are a result of virtual "hopping" events of particles to neighboring lattice sites, in which a particle tunnels to an adjacent site and then the same particle—or its neighbor—returns to the original lattice position. For two spin-polarized fermions, such hopping is suppressed by the Pauli principle, whereas for two fermions with opposing spin directions, the hopping is allowed and leads to

Johannes-Gutenberg Universität, 55118 Mainz, Germany. E-mail: bloch@uni-mainz.de

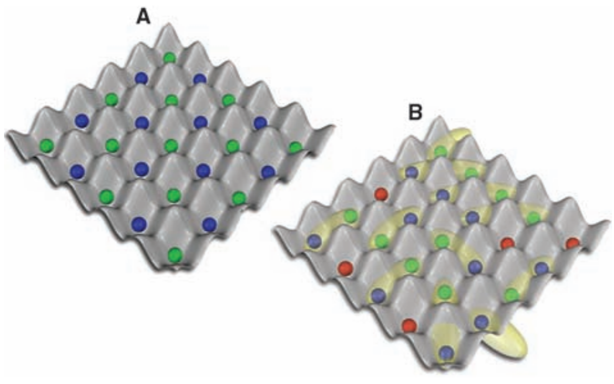

Fig. 2. For two-component Fermi gases (blue and green spheres) in an optical lattice, an antiferromagnetic ground state is expected at half-filling (**A**). Upon doping (red spheres), such an antiferromagnetic order is expected to be destroyed and, in some theories, a spin-liquid state emerges (**B**), which can form the basis for a high-Tc superconducting phase [see (*17*, *29*)]. Researchers with ultracold atoms are currently trying to establish both phases in the experiment.

a decrease of the total energy of the particles. For fermionic atoms, an antiferromagnetic orientation of the atoms with alternating spins is thus favorable over a ferromagnetic one. Exchange and superexchange interactions between ultracold atoms have been observed, showing that they can be controlled to a high degree and that their coupling strengths can be in the kHz range (*20*, *21*). The temperatures of the Fermi gases produced in the experiments so far also seem to be initially low enough to enter an antiferromagnetically ordered phase, such that one can expect to observe these phases in upcoming experiments (*22*). Whether the temperatures one needs are, however, already also low enough to observe a *d*-wave superconducting phase is unclear, especially as current theories do not permit a precise estimate of the critical temperature for entering the superconducting phase. By comparing to critical temperatures observed with typical high-Tc superconductors, one can estimate the required temperatures to be a fraction of the superexchange coupling. Progress in this direction might therefore be crucially linked to novel approaches for cooling the quantum gases to even lower temperatures in the lattice (*23*).

Control over the effective spin-spin interactions between neighboring atoms could also open up a new avenue for the simulation of quantum magnetism with cold atoms or molecules. Both atoms and molecules offer the ability to implement arbitrary spin Hamiltonians on a lattice (*24*). For atoms, the spin-spin interactions are generated by superexchange couplings (*25*, *26*), whereas for ultracold molecules the electric dipole-dipole interaction can mediate even stronger spin coupling between individual molecules on neighboring sites (*27*). Heteronuclear Feshbach molecules have recently already been formed in optical lattices; however, in order for them to possess a strong electric dipole moment, the molecules that are up to now created in highly excited vibrational states will have to be brought to the ground state in a controlled way. Laser spectroscopy in molecules could do exactly this: by using a single pulse or a sequence of two-photon Raman transitions, the high excitation energy could be removed and the molecules could be brought to the ground state. If all this can still be done in such a gentle fashion that does not heat the atoms and molecules too much, and whether one will be able to ultimately obtain a degenerate gas of heteronuclear molecules stable enough to carry out experiments, remain to be seen.

Finally, several research teams are currently trying to find ways to address and observe single atoms on single lattice sites (*28*). In the optical regime, this requires a demanding optical microscope, as the atoms in the lattice are only spaced by half a micrometer or less. However, if successful, one not only would be able to observe but also could control a spin system in two dimensions with 10,000 particles simultaneously in view, all with single-site and single-atom resolution. Observing dynamical evolutions in these systems, probing their spatial correlations, and finally implementing quantum information processing in such a truly large-scale system would offer exciting prospects for future research.

References

1. L. Pitaevskii, S. Stringari, *Bose-Einstein Condensation* (Oxford Univ. Press, Oxford, 2003).
2. D. Jaksch, P. Zoller, *Ann. Phys.* **315**, 52 (2005).
3. I. Bloch, J. Dalibard, W. Zwerger, http://arXiv.org/pdf/0704.3011.
4. R. P. Feynman, *Int. J. Theor. Phys.* **21**, 467 (1982).
5. M. P. A. Fisher, P. B. Weichman, G. Grinstein, D. S. Fisher, *Phys. Rev. B* **40**, 546 (1989).
6. D. Jaksch, C. Bruder, J. I. Cirac, C. W. Gardiner, P. Zoller, *Phys. Rev. Lett.* **81**, 3108 (1998).
7. M. Greiner, O. Mandel, T. Esslinger, T. W. Hänsch, I. Bloch, *Nature* **415**, 39 (2002).
8. C. A. Regal, M. Greiner, D. S. Jin, *Phys. Rev. Lett.* **92**, 040403 (2004).
9. M. W. Zwierlein *et al.*, *Phys. Rev. Lett.* **92**, 120403 (2004).
10. M. Bartenstein *et al.*, *Phys. Rev. Lett.* **92**, 120401 (2004).
11. T. Bourdel *et al.*, *Phys. Rev. Lett.* **93**, 050401 (2004).
12. A. Rapp, G. Zarand, C. Honerkamp, W. Hofstetter, *Phys. Rev. Lett.* **98**, 160405 (2007).
13. P. Fulde, R. A. Ferrell, *Phys. Rev.* **135**, A550 (1964).
14. A. I. Larkin, Y. N. Ovchinnikov, *Sov. Phys. JETP* **20**, 762 (1965).
15. E. Wille *et al.*, http://arXiv.org/pdf/0711.2916.
16. M. Taglieber, A.-C. Voigt, T. Aoki, T. W. Hänsch, K. Dieckmann, *Phys. Rev. Lett.* **100**, 010401 (2008).
17. P. A. Lee, N. Nagaosa, X.-G. Wen, *Rev. Mod. Phys.* **78**, 17 (2006).
18. W. Hofstetter, J. I. Cirac, P. Zoller, E. Demler, M. D. Lukin, *Phys. Rev. Lett.* **89**, 220407 (2002).
19. A. Auerbach, *Interacting Electrons and Quantum Magnetism* (Springer, Berlin, 2006).
20. M. Anderlini *et al.*, *Nature* **448**, 452 (2007).
21. S. Trotzky *et al.*, *Science* **319**, 295 (2008).
22. F. Werner, O. Parcollet, A. Georges, S. R. Hassan, *Phys. Rev. Lett.* **95**, 056401 (2005).
23. A. Greissner, A. J. Daley, S. R. Clark, D. Jaksch, P. Zoller, *Phys. Rev. Lett.* **97**, 220403 (2006).
24. M. Lewenstein *et al.*, *Adv. Phys.* **56**, 243 (2007).
25. L.-M. Duan, E. Demler, M. D. Lukin, *Phys. Rev. Lett.* **91**, 090402 (2003).
26. A. B. Kuklov, B. V. Svistunov, *Phys. Rev. Lett.* **90**, 100401 (2003).
27. A. Micheli, G. K. Brennen, P. Zoller, *Nat. Phys.* **2**, 341 (2006).
28. K. D. Nelson, X. Li, D. S. Weiss, *Nat. Phys.* **3**, 556 (2007).
29. P. W. Anderson, *Science* **235**, 1196 (1987).

10.1126/science.1152501

PERSPECTIVE

Quantum Liquids

A. J. Leggett

Quantum liquids are systems in which not only the effects of quantum mechanics but also those of the characteristic indistinguishability of elementary particles are important. The most spectacular of these are the systems of bosons (liquid ^{4}He, the Bose alkali gases), which undergo the phenomenon of Bose condensation, and the fermion systems (liquid ^{3}He, the electrons in some metals), which display the related phenomenon of Cooper pairing. I discuss these phenomena and the relation between them.

A quantum liquid may be defined as a many-particle system that shows not only the effects of quantum mechanics but also those of quantum statistics. As is well known, general considerations concerning the rotation group limit the possible values of total angular momentum that can be possessed by any "elementary" particle, or any compound system such as an atom made up of elementary particles, to $n\hbar/2$ where n is an integer (including zero). If n is even, the particle or system in

Department of Physics, University of Illinois at Urbana-Champaign, Urbana, IL 61801, USA.

question is called a boson; if odd, a fermion. A famous theorem of quantum field theory, the spin-statistics theorem (*1*), then states that the total wave function of any many-particle system must be even under the interchange of all the coordinates of any two bosons of identical type, and odd under interchange of those of any two identical fermions. Formally, if i and j label two particles of the same species, and $\mathbf{r}_i\boldsymbol{\sigma}_i$, etc., their space and spin (or other internal) coordinates, then we must have

$$\Psi(\mathbf{r}_1\boldsymbol{\sigma}_1, \mathbf{r}_2\boldsymbol{\sigma}_2..., \mathbf{r}_i\boldsymbol{\sigma}_i..., \mathbf{r}_j\boldsymbol{\sigma}_j ...\mathbf{r}_N\boldsymbol{\sigma}_N) = \pm\Psi(\mathbf{r}_1\boldsymbol{\sigma}_1, \mathbf{r}_2\boldsymbol{\sigma}_2...\mathbf{r}_j\boldsymbol{\sigma}_j ...\mathbf{r}_i\boldsymbol{\sigma}_i...\mathbf{r}_N\boldsymbol{\sigma}_N) \quad (1)$$

with the ± sign applying to bosons (fermions). In the special case of free particles that can occupy plane wave states with momentum $\boldsymbol{k}$, spin projection σ, and energy $\varepsilon_{k\sigma}$, the condition (Eq. 1) leads, for a collection of identical particles of a single species in thermal equilibrium at temperature $T \equiv 1/k_B\beta$ (where k_B is Boltzmann's constant), to a single-particle distribution $n_{k\sigma}$ of the form

$$n_{k\sigma}(T, \mu) = [\exp\beta(\varepsilon_{k\sigma} - \mu) \mp 1]{-1} \quad (2)$$

where μ is the chemical potential and the ∓ sign now refers to bosons (fermions). The distribution (Eq. 2) with the minus sign is known as the Bose-Einstein distribution and that with the plus sign as the Fermi-Dirac distribution; hence, bosons (fermions) are often said to satisfy Bose (Fermi) statistics. However, it needs to be emphasized that the "statistics" of Eq. 2 apply only to a very special case, whereas the requirement of Eq. 1 of symmetry or antisymmetry of the many-particle wave function is much more general. I will follow the inaccurate but conventional practice of referring to the consequences of Eq. 1 as those of (quantum) statistics.

