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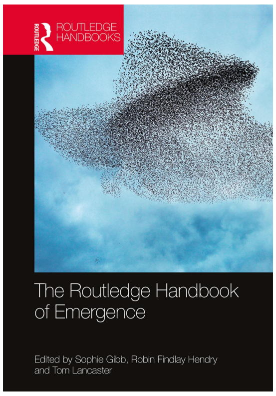

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22

THE EMERGENCE OF EXCITATIONS IN QUANTUM FIELDS

Quasiparticles and topological objects

Tom Lancaster

Introduction

Quantum field theory (QFT) is one of the most successful and rigorously tested areas in all of physics. Central to this discipline is the notion of a *particle* as an excitation in a field. In fact, the central idea of QFT may be stated as: *all particles are excitations in a quantum field* (Lancaster and Blundell 2014). QFT is the only known method of combining quantum mechanics and special relativity. Moreover, as soon as quantum mechanics and relativity are combined, it becomes impossible to deal with single particles, and we are forced to consider the physics of collections of many particles. As a result, QFTs describe not only the relativistic realm of the standard model of particle physics but also the multiparticle realm of condensed matter physics, including the sub-disciplines of (quantum mechanical) hard matter and (the classical statistics of) soft matter. This variety allows a number of insights into cases where candidate emergent phenomena might be found. In this chapter we discuss the emergence of the properties of interacting particles, followed by the emergence of qualitatively new forms of particles. Our aim is to track how *interactions* lead directly to the physical masses and charges of particles in QFT and how they bring about states of matter that support the existence of novel forms of particles.

A particle in classical mechanics is known in Russian, very appropriately, as a *material point* (Landau and Lifshitz 1976). That is, a massive object localized at a point in real space with no extension in space. In conventional quantum mechanics (QM), where position-momentum uncertainty limits our simultaneous knowledge of a particle's position and momentum, it is often more convenient to describe a particle using a basis of states where particles are eigenstates of the momentum operator, rather than the position operator. The quantum particle is then an object defined by its well-defined momentum, with its position in real space being subject to uncertainty. In QFT, in addition to this choice of how to describe particles, their properties are encoded in a small number of parameters known as *coupling constants*, which could include the particle's mass m and charge q. These feature in the *Lagrangian*, which is a mathematical expression giving the difference between kinetic and potential energy densities, and which is built from fields $\phi(x)$. The Lagrangian defines a quantum field theory such that we often speak of "a theory", meaning the state of affairs described by a particular Lagrangian. An example Lagrangian density describing

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scalar fields is $\mathcal{L} = \frac{1}{2}\partial_\mu \phi^2 - \frac{m^2}{2}\phi^2 - \frac{g}{4!}\phi^4$, that is, a fairly simple polynomial expression in the field and its derivatives. Here, the parameter m will turn out to give the mass of the particle excitations, and the parameter g (which is rather like a charge) will tell us how strongly the particles interact.

In practice, QFT involves the process of writing down a Lagrangian, built from classical fields, and then *quantizing* it. Technically, quantization involves setting up commutation relations between the fields in a theory and then expanding the fields in terms of operators that create and destroy particles. The aim of quantization is to express the properties of a system (such as its total energy, or total momentum) in terms of the number of particle excitations present. This procedure is most easily carried out for so-called non-interacting theories. Such theories are frequently mathematically solvable and, as the name suggests, have properties expressible in terms of particles that don't interact with each other. One might think of the particles in a non-interacting theory as passing, ghost-like, through one another, rather than undergoing scattering collisions with one another (as interacting particles do). Of course, real systems involve interactions (we would not be able to measure or experience them in any way if they didn't), and when interactions are included in the description there are frequently two dramatic consequences to our analysis: (i) the theory is often no longer exactly solvable and must be approximated somehow; and (ii) the interactions have a profound effect on the properties of the particles themselves. It is this latter point, and the possibility of emergent properties that it presents, with which this chapter is concerned.

In order to understand the properties of particles in interacting theories, we commonly employ a thought experiment. This involves focusing now on the *Hamiltonian*[1] (rather than the Lagrangian, to which it bears a straightforward mathematical relationship) and splitting up this Hamiltonian into a sum of non-interacting (H_0) and interacting (H') parts: $H = H_0 + \lambda H'$. We then imagine what happens to a non-interacting theory (with the Hamiltonian H_0) when we slowly turn on the interaction Hamiltonian (H'). To do this we simply multiply it by the "turning-on" function λ that starts vanishingly small and slowly grows to unity, at which point the previously non-interacting theory becomes an interacting one that describes a bubbling cauldron of interactions.

When λ is zero the non-interacting system described by the theory is simple. The ground state of the theory is known as a vacuum (mathematically given the symbol $|0\rangle$ to represent a state in which there is nothing at all). The excitations from this ground state are particles $|p\rangle$, labelled by their momentum. These are known as non-interacting, or *bare*, particles and have the masses and charges that may be read off directly from the Lagrangian. In order to create an excitation mathematically, we use the particle creation operator $\hat{a}_{\boldsymbol{p}}^\dagger$ which acts on the vacuum to create a particle with momentum $\boldsymbol{p}$ thus: $\hat{a}_{\boldsymbol{p}}^\dagger |0\rangle = |p\rangle$.

As we turn on the interactions the ground state vacuum changes from $|0\rangle$ into another state we will call $|\Omega\rangle$. If we attempt to excite the interacting system