A necessary condition to see nontrivial effects of Eq. 1 is that the system should show appreciable effects of quantum mechanics in the first place, i.e., should deviate appreciably from the behavior predicted by a purely classical description. Crudely speaking, this is likely to happen when the thermal energy k_BT falls below a typical single-particle excitation energy. For example, if we describe an insulating crystalline solid by the Einstein model, in which each atom vibrates in the potential field of its neighbors with a frequency ω_o, then the condition is about $k_BT \lesssim \hbar\omega_o$; a more sophisticated (Debye) model confirms this result, in the sense that (for example) the specific heat of a crystalline solid falls below the classical "equipartition" value of $3k_b$ per atom when $k_BT \lesssim \hbar\omega_D$, where the Debye frequency, ω_D, is of the same order as ω_o. More generally, a very rough order of magnitude for the temperature T_d at which quantum-mechanical effects begin to show up in an important way may be obtained by imagining each particle of mass m to move in a cage of side $a \sim n^{-1/3}$ (n = particle density) formed by its neighbors; because the typical single-particle excitation energy is then of order $\hbar^2/ma^2$, the criterion is

$$T_d \sim \frac{\hbar^2 n^{2/3}}{mk_B} \quad (3)$$

This criterion is approximately correct for liquids and gases (in the latter case, the temperature T_d, which can be estimated more rigorously from an analysis of Eq. 2, is often called the degeneracy temperature); it underestimates T_d somewhat for solids because there the effective size of the "cage" tends to be substantially less than the interatomic spacing. In any case, on putting in the values of n and m, it is clear that the condition $T \lesssim T_d$ is satisfied for electrons in any liquid or solid below the vaporization temperature and for atoms in any liquid or solid at cryogenic, but nowadays relatively easily attainable, temperatures; it can also be satisfied for ultracold atomic gases, as in these systems, although the density is many orders of magnitude smaller than that in a typical solid or liquid, the temperature is also much less.

However, although the criterion $T \lesssim T_d$ is certainly a necessary condition to see the effects of quantum statistics (i.e., the constraint Eq. 1), it is by no means sufficient. In fact, if we compare the behavior of a carbon crystal made of the "common" isotope ^{12}C (a boson) with that of one composed of the rare isotope ^{13}C (a fermion), the only difference is a trivial one associated with the slightly different masses. In order for the statistics to have an effect, it is essential that identical particles are able to change places. A nice example of this principle (*2*) is seen in the structure of the vibrational and rotational levels of diatomic molecules composed of chemically identical atoms, such as C_2: if we consider a heteronuclear molecule such as ^{12}C–^{13}C, then, because there is no question of exchange of "identical" particles, the constraint (Eq. 1) has no effect and all levels are allowed. If now we replace (for example) the ^{13}C atom by a second ^{12}C, we find that the vibrational levels are unaffected (except trivially, via the difference in reduced mass) but that the odd-angular-momentum rotational levels are missing! This underlines spectacularly the difference between a process such as rotation, in which the identical atoms physically change places, and one such as vibration where they do not. Note that for the "statistics" to have an effect it is not necessary that the two identical atoms ever occupy the same position at the same time. [For a more detailed discussion, see for example (*2*), section 1.1.]

Thus, a quantum liquid is a many-particle system in which (i) the temperature is less than or of the order of the T_d defined by Eq. 2 and (ii) the particles can change places relatively easily. As so defined, the category of quantum liquids includes the ultracold dilute atomic gases as a special case; however, because they are the subject of another essay in this issue (*3*), I will confine myself here to systems occurring at typical liquid or solid densities. The category then includes the electrons in metals, the two stable isotopes of helium (the only element that remains liquid under its own vapor pressure in the limit $T \to 0$), and if we are willing to leave the terrestrial sphere, the neutrons in neutron stars and possibly more exotic forms of matter such as quark stars.

Rather generally, the properties of any quantum liquid are likely to be quantitatively and even qualitatively different from those of the corresponding classical system; this has long been known in, for example, the case of the electrons in metals at temperatures of the order of room temperature ($T << T_d$), which are actually often described surprisingly well by the simple (Sommerfeld) model of noninteracting fermions, which leads to Eq. 2. However, the most spectacular manifestations of the effects of quantum statistics are associated with the phenomenon of Bose-Einstein condensation (BEC) and the related phenomenon of Cooper pairing occurring in Fermi systems. For a noninteracting gas of bosons described by the distribution of Eq. 2, a straightforward analysis originally carried out by Einstein (*4*) shows that below a temperature T_0 of the order of T_d a nonzero fraction of all the N particles, that is, a macroscopic number $N_0(T) \sim N$, occupies the lowest single-particle state (in free space, this is the zero-momentum state). It has long been believed that a similar phenomenon occurs, in thermal equilibrium, in a system of interacting bosons, provided that the interaction is overall positive (repulsive) and that it is just this that is happening in the "superfluid" (He-II) phase of the bosonic liquid ^4He. In recent years, direct evidence for BEC has been obtained in dilute ultracold atomic gases such as ^{87}Rb and ^{23}Na; these gases are actually confined in a harmonic trap, and in the absence of interactions BEC would show up as a much-enhanced ($N_0 \sim N$) population of the harmonic ground state, leading to a sharp spike in the density distribution around the origin. In real life, this spike is somewhat broadened by the repulsive interatomic interactions, but it can still be clearly seen in the experiments (*5*). An interesting feature of BEC in the atomic gases is that theory suggests, and experiment confirms, that it can occur even when the system is far out of thermal equilibrium and the macroscopically occupied state is thus strongly time-dependent.

In qualitative terms, the BEC state is characterized by the fact that a macroscopic number N_0 of particles are forced to occupy the same single-particle state and thus to behave in exactly the same way [compare with (*6*)]. This property leads to a variety of spectacular effects, including the complex of phenomena known as superfluidity, which is observed to occur in the He-II phase of liquid ^{4}He, and a variety of interference phenomena, which it has become possible to observe in the ultracold atomic gases [(*2*), section 2.5].

Turning now to Fermi systems, we see from Eq. 2 that in this case the value of $\langle n_{k\sigma} \rangle$ can never exceed 1 (the Pauli principle), so that the direct analog of BEC certainly cannot occur. However, there is no reason why a complex made up of an even number of fermions (a boson) cannot undergo BEC (indeed, this is exactly what is happening in, for example, ^{87}Rb at ultralow temperatures), and in particular there is every reason to believe (compare below) that if two spin-1/2 fermions are coupled by an attractive interaction sufficiently strong to bind them into a spin-0 bosonic molecule, these molecules will indeed undergo BEC; such a scenario might be imagined to describe, for example, liquid D_2, if we could exclude crystallization. Imagine now that we gradually weaken the interfermion attraction to the point where (in the two-body problem) the molecule is no longer stable and even beyond. Can a sort of BEC still persist under these conditions? In their epoch-making work (*7*) in 1957, Bardeen, Cooper, and Schrieffer (BCS) showed that the answer is yes: A degenerate system of fermions with an arbitrarily weak attraction will, at sufficiently low temperatures (exponentially low compared to T_d), form "Cooper pairs," a sort of giant diatomic (or more accurately dielectronic, because BCS were dealing explicitly with the electrons in a metal), spin-0 molecules, and the latter will then in effect automatically undergo the phenomenon of BEC. In contrast to the case of (hypothetical) liquid D_2, however, the size of the "molecules" is now large compared with their average separation, so that the theory of Cooper pairing is quantitatively and even qualitatively quite different from that of BEC of tightly bound diatomic molecules.

When Cooper pairing occurs in an electrically neutral system of fermions such as liquid ^{3}He, the consequences are qualitatively similar to those of BEC in a bosonic system such as ^{4}He. Indeed, it is almost universally believed that the anomalous phases of liquid ^{3}He that occur below 3 mK (note this is $\ll T_d \sim 1$ K), which show many of the manifestations of superfluidity, are indeed characterized by the onset of Cooper pairing. When the latter occurs in the electrically charged system of electrons in metals, the effects are even more spectacular: In particular, the metal in question will exhibit the two major effects characterizing superconductivity, namely, persistent flow of currents in a ring and the exclusion of magnetic flux (Meissner effect), which leads to the possibility, inter alia, of static magnetic levitation. Superconductivity, originally thought to be an intrinsically low-temperature effect, has in the past 20 years been observed to occur in a class of cuprate materials up to around half of room temperature; although the detailed explanation of this high-temperature superconductivity is still furiously debated, there seems little doubt that its fundamental origin lies in the phenomenon of Cooper pairing.

As explained above, it is somewhat natural to think of the phenomenon of Cooper pairing in a system of fermions with weak attraction on the one hand and BEC in the system of diatomic molecules formed from them on the other as opposite ends of the same spectrum, and it has long been speculated that by "tuning" the strength of the attraction one might be able to realize a continuous transition between the two situations; this is known as the "BEC-BCS crossover." In the past 4 years, by using the phenomenon of Feshbach resonance, it has become possible to study the BEC-BCS crossover experimentally in ultracold atomic gases, and it indeed appears to be continuous as tentatively predicted by theory: See the article by I. Bloch in this issue (*3*). Thus, we now have a very satisfying unification of the concepts of BEC in a bosonic system and Cooper pairing in a fermionic one.

References and Notes

1. R. F. Streater, A. S. Wightman, *PCT, Spin and Statistics and All That* (Benjamin, New York, 1964).
2. A. J. Leggett, *Quantum Liquids: Bose Condensation and Cooper Pairing in Condensed-Matter Systems* (Oxford Univ. Press, Oxford, 2006).
3. I. Bloch, *Science* **319**, 1202 (2008).
4. A. Einstein, *Abh. Preuss. Akad. Wiss.* **3**, 18 (1925).
5. M. R. Andrews *et al.*, *Science* **273**, 84 (1996).
6. Molecule of the Year for 1995, *Science* **270** (1995).
7. J. Bardeen, L. N. Cooper, J. R. Schrieffer, *Phys. Rev.* **108**, 1175 (1957).

8. This work was supported by the NSF under grant no. NSFDMR-03-50842.

10.1126/science.1152822

PERSPECTIVE

Quantum Critical Electron Systems: The Uncharted Sign Worlds

J. Zaanen

Phases of classical matter, such as solids and liquids, are ruled by emergence principles that are well understood. Although the same principles govern forms of quantum matter that have no secrets for physicists, such as the superfluids, having to deal with fermions and the associated Fermi sign problem shatters this analogy. This Perspective addresses the Fermion sign problem and describes experiments on metals undergoing quantum phase transitions exhibiting scale-invariant electronic behavior, a description of which is at odds with established quantum theory.

Ice is different from water, and water is different from steam, although these phases of matter are all made from the same water molecules. Countless numbers of molecules are required to make this work, and the various phases of matter are said to "emerge." Emergence is at its best when the transition between such phases is continuous and the system no longer has any sense of preference for one or the other phase. This lack of "executive power" has the consequence that the system spontaneously adopts the powerful symmetry of scale invariance. In this "critical state" the system looks on average the same, regardless of the amplification factor that is used to observe it (*1*).

The concept of emergence is so powerful that it transcends the classical-quantum divide, and this apparently includes the generation of scale invariance: Quantum critical states where quantum fluctuations drive a phase transition at zero temperature are now routinely observed (*2*–*6*). However, dealing with emergence in quantum physics requires one to consider the organizational principles of quantum statistics, as discussed by Leggett (*7*), whereby quantum particles are either bosons or fermions. Despite its underlying quantum properties, bosonic matter is ruled by the same emergence principles as classical matter (*2*). In stark contrast, Fermi statistics wreck this analogy, and the emergence principles governing fermionic matter are among the great mysteries of modern physics. Fortunately, experimentation can help: Electrons in solids are relatively easy to probe, and they form systems of countless numbers of strongly interacting fermions. The recently observed quantum phase transitions in a variety of metals (*3*–*6*) reveal that fermionic quantum matter can exhibit unexpected behavior: Particles tend to acquire an infinite mass, and the scale-invariant fermionic phases that take over appear to be the birthplace of new forms of stable quantum matter.

Regular matter is formed from a large number of quantum particles, electrons, quarks, and

Instituut-Lorentz for Theoretical Physics, Leiden University, Leiden 2333 CA, Netherlands. E-mail: jan@lorentz.leidenuniv.nl

so forth. To answer the question of where classical integrity comes from, we look to Feynman's path integral formalism (*2*), a rather pictorial view of the quantum world but one that does reproduce accurately all the known facts about quantum systems in equilibrium. Within this formalism, quantum systems resemble classical matter living in a higher-dimensional "Euclidean space-time" (Fig. 1) having an extra axis, "imaginary time" (our time times the square root of –1). The maximal duration of this imaginary time is the ratio of Planck's constant divided by temperature: When temperature is lowered, "more time is available" to see quantum behavior. The strength of the quantum fluctuations is analogous to temperature in classical physics, having the effect of "heating up" the "stuff" inside space-time. When temperature is lowered and the quantum fluctuations are sufficiently vigorous, this frozen quantum matter might melt. The resulting space-time liquid will appear to our eyes as, for instance, a superconductor. When this melting transition taking place in space-time at zero temperature generates scale invariance, the quantum critical state is realized.

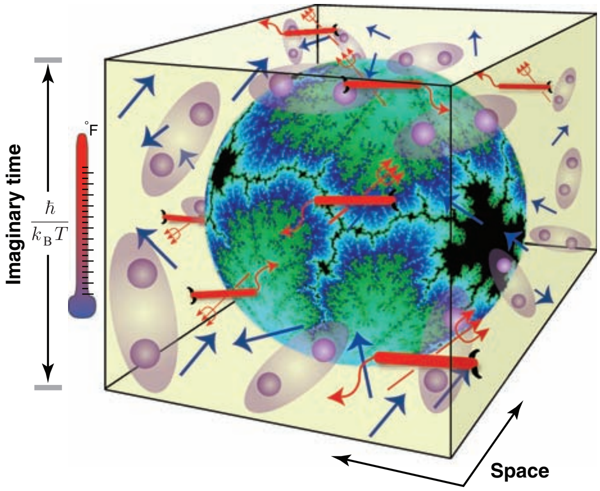

Fig. 1. Illustrating the Feynman path integral, the mathematical tool of choice to address emergence phenomena in many-particle quantum systems (*2*). Near a quantum phase transition, the world inside space-time turns scale-variant at shorter scales, like the Julia set of this cartoon, whereas at larger scales a stable form of quantum matter takes over. Dealing with fermions, the devilish minus signs obscure, however, any detailed understanding of these space-time worlds. The duration of imaginary time is determined by $\hbar$ (Planck's constant divided by 2π) and the product of Boltzmann's constant k_B and absolute temperature T.

These quantum critical states have in fact become quite ubiquitous in the laboratory. One variety is formed by the "designer quantum critical states," where the theorists have so much understanding that they can now guide the experimentalists to where to look. Prominent examples are the cold atoms and the spin systems, as highlighted by Bloch (*8*) and Lloyd (*9*). For bosons and some spin systems, the stuff filling up the space-time of the path integral is similar to classical matter, but this is not at all the case for fermions. The culprit for this deviation is the infamous "fermion sign problem." When fermions come into play, it turns out that in the statistics underlying the description of matter, one must deal with "negative probabilities," and this detaches the many-fermion problem from any classical analog. We have in fact no understanding at all of what is going on in space-time, because we need mathematics to look around and the sign problem is "NP hard" (*10*), meaning that the problem is mathematically unsolvable. The only fermionic substance that we can handle mathematically is the Fermi liquid, the state of electrons in normal metals. Although different from any form of classical matter, this state is at first sight deceptively simple: The electrons turn cooperatively into noninteracting "quasi-electrons" that only communicate via the Pauli exclusion principle. This pushes them into states of high quantum zero-point energy (*7*), and their average "Fermi energy" is typically on the order of 10,000 K in standard metals. This Fermi liquid is in fact a purely empirical construction. Although it is observed in experiments, theoretical physics has failed to explain its existence in general terms, despite countless attempts.

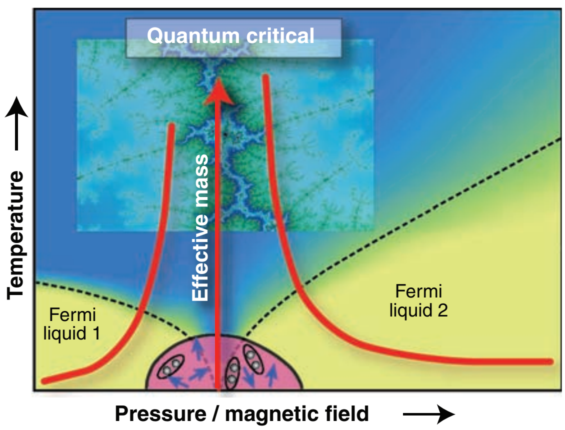

Fig. 2. Typical phase diagram observed in the heavy-fermion metals in the proximity of a quantum phase transition (*3*–*6*). The thermal phase transition to a magnetic state is driven to zero temperature by varying a magnetic field or pressure, and this is the anchor point of a regime of finite-temperature quantum critical fluid behavior fanning out for increasing temperature. The fermionic weirdness manifests itself through the effective mass of the quasi-electrons in the Fermi liquids on both sides, which increases without bound approaching the quantum phase transition. Invariably one finds that at a low temperature, an exotic superconductor (or even a quantum liquid crystal state) takes over at the last minute.

But electrons are also known to form "non–Fermi liquid" states, as can be found in metals containing rare earth, actinide, and transition metal ions (*3*–*6*). At ambient conditions, these show a phase transition to some magnetic state at a low temperature, and by applying pressure or magnetic fields, this transition can be driven to zero temperature. For example, in the diagram of pressure/magnetic field versus temperature (Fig. 2), one finds a V-shaped region anchored where the magnetic transition approaches zero temperature, and the regime inside the V shows the telltale signs of the quantum critical fluid. Actually, this fanning out of the quantum critical region for increasing temperature is just mapping out the "scale invariance geography" in space-time. When the system is close but not at the phase transition, it will show the physics of the stable phase at large scales (Fig. 1). However, upon zooming in, the system will forget its preferred state, and at a characteristic scale the system will reenter the scale-invariant regime. The increase of temperature is like the magnification factor of a microscope, and the V reveals that the scale where the system takes the decision to become a stable phase shifts to shorter times when one moves away from the quantum critical point.

But now the fermion signs hit hard: The experiments give away the workings of quantum scale invariance in space-time, but we have no clue whatever about the nature of the stuff creating the scale invariance! The stable states that are found outside the V in the proximity of the quantum phase transition are Fermi liquids, and because we have a phenomenological understanding of these states, they tell us something. When interactions are weak, one can do controlled calculations, and these reveal a peculiar Fermi liquid rule: The interactions between real electrons have the effect of increasing the mass of the quasi-electrons. This mass enhancement effect is quite modest when the calculations can be trusted, but in the approach to the metallic quantum phase transition, one finds that the effective mass of these quasi-electrons easily exceeds 1000 times the electron mass, to increase indefinitely upon getting closer and closer to the quantum critical point (*3*, *4*). Again the only hold we have is quantum scale invariance: The Fermi energy is a scale, but because the quantum critical state forbids any scale, it has to disappear. The Fermi energy is the average zero-point motion energy, and the only way to remove it is by making the mass of the quasi-electrons infinite!

Somehow there is something badly wrong with these infinitely heavy quasi-electrons. Nature seems to share this concern: Without exception, one observes that eventually some other stable quantum matter state takes over (Fig. 2). These phenomena are currently under intense investigation, and it is clear that they can be quite strange. Recently a quantum version of a liquid crystal was discovered (*5*) but generically strange forms of superconductivity were found (*3*, *4*), including a superconductor that appears to be indestructible by magnetic fields (*6*).

These observations beg for an explanation in terms of a triumphant mathematical theory, but the efforts of the theorists have gotten stuck in running variations on the established themes of bosonic matter and the Fermi liquid (*4*): One finds the fermion signs, in one or the other disguise, as the proverbial brick wall blocking any progress. The "heavy fermion" quantum criticality highlighted here is quite instrumental in forcing us to face the fact that there is still a vast quantum territory lying behind our intellectual horizon that awaits further exploration. The 20-year-long struggle of the physics community with superconductivity at high temperatures, as found in copper oxides, might well be rooted in the sign problem: Although the empirical situation is less clear, there are indications that this high-T_c superconductivity is born from a quantum critical state (*11*). But the fermion signs infest all of physics. In high-energy physics this is well recognized in the context of quark matter, but it might even be consequential in the most fundamental realms (*12*). The modern way of thinking about the ultimate origin of space-time (and everything else) has quantum emergence as a common denominator, but even string theory rests in this regard on intuitions originating in the earthly realms. There are plenty of fermions in such theories, but they are instinctively taken to be of the Fermi-liquid kind, and there is plenty of room for big surprises caused by the fermion signs at the very bottom.

References and Notes

1. J. Cardy, *Scaling and Renormalization in Statistical Physics* (Cambridge Univ. Press, Cambridge, 1996).
2. S. Sachdev, *Quantum Phase Transitions* (Cambridge Univ. Press, New York, 1999).
3. P. Coleman, A. J. Schofield, *Nature* **433**, 226 (2005).
4. H. von Löhneisen, A. Rosch, M. Vojta, P. Wölfle, *Rev. Mod. Phys.* **79**, 1015 (2007).
5. R. A. Borzi *et al.*, *Science* **315**, 214 (2007); published online 22 November 2006 (10.1126/science.1134796).
6. F. Levy, I. Sheikin, A. Huxley, *Nature Phys.* **3**, 460 (2007).
7. A. J. Leggett, *Science* **319**, 1203 (2008)
8. I. Bloch, *Science* **319**, 1202 (2008).
9. S. Lloyd, *Science* **319**, 1209 (2008).
10. M. Troyer, U. J. Wiese, *Phys. Rev. Lett.* **94**, 170201 (2005).
11. J. Zaanen, *Nature* **430**, 513 (2004) and references therein.
12. H. Georgi, *Phys. Rev. Lett.* **98**, 221601 (2007).
13. I thank F. Krueger for his work on the figures.

10.1126/science.1152443

PERSPECTIVE

Supersolidity

M. H. W. Chan

The observation of nonclassical rotational inertia (NCRI) by the torsional oscillator in 2004 gave rise to a renaissance in the study of solid helium-4. Recent theoretical and experimental studies found evidence that disorder in the solid plays a key role in enabling superfluidity. A recent experiment found a marked increase in the shear modulus that shares the same temperature and helium-3 impurity concentration dependence as that of NCRI. This correlation indicates that the onset of superfluidity requires the pinning and stiffening of the dislocation network by helium-3.

Shortly after the discovery of superfluidity in liquid ^{4}He (*1*, *2*), the possibility of the same phenomenon occurring in solid helium was raised by Wolfke (*3*). Careful theoretical consideration of the problem (*4*–*7*) suggested that the possible presence of quantum mechanically induced or zero-point lattice vacancies could facilitate such a "supersolid." In this scenario the superfluid fraction, which reaches 100% in liquid helium, may be immeasurably small. Nevertheless, the suggestion spurred considerable experimental effort in search of evidence for the supersolid phase. Other than some interesting anomalies in the ultrasound experiments (*8*), these efforts were unsuccessful (*9*). The situation changed in 2004 when we reported (*10*, *11*) superfluid-like behavior of solid helium samples housed within a torsional oscillator (TO).

In an ideal TO the resonant period is given by $2\pi(I/G)^{1/2}$, where G is the torsional spring constant of the torsion rod and I the rotational inertia of the torsion bob. We observed that below 200 mK, the resonant period of such an oscillator drops rather abruptly. A number of control experiments led us to conclude that the period drop was due to the solid ^{4}He confined inside the torsion bob oscillating with an effective moment of inertia I that is smaller than the high-temperature, classical value. This is known as nonclassical rotational inertia (NCRI), where $I(T) = I_{classical}[1 - f_s(T)]$ and $f_s(T)$ is the superfluid fraction (*7*). f_s becomes distinguishable from noise at an onset temperature, T_0 ~ 200 mK, and grows at first gradually and then more rapidly with decreasing temperature before saturating below ~50 mK (Fig. 1). We found f_s ~ 1% in the low-temperature limit for solid samples grown inside an annulus of 1 mm in width, as well as for those confined within porous structures having characteristic lengths from nanometers (*10*) to half a micrometer (*12*). The measured value of f_s is attenuated when the oscillation speed exceeds a value corresponding to several quanta of circulation, suggesting that the important excitations in the system are vortices. The phenomenon is immensely sensitive to ^{3}He impurities, even down to a concentration of x_3 in the 1 part per billion (ppb) level. The temperature at which f_s reaches half its saturated value, $T_{1/2}$, increases smoothly from 30 mK at x_3 = 1 ppb to 500 mK at x_3 = 85 parts per million (ppm) (*10*, *13*).

A phenomenological model that captures a great deal of the experimental findings is the vortex liquid model proposed by Anderson (*14*). The attenuation of f_s with oscillation speed is attributed to the nonlinear susceptibility of the entangled collection of many thermally activated vortices. The ability of the vortices to move counter to the time-dependent superflow (relative to the cell's oscillation) results in the screening of the supercurrents. As the temperature is lowered, the motion and number of vortices are reduced so that f_s becomes finite. One prediction of the model is an increase in T_0 with increasing measurement frequency, and it was confirmed for the same sample (x_3 = 0.3 ppm) that T_0 ~ 160 mK at 496 Hz and T_0 ~ 240 mK at 1173 Hz (*15*).

The observation of NCRI has now been replicated in at least three other laboratories (*16*–*19*). Although the temperature dependence of $f_s(T)$ is entirely reproducible, its magnitude varies substantially. The low-temperature supersolid fraction ranges from as little as 0.015% to as much as 20%, the latter of which was reported by Rittner and Reppy in their studies of extremely narrow annuli (0.15 mm width) of solid helium (*17*). They also found f_s to be substantially reduced by thermally annealing the sample (*16*). The large variation in f_s, the effects of annealing, and the lack of evidence for zero-point vacancies (*20*) in the T = 0 limit support the theoretical consensus that superfluidity does not exist in a perfect crystal (*20*, *21*).

Three types of disorder have been considered to be responsible for the phenomenon: glassy regions, grain boundaries, and dislocation lines. Glassy regions have been proposed primarily because they lack crystalline order, thus making them more amenable for superfluidity. Indeed, a quantum Monte Carlo simulation (*22*) found that when disorder is quenched into the

Department of Physics, Pennsylvania State University, University Park, PA 16802, USA. E-mail: chan@phys.psu.edu

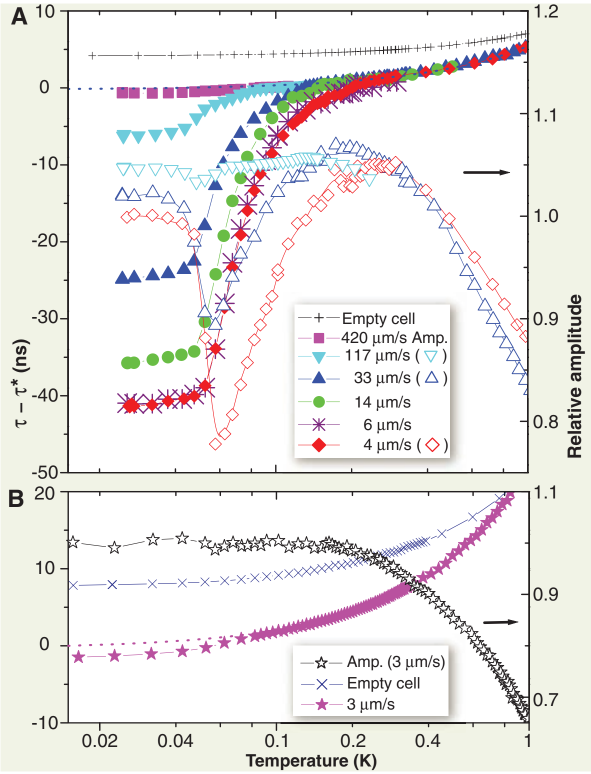

Fig. 1. (**A**) The period shift (left scale, filled symbols) and the relative oscillation amplitude (right scale, open symbols) of the TO for different maximum oscillation speeds as measured by Kim and Chan (*11*). The introduction of solid ^{4}He into the annular open space in the torsion cell increased the resonant period by 3012 ns. (**B**) The period shift was greatly reduced when measurements were carried out in a cell with a barrier inserted in the annulus. [Figure reproduced from (*11*)]

solid phase of ^{4}He it exhibits superfluid characteristics, indicating that NCRI may be the consequence of percolating "superglass" regions. The major difficulty with this idea is that such a glass phase has never been detected in x-ray or other diffraction studies, past or present. Results from a recent high-precision specific-heat study are also inconsistent with glassy behavior (*23*). Instead of the expected linear dependence on T for glasses, a peak with a maximum height of ~20 μJ/mol-K (2.5×10^{-6} k_B per ^{4}He atom, where k_B is Boltzmann's constant) is found. The peak is centered near 75 mK (T_0 of NCRI in 1 ppb ^{4}He) in all solid samples studied. This peak suggests that there is indeed a genuine thermodynamic phase transition separating the normal and the supersolid phases.

In the grain boundary model, it is proposed that liquid superfluid films flow along the interfaces of small crystalline grains and give rise to NCRI (*24*–*26*). This indicates that f_s scales with the total area of the grain boundaries. It is known that samples grown by the blocked capillary method, the method used in all but one TO experiment to date, commonly result in crystal grains with linear dimensions larger than 0.1 mm (*27*). In contrast, solid helium confined in porous Vycor glass (having a typical pore size of 7 nm) has a surface area per unit volume that is roughly 10^4 times as large, and yet f_s is on the same order (~1%) as that of many bulk solid samples.

Dislocation lines in solid ^{4}He form a three-dimensional network consisting of a vast number of dislocation segments and nodes, the latter of which are essentially immobile. Ultrasound measurements indicated that when an oscillating stress field is imposed, the dislocation segments vibrate with little or no damping below 1 K (*28*). The network is characterized by the total dislocation line length per unit volume, λ, and network loop length, L_N, between nodes. In single crystals it was found that $0.1 < \lambda^* L_N^2 < 0.3$, where $\lambda \sim 1 \times 10^6$ cm^{-2} and $L \sim 5$ μm. The dislocation lines can also be pinned by ^{3}He impurities that condense onto them (*29*). The average distance, L_3, between the ^{3}He atoms on a dislocation line is determined by the binding energy, $E_b/k_B \sim 0.5$ K (*13*, *29*, *30*), and the temperature. At a fixed x_3, there is a specific crossover from network pinning to impurity pinning when L_3 becomes shorter than L_N. The characteristic temperatures, such as T_0 and $T_{1/2}$, of samples of different x_3 are found to track the crossover temperature when impurity pinning dominates (*13*), implying that the appearance of NCRI is related to the stiffening of the dislocation network. A direct measurement of the shear modulus has recently confirmed this interpretation. Day and Beamish (*31*) found a marked increase (between 5 and 20%) of the shear modulus, μ, of solid helium (with $x_3 = 0.3$ ppm) below 250 mK. The temperature dependence of μ resembles that of f_s found in TO measurements. When the measurements were repeated with just 1 ppb of ^{3}He impurities, the increase in μ shifted to a lower temperature, consistent with the TO results (Fig. 2).

The stiffening of the dislocation network and the onset of NCRI are clearly related. However, there is as yet no understanding of how these phenomena are correlated. It has been suggested that the long-range phase coherence inherent in supersolidity requires a rigid dislocation network that is pinned by ^{3}He impurities (*13*, *31*). It has also been suggested that superflow takes place along the dislocation lines (*32*, *33*) and that supersolidity appears when these dislocation lines are cross-linked into a three-dimensional network. The problem with this latest idea is that typical densities of dislocations are three orders of magnitude too low to support $f_s \sim 1\%$, let alone 20%.

There is also the possibility of a more mundane connection between the two phenomena without invoking supersolidity. In a real TO, the resonant period depends on the exact dimensions, densities, and elastic moduli of all its constituent parts (*34*). The stiffening of solid helium inside the torsion cell will lead to an enhancement of the overall rigidity of the system and therefore lower the resonant period, mimicking mass decoupling. A careful simulation study by means of the finite element method of the annular TO used by us (*11*) indicates that the reduction in

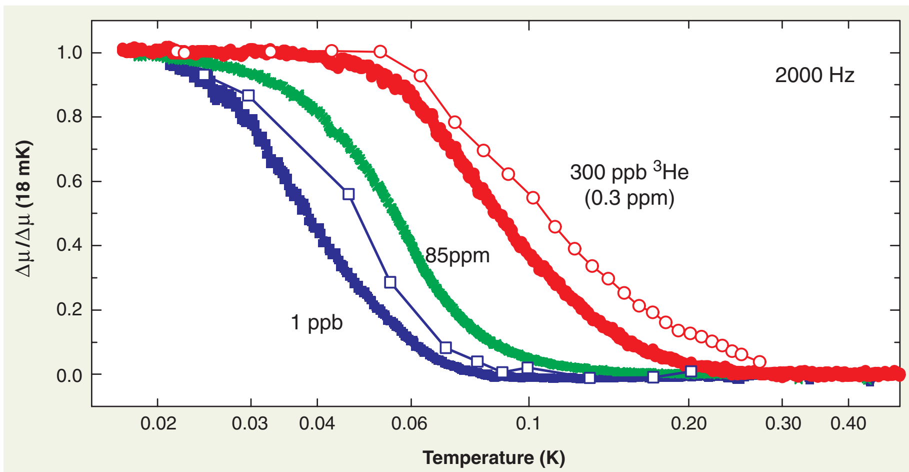

Fig. 2. Shear modulus anomaly in solid ^{4}He with 1 ppb, 0.3 ppm, and 85 ppm of ^{3}He impurities as measured by Day and Beamish (*31*). Changes in shear modulus, Δμ, have been scaled by the values at 18 mK in order to compare temperature dependence. Open circles with lines are similarly scaled NCRI data from TO measurements on 1 ppb (*27*) and 0.3 ppm (*11*) samples.

the resonant period (~10^6 ns) due to a 10% increase in μ of solid helium is less than 0.5 ppm or 0.5 ns (*35*). This decrease is a factor of 100 less than the period drop observed experimentally (Fig. 1). In addition, it is difficult to correlate dislocation stiffening to NCRI for solid helium confined in porous gold (*12*) and particularly in Vycor glass (*10*), because the dimensions of solid helium are much smaller than the micrometer-sized dislocation segments.

An important test that should clarify the relation and the possible causality of the two phenomena would be repeating both the shear modulus and TO experiments with the same sample of ultrahigh-purity ^{3}He. Ultrasound measurements indicate that the dislocation network in ultrahigh-purity ^{3}He, particularly in the high-pressure, hexagonal-close-packed (hcp) phase ($P > 10^7$ Pa), responds to isotopic impurities in much the same way as does hcp ^{4}He (*36*), the solid helium phase of current interest. A likely outcome would be the appearance of a similar increase in the shear modulus without any (or a greatly reduced) concomitant drop in the period of the TO because ^{3}He is a fermion. However, given the history of the last 3.5 years, there may well be other unforeseen surprises.

References and Notes

1. P. Kapitza, *Nature* **141**, 74 (1938).
2. J. F. Allen, A. D. Misener, *Nature* **141**, 75 (1938).
3. M. Wolfke, *Ann. Acad. Sci. Techn. Varsovie* **6**, 14 (1939).
4. L. Reatto, *Phys. Rev.* **183**, 334 (1969).
5. A. F. Andreev, I. M. Liftshitz, *Sov. Phys. JETP* **29**, 1107 (1969).
6. G. V. Chester, *Phys. Rev. A* **2**, 256 (1970).
7. A. J. Leggett, *Phys. Rev. Lett.* **25**, 1543 (1970).
8. J. M. Goodkind, *Phys. Rev. Lett.* **89**, 095301 (2002).
9. M. W. Meisel, *Physica B (Amsterdam)* **178**, 121 (1992).
10. E. Kim, M. H. W. Chan, *Nature* **427**, 225 (2004).
11. E. Kim, M. H. W. Chan, *Science* **305**, 1941 (2004).
12. E. Kim, M. H. W. Chan, *J. Low Temp. Phys.* **138**, 859 (2005).
13. E. Kim *et al.*, *Phys. Rev. Lett.* **100**, 065301 (2008).
14. P. W. Anderson, *Nature Phys.* **3**, 160 (2007).
15. Y. Aoki, J. C. Graves, H. Kojima, *Phys. Rev. Lett.* **99**, 015301 (2007).
16. A. S. C. Rittner, J. D. Reppy, *Phys. Rev. Lett.* **97**, 165301 (2006).
17. A. S. C. Rittner, J. D. Reppy, *Phys. Rev. Lett.* **98**, 175302 (2007).
18. A. Penzev, Y. Yasuta, M. Kubota, *J. Low Temp. Phys.* **148**, 677 (2007).
19. M. Kondo, S. Takada, Y. Shibayama, K. Shirahama, *J. Low Temp. Phys.* **148**, 695 (2007).
20. B. K. Clark, D. M. Ceperley, *Phys. Rev. Lett.* **96**, 105302 (2006).
21. See, for example, the review in (*37*).
22. M. Boninsegni, N. V. Prokof'ev, B. V. Svistunov, *Phys. Rev. Lett.* **96**, 105301 (2006).
23. X. Lin, A. C. Clark, M. H. W. Chan, *Nature* **449**, 1025 (2007).
24. E. Burovski, E. Kozik, A. Kuklov, N. V. Prokof'ev, B. V. Svistunov, *Phys. Rev. Lett.* **94**, 165301 (2005).
25. S. Sasaki, R. Ishiguro, F. Caupin, H. J. Maris, S. Balibar, *Science* **313**, 1098 (2006).
26. L. Pollet *et al.*, *Phys. Rev. Lett.* **98**, 135301 (2007).
27. A. C. Clark, J. T. West, M. H. W. Chan, *Phys. Rev. Lett.* **99**, 135302 (2007).
28. R. Wanner, I. Iwasa, S. Wales, *Solid State Commun.* **18**, 853 (1976).
29. I. Iwasa, K. Araki, J. Suzuki, *J. Phys. Soc. Jpn.* **46**, 1119 (1979).
30. M. A. Paalanen, D. J. Bishop, H. W. Dail. *Phys. Rev. Lett.* **46**, 664 (1981).
31. J. Day, J. R. Beamish, *Nature* **450**, 853 (2007).
32. S. I. Shevchenko, *Sov. J. Low Temp. Phys.* **13**, 61 (1987).
33. M. Boninsegni *et al.*, *Phys. Rev. Lett.* **99**, 035301 (2007).
34. Z. Nussinov, A. V. Balatsky, M. J. Graf, S. A. Trugman, *Phys. Rev. B* **76**, 014530 (2007).
35. A. C. Clark, J. D. Maynard, M. H. W. Chan, http://arxiv.org/abs/0711.3619v2 (2008).
36. J. R. Beamish, J. P. Franck, *Phys. Rev. B* **28**, 1419 (1983).
37. N. V. Prokof'ev, *Adv. Phys.* **56**, 381 (2007).

38. I acknowledge the NSF for support and J. R. Beamish, J. D. Maynard, J. D. Reppy, E. Kim, A. C. Clark, X. Lin, and J. T. West for informative discussions. I thank J. R. Beamish for providing Fig. 2.

10.1126/science.1155302

PERSPECTIVE

Quantum Information Matters

Seth Lloyd

This Perspective discusses the role that quantum information plays in determining the quantum-mechanical aspects of matter. Beginning with the entwined concepts of information and entropy, the article discusses how quantum information theory can supply us with novel concepts and techniques for understanding how matter behaves at the most microscopic of levels.

At first glance, the relationship between quantum information and quantum matter seems tenuous. Information is not very material: It is more concept than thing. Quantum information is even more ethereal than classical information. Matter, by contrast, is solid stuff, reliable and down to earth. The word for matter comes from the Latin *materia*: "wood for building, construction materials." Quantum matter is particularly solid: Quantum mechanics guarantees the stability of the elementary particles and atoms that make up the building blocks of nature. (A hydrogen atom constructed according to the laws of classical electromagnetism would explode in a burst of radiation in less than a trillionth of a second. A hydrogen atom constructed according to the laws of quantum mechanics can last the age of the universe.) Matter seems to be about energy and stability; when it comes to discussing its properties, why should quantum information matter?

Department of Mechanical Engineering, Massachusetts Institute of Technology, MIT 3-160, Cambridge, MA 02139, USA. E-mail: slloyd@mit.edu

In fact, when it comes to matter, quantum information matters a lot. First of all, information is not as immaterial as it might seem. By the end of the 19th century, the great statistical mechanicians Maxwell, Boltzmann, and Gibbs had firmly established that the physical quantity called entropy, which limits the efficiency of heat engines, was in fact a form of information—information about the microscopic motions of atoms and molecules. The very first paper about quantum matter, Planck's 1901 paper on black-body radiation, was also fundamentally about information and quantum mechanics (*1*). In that paper, Planck not only introduced his famous constant to establish the relationship between energy and frequency ($E = h\nu$), he also established the constant of proportionality between information (defined statistically) and entropy. This constant, now called Boltzmann's constant, $k_B = 1.3806503 \times 10^{-23}$ J/K, can be thought of as establishing the relationship between information and entropy: One bit corresponds to an amount of entropy equal to k_B times the logarithm of 2. Planck's paper established that the universe was, at bottom, digital.

Quantum information theory studies the consequences of the digital nature of the universe. Quantum computers are devices that store information at the level of individual quanta (photons, electrons, atoms, etc.) and process that information in a way that preserves quantum coherence (*2*). Quantum communication systems transmit information at the ultimate rates allowed by the laws of quantum mechanics. Although, as noted above, information has played an important role in quantum mechanics since the very beginning, quantum information theory as a distinct discipline is a young field. Before Shor's 1994 discovery that quantum computers could in principle factor large numbers and so break commonly used codes (*2*), quantum information mattered to only a handful of scientists.

During the 1960s and 1970s, Richard Feynman was involved in attempts to use classical digital computers to evaluate the consequences of quantum field theory. He observed that quantum mechanics was hard to program on a classical digital computer. The reason for this difficulty was straightforward: Quantum mechanics possesses a variety of strange and counterintuitive features, and features that are hard for human beings to comprehend are also hard for classical computers to represent at the level of individual classical bits. Consider that a relatively small quantum system consisting of a collection of 300 electron spins "lives" in $2^{300} \approx 10^{90}$ dimensional space. As a result, merely writing down the quantum state of the spins in a classical form as a sequence of bits would require a computer the size of the universe, and to compute the evolution of that state in time would require a computer much larger than that.

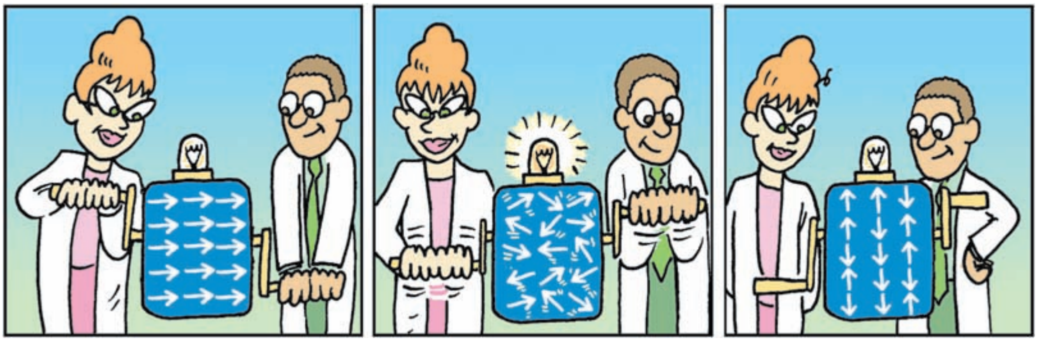

Fig. 1. This figure shows, in an impressionistic fashion, the state of a spin system performing an adiabatic quantum computation. The first panel shows the initial state of the spins. They are all aligned along the *x* axis by a strong magnetic field. The magnetic field is gradually turned off, and, at the same time, a "problem" Hamiltonian whose ground state encodes the answer to some hard problem is gradually turned on. In the center panel, both the magnetic field and the problem Hamiltonian are of equal strength, and the spins are in a fully entangled state. In the final panel, the magnetic field has gone to zero and the spins are in the ground state of the problem Hamiltonian. The answer to the problem can now be obtained by reading off the states of the spins with the convention that spin up = 0 and spin down = 1.

In 1982, Feynman noted that if one has access to a quantum-mechanical device for representing the state of the spins and for transforming that state, rather than a classical device, then the computation of the time evolution of such a system can be much more economical (*3*). Consider a collection of 300 two-level quantum systems, or qubits, one for each electron spin. Suppose that one sets up or programs the interactions between those qubits to mimic the dynamics of the collection of spins. The resulting device, which Feynman called a universal quantum simulator, will then behave as a quantum analog computer, whose dynamics form an analog of the spin dynamics. Since Feynman's proposal, researchers in quantum information have created detailed protocols for programming quantum analog computers, including reproducing the behavior of fermions (*4*, *5*) and gauge fields. Large-scale quantum simulators have actually been constructed out of crystals of calcium fluoride (*6*). Each crystal contains a billion billion spins, which can be programmed using techniques of nuclear magnetic resonance to simulate the behavior of a wide variety of solid-state systems. These solid-state quantum simulators have already revealed a variety of previously unknown quantum phenomena, including spin transport rates that are startlingly higher than the rates predicted by semiclassical theory.

The chameleon-like ability of quantum computers to change their behavior not only allows them to simulate other quantum systems but also gives rise to novel methods for solving computational problems. Many classically hard problems take the form of optimization problems—for example, the traveling salesman problem, in which one aims to find the shortest route connecting a set of cities. Such optimization problems can be mapped onto a physical system, in which the function to be optimized is mapped onto the energy function of the system. The ground state of the physical system then represents a solution to the optimization problem. A common classical technique for solving such problems is simulated annealing: One simulates the process of gradually cooling the system in order to find its ground state (*7*). However, simulated annealing is bedeviled by the problem of local minima, states of the system that are close to the optimal states in terms of energy but very far away in terms of the particular configuration of the degrees of freedom of the state. To avoid getting stuck in such local minima, one must slow the cooling process to a glacial pace in order to ensure that the true ground state is reached in the end.

Quantum computing provides a method for getting around the problem of local minima. Rather than trying to reach the ground state of the system by cooling, one uses a purely quantum-mechanical technique for finding the state (*8*). One starts the system with a Hamiltonian, or energy functional, whose ground state is simple to prepare (for example, "all spins are sideways"). Then one gradually deforms the Hamiltonian from the simple dynamics to the more complex dynamics, whose ground state encodes the answer to the problem in question. If the deformation is sufficiently gradual, then the adiabatic theorem of quantum mechanics guarantees that the system remains in its ground state throughout the deformation process. When the adiabatic deformation is complete, then the state of the system can be measured to reveal the answer.

Adiabatic quantum computation (also called quantum annealing) represents a purely quantum way to find the answer to hard problems (Fig. 1). How powerful is adiabatic quantum computation? The answer is nobody knows for sure. The key question is what is sufficently gradual deformation? That is, how slow does the deformation have to be to guarantee that the transformation is adiabatic? The answer to this question lies deep in the heart of quantum matter. As one performs the transformation from simple to complex dynamics, the adiabatic quantum computer goes through a quantum phase transition. The maximum speed at which the computation can be performed is gov-

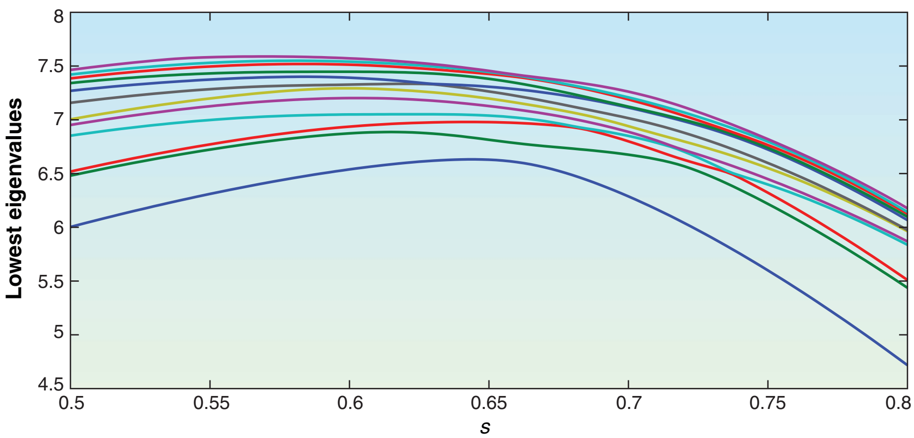

Fig. 2. The figure shows the 12 lowest energy levels of a 16-qubit quantum computer undergoing an adiabatic quantum computation. Over the course of the computation, the system's Hamiltonian goes from a simple form to a complex form whose ground state encodes an instance of a hard computational problem. The parameter *s* tells how simple or complex the Hamiltonian is. For s = 0.05, the Hamiltonian is relatively simple. At the point s = 0.8, the Hamiltonian is complex and its ground state encodes the answer to the problem. At the intermediate point s = 0.66, the Hamiltonian undergoes a phase transition between its simple and complex forms. At this point, the gap between the ground-state energy and the first excited-state energy reaches its minimum value and the state of the system is maximally entangled. The time it takes to perform the computation is inversely proportional to the size of this gap. Just how the gap scales with the size of the system is an open question in both quantum computation and the theory of quantum phase transitions.

erned by the size of the minimum energy gap of this quantum phase transition (Fig. 2). The smaller the gap, the slower the computation must be. The scaling of gaps during phase transitions (gapology) is one of the key disciplines in the study of quantum matter (*9*). Although scaling of the gap has been calculated for many familiar quantum systems, such as Ising spin glasses, calculating the gap for adiabatic quantum computers that are solving hard optimization problems seems to be just about as hard as solving the problem itself.

Few quantum computer scientists believe that adiabatic quantum computation can solve the traveling salesman problem. Nonetheless, there is good reason to believe that adiabatic quantum computation can outperform simulated annealing on a wide variety of hard optimization problems.

How do you build an adiabatic quantum computer? You need a quantum system that can be programmed to enact different Hamiltonians whose ground states instantiate the solutions to hard problems. The dynamics of the system must be sufficiently flexible to be adiabatically deformed from simple to complex Hamiltonians. Finally, you have to be able to measure the final state of the system to read out the answer to the problem.

Quantum systems found in nature are typically too inflexible or too uncontrollable to meet the requirements for adiabatic quantum computation. Human-made quantum systems, in contrast, can be designed to meet those specifications. The solution arises from one of quantum information processing's most significant contributions to the study of quantum matter: the demonstration of macroscopic quantum coherence. Macroscopic quantum coherence arises when a relatively macroscopic system, such as a supercurrent containing billions of electrons, is able to exhibit collective quantum behavior. In 2001, for example, two groups were able to put superconducting loops in a macroscopic quantum superposition of supercurrent flowing clockwise and counterclockwise simultaneously (*10*, *11*). (Don't try to visualize this phenomenon; it is one of those quantum effects mentioned above that resist all intuitive explanation.) In 2003, my colleagues and I designed superconducting circuits that used macroscopic quantum coherence to implement adiabatic quantum computation (*12*). Although 16-qubit and 28-qubit devices based on this design have been constructed, claims that adiabatic quantum computers can solve all sorts of hard problems in a completely quantum-mechanical fashion have yet to be borne out (*13*, *14*). Even if adiabatic quantum computers fail to solve hard problems, such devices still constitute artificial systems that, as Feynman envisaged, can simulate the behavior of strange computational quantum matter.

Quantum information offers a wide variety of techniques for understanding quantum matter. One of the primary contributions of the field is a detailed picture of the weird form of quantum correlation known as entanglement (*2*). Entanglement initially made its way into quantum mechanics as a particularly egregious example of a counterintuitive quantum phenomenon. (Einstein referred to it as "spukhafte Fernwirkung," or "spooky action at a distance.") Quantum information theory has shown that, far from being exotic, entanglement is ubiquitous. Entanglement underlies the stability of the covalent bond; entanglement is a key feature of ground states of solid-state systems (*15*); even the vacuum of space is entangled (*16*)! In fact, the Hawking radiation emerging from black holes can be thought of as a particularly exotic form of vacuum entanglement (*17*).

Over the past few decades, quantum information theory has transcended its earlier role as an esoteric study of the foundations of quantum mechanics, to become an integral part of the science of quantum matter in all its forms. As its relationship with the theory of quantum matter becomes more elaborate and more intimate, quantum information theory has the potential to unravel some of the deep mysteries of physics. For example, the connection of entanglement with Hawking radiation suggests that quantum information may have a key role to play in understanding quantum gravity. Meanwhile, quantum information theory has suggested a wide variety of fundamental experiments in quantum matter, such as the demonstration of macroscopic quantum coherence mentioned above (*10*, *11*). Unexpected connections have arisen between transport theory in quantum matter and quantum information theory: Hopping electrons might be able to discern winning strategies for chess or Go more efficiently than classical chess or Go masters (*18*), or the efficiency of photosynthesis might arise because excitons moving through photocenters effectively implement a quantum computer algorithm (*19*). Finally, if we can continue to realize Feynman's vision of quantum simulation, we may someday be able to simulate the behavior not only of electrons and elementary particles but of the universe itself (*20*). The fact that simulating the entire universe requires a quantum computer as big as the universe should not stand in our way!

References

1. M. Planck, *Ann. Phys.* **4**, 553 (1901).
2. M. A. Nielsen, I. L. Chuang, *Quantum Computation and Quantum Information* (Cambridge Univ. Press, Cambridge, 2000).
3. R. P. Feynman, *Int. J. Theor. Phys.* **21**, 467 (1982).
4. S. Lloyd, *Science* **273**, 1073 (1996).
5. S. Bravyi, A. Kitaev, *Ann. Phys.* **298**, 210 (2002).
6. C. Ramanathan, S. Sinha, J. Baugh, T. F. Havel, D. G. Cory, *Phys. Rev. A* **71**, 020303(R) (2005).
7. S. Kirkpatrick, C. D. Gelatt, M. P. Vecchi, *Science* **220**, 671 (1983).
8. E. Farhi, J. Goldstone, S. Gutmann, *Science* **292**, 472 (2001).
9. S. Sachdev, *Quantum Phase Transitions* (Cambridge Univ. Press, Cambridge, 1999).
10. C. H. van der Wal *et al.*, *Science* **290**, 773 (2000).
11. R. Friedman, V. Patel, W. Chen, S. K. Tolpygo, J. E. Lukens, *Nature* **406**, 43 (2000).
12. W. M. Kaminsky, S. Lloyd, in *Quantum Computing and Quantum Bits in Mesoscopic Systems*, B. Ruggiero, P. Silvestrini, A. Leggett, Eds. (Kluwer Academic, Netherlands, 2003), pp. 229–236.
13. P. Judge, "D-Wave's quantum computer ready for latest demo," CNET News.com, 7 November 2007 (www.news.com/D-Waves-quantum-computer-ready-for-latest-demo/2100-1010_3-6217842.html).
14. S. K. Moore, "Prototype commercial quantum computer demo'ed," *IEEE Spectrum Online*, 13 February 2007 (www.spectrum.ieee.org/feb07/comments/1710).
15. M. M. Wolf, G. Ortiz, F. Verstraete, J. I. Cirac, *Phys. Rev. Lett.* **97**, 110403 (2006).
16. M. Cramer, J. Eisert, M. B. Plenio, J. Dreissig, *Phys. Rev. A* **73**, 012309 (2006).
17. S. W. Hawking, *Nature* **248**, 30 (1974).
18. E. Farhi, J. Goldstone, S. Gutmann, preprint available at http://xxx.lanl.gov/abs/quant-ph/0702144.
19. G. S. Engel *et al.*, *Nature* **446**, 782 (2007).
20. S. Lloyd, *Phys. Rev. Lett.* **88**, 237901 (2002).

10.1126/science.1154732

PERSPECTIVE

Looking to the Future of Quantum Optics

Ian A. Walmsley

Light has provided both fundamental phenomenology and enabling technology for scientific discovery for many years, and today it continues to play a central role in fundamental explorations and innovative applications. The ability to manipulate light beams and pulses with the quantum degrees of freedom of optical radiation will add to those advances. The future of quantum optics, which encompasses both the generation and manipulation of nonclassical radiation, as well as its interaction with matter, lies in the rich variety of quantum states that is now becoming feasible to prepare, together with the numerous applications in sensing, imaging, metrology, communications, and information processing that such states enable.

The main attributes of quantum light that distinguish it from classical light are the nature of its fluctuations and of its correlations. These distinguishing features have been known for many years, and precise formulations of what constitutes nonclassical light were developed in the 1960s. However, the past decade has seen marked advances in our ability to prepare specific quantum states of light that elucidate these distinctions and make use of them. Developing this ability further—to increase the distance over which quantum correlations (or entanglement) may be distributed, to increase the number of light quanta (photons) that occupy these entangled states, and to understand how to characterize, store, and exploit those states—is the direction of future research.

Classical light has a certain degree of randomness to it, even though the intensity and phase of a classical light beam or pulse may both be very precisely defined. This is not so for a quantum light beam, the photon number and phase of which cannot both be specified simultaneously. One consequence is that the number of photons in a classical light pulse fluctuates, whereas, for a quantum pulse, it may not. Similarly, the amplitudes of a quantum light beam may be less noisy than those of its classical counterpart.

Classical light is also restricted in the types of correlations that separate beams may have. For example, two light beams may be highly correlated in their wavelength or direction, though each beam individually may not have a well-defined wavelength or direction, respectively. Similarly, the beams may be correlated in the conjugate variables of time or position, or in two orthogonal polarizations. The difference for quantum light is that there may be strong correlations in both pairs of conjugate degrees of freedom at the same time. That is, two light beams in a pure quantum state may be highly correlated in both their direction and their position or in both wavelength and time. This is an example of quantum entanglement.

There are important consequences stemming from the fact that only the joint properties of the two (or more) beams can be said to have definite values. A well-known example of this is in secure communications, for which entanglement is a key property that allows the presence of an eavesdropper to be detected (*1*, *2*). Entanglement among many photons in many different modes is now becoming possible, and optics provides a means to explore the rich structure and dynamics of these complex quantum states.

Today, it is possible to create a broad range of quantum states of light that have no classical analog. There are two main routes to achieve this. The first is to use single emitters, such as individual atoms or excitons (an electron-hole pair), to emit single photons (*3*–*5*) (Fig. 1, A and B). It is possible to use this approach to entangle two emitter atoms by detecting one or more photons from each, provided the apparatus is arranged so that it is impossible to tell from which atom the detected photon has come (*6*–*8*). It will be possible to "daisy chain" this procedure: thus, to arrange that large numbers of atoms, each of which are remote from the others, eventually become entangled, which leads to a highly nonclassical state of many atoms that is mediated by photons.

The integration of quantum optics with technologies such as ion and atom traps, and indeed photonic microstructured cavities and waveguides, provides a deterministic way to swap entanglement between light and matter. One important avenue will be the implementation of quantum memories for photons, which can store information coded in these "flying qubits" in long-lived matter states or "stationary qubits" (*9*–*12*). This is a key technology for quantum communications and computational networks (Fig. 2). Reading out quantum correlations from matter provides a creative approach to sensing and measurement: It may be used to probe the dynamics of complex entangled states of matter with minimal back-action on the atoms. A recent example of this shows how the extraordinarily strong coupling between individual atoms injected into an optical cavity and the microwave field inside the cavity may be used to implement a true quantum nondemolition detector for photons, one that does not absorb the photon in order to measure it (*13*).

A different avenue is that of quantum optomechanics. Recent work has shown that it is possible to cool mirrors in optical cavities with the light force associated with the beams contained in the cavity; effectively, the momentum fluctuations of the mirror are modified by its collisions with the photons (*14*, *15*). In principle, it is possible to produce nonclassical states of the mirror motion in this way; perhaps this will eventually enable us to understand how to entangle a light beam with a truly macroscopic object.

The second route to generating quantum light is by means of nonlinear optics, such as parametric downconversion (Fig. 1, C and D), in which a light beam at one frequency generates two new light beams (the signal and idler), each near half the frequency of the pump. The "splitting" of the pump photon into two siblings

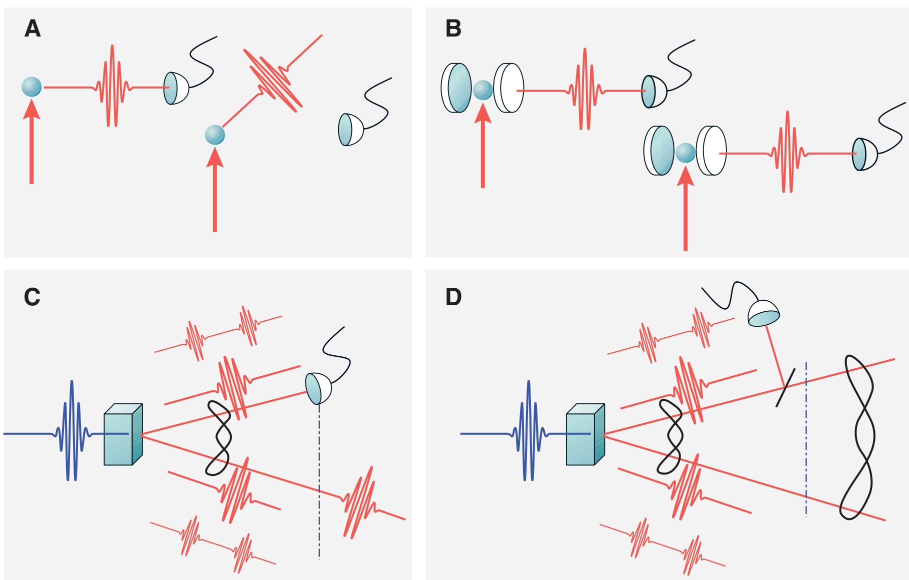

Fig. 1. Approaches to the generation of quantum light. (**A**) A single atom (or single semiconductor quantum dot or other quantum system containing a single electronic excitation), indicated by the sphere, is held in isolation. When it is excited by a laser (red arrow), it emits a single photon into the whole of space. Infrequently, the photon is emitted in the direction of a small photodector (hemispherical bucket). (**B**) A refinement of this approach, which determines the direction of photon emission, places the atom in an optical cavity, consisting of a pair of mirrors (discs). This approach to "shaping" the photons (so that they can be used with conventional optics and detectors) is called cavity quantum electrodynamics and is currently one of the most promising avenues to building multiatom and multiphoton nonclassical states. (**C**) The nonlinear optical process of parametric downconversion, in which a blue photon is split into a pair of red photons in the nonlinear optical microstructure (shaded cube). The signal and idler beams are quantum correlated (or entangled, as indicated by the black lasso) in their in-phase and in-quadrature amplitudes, in their time of generation and frequency, and in their position and direction, providing for the preparation of nonclassical states based on the detection of photons from one of the beams, such as "heralding" a single photon. (**D**) Conditional-state preparation may also increase the entanglement between signal and idler, if (for example) photons are subtracted from one of the beams. This is a component of entanglement distillation and enables Schrödinger cat–like states to be generated.

Department of Physics, University of Oxford, Clarendon Laboratory, Parks Road, Oxford OX1 3PU, UK. E-mail: walmsley@physics.ox.ac.uk

means that the signal and idler are highly entangled in photon number, in their amplitudes, in time-energy, and in position-direction. This provides several different capabilities. For example, it is possible to prepare new kinds of nonclassical states by means of conditional detection, such as one- and two-photon states (*16*), as well as "Schrödinger kitten" states—that is, states whose amplitudes have two distinct values at the same time—named for the famous "cat" that is simultaneously dead and alive (*17*). In terms of applications, fields that are specified by continuous variables such as amplitudes (as opposed to discrete variables such as photon number) have a very large information capacity. Further, even simple optical elements such as a beamsplitter (for example, a half-silvered mirror that combines two beams of light) can be used to entangle and disentangle the signal and idler beams, and photodetectors for amplitude may be very efficient. A goal for the near future is to demonstrate that the approach can be used to combine a large number of weakly entangled beams into a smaller number of highly entangled beams, a process known as entanglement "distillation" (*18*).

Conditional-state preparation based on detection, when combined with feedforward control, also forms the basis of an optical quantum computer (*19*). Here, future measurement sequences are determined by the outcomes of past measurements, enabling information processing with the use of complex entangled states prepared beforehand (*20*) (Fig. 2). Indeed, certain metrology can be brought to a precision of the tightest known quantum bound by this means (*21*). How far this can be pushed remains an open question.

A task for the immediate future is to make quantum states robust against the loss of photons, which is inevitable in propagating light over distance through optical systems (*22*). This would enable the distribution of entanglement over longer distances and address important questions regarding how entanglement scales in the presence of decoherence. Is it possible to find states of many photons or large amplitudes that do not decohere rapidly or that can be protected against noise by experimentally feasible control schemes?

In quantum optics, there is a symbiosis between innovative technology and science. Current exciting developments in optical technologies (such as materials, detectors, and lasers) and structures (such as waveguides, fibers, and photonic crystal microcavities) will play an increasingly important role in enabling the exploration of fundamental quantum phenomena and their applications in new technologies. The future of the field lies in the development of robust approaches for creating clean, pure-state quantum light beams and the precise control of their interaction with matter, both at the level of individual atoms and with macroscopic objects. Pushing this frontier will lead both to scientific discoveries, unraveling the enigmatic character of quantum correlations, and to innovative quantum technologies, many of which are yet to be envisioned.

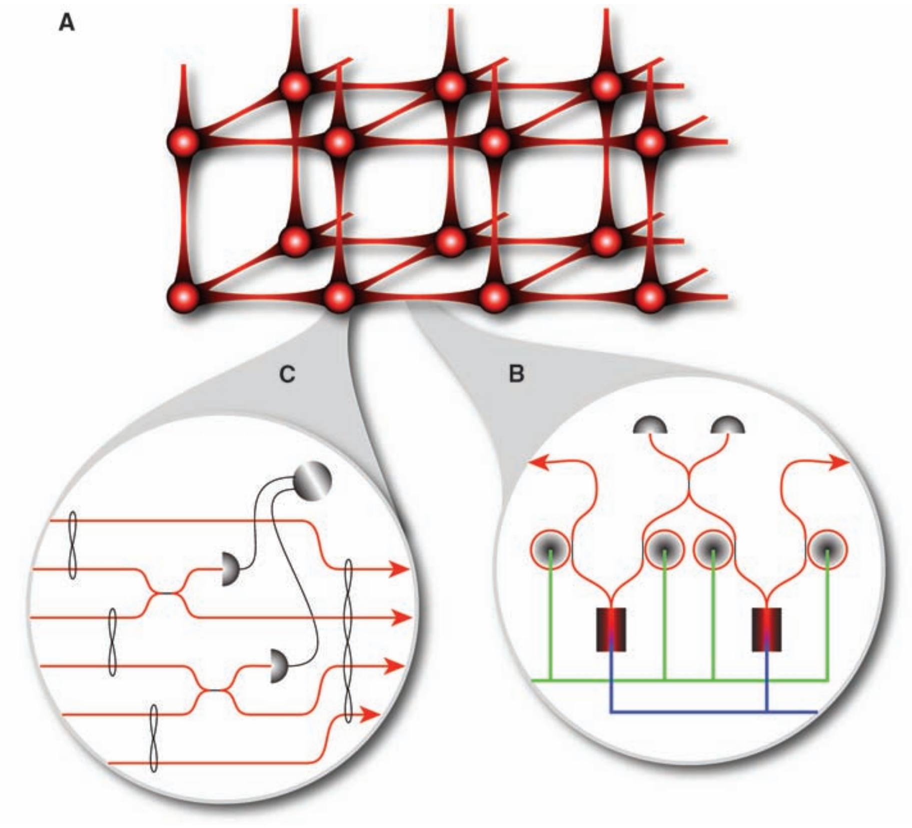

Fig. 2. A quantum network, as shown in (**A**), enables the generation of large-scale entanglement among many photons by means of feedforward. Sources of nonclassical light are distributed between the nodes of the network. The photons' quantum states generated by these sources are stored in quantum memories until required, illustrated in (**B**). The light is released into the network, and separate nodes are entangled by means of measurements on subsets of the photons, shown in (**C**), which herald the production of entangled quantum states among the other nodes. Classical communications are used to signal the network preparation and to control the operations that are then enacted on the output quantum states. This provides a possible scenario for building up many-photon entangled states over large distances. Aside from elucidating fundamental issues such as the robustness of scalability of entanglement, these networks will have applications in quantum communications, computation, and imaging.

References and Notes

1. N. Gisin, G. Ribordy, W. Tittle, H. Zbinden, *Rev. Mod. Phys.* **74**, 145 (2002).
2. A. Acin, N. Gisin, L. Masanes, *Phys. Rev. Lett.* **97**, 120405 (2006).
3. M. Hijlkema *et al.*, *Nat. Phys.* **3**, 253 (2007).
4. J. McKeever *et al.*, *Science* **303**, 1992 (2004).
5. R. M. Stevenson *et al.*, *Nature* **439**, 179 (2006).
6. S. Bose, P. L. Knight, M. B. Plenio, V. Vedral, *Phys. Rev. Lett.* **83**, 5158 (1999).
7. L. M. Duan, M. D. Lukin, J. I. Cirac, P. Zoller, *Nature* **414**, 413 (2001).
8. D. L. Moehring *et al.*, *Nature* **449**, 68 (2007).
9. M. Fleischhauer, M. D. Lukin, *Phys. Rev. A* **65**, 022314 (2002).
10. B. Julsgaard, J. Sherson, J. I. Cirac, J. Fiuráse, E. S. Polzik, *Nature* **432**, 482 (2004).
11. J. Nunn *et al.*, *Phys. Rev. A* **75**, 011401 (2007).
12. A. V. Gorshkov, A. Andre, M. Fleischhauer, A. S. Sorensen, M. D. Lukin, *Phys. Rev. Lett.* **98**, 123601 (2007).
13. S. Gleyzes *et al.*, *Nature* **446**, 297 (2007).
14. D. Vitali *et al.*, *Phys. Rev. Lett.* **98**, 030405 (2007).
15. D. Kleckner, D. Bouwmeester, *Nature* **444**, 75 (2006).
16. D. Achilles, Ch. Silberhorn, I. A. Walmsley, *Phys. Rev. Lett.* **97**, 043602 (2006).
17. A. Ourjoumtsev, R. Tualle-Brouri, J. Laurat, P. Grangier, *Science* **312**, 83 (2006).
18. D. E. Browne, J. Eisert, S. Scheel, M. B. Plenio, *Phys. Rev. A* **67**, 062320 (2003).
19. E. Knill, R. Laflamme, G. Milburn, *Nature* **409**, 46 (2001).
20. R. Prevedel *et al.*, *Nature* **445**, 65 (2007).
21. B. L. Higgins, D. W. Berry, S. D. Bartlett, H. M. Wiseman, G. J. Pryde, *Nature* **450**, 393 (2007).
22. W. Wasilewski, K. Banaszek, *Phys. Rev. A* **75**, 042316 (2007).

23. I acknowledge support from the European Commission through the Integrated Project Qubit Applications (QAP) and from the UK Engineering and Physical Sciences Research Council.

10.1126/science.1152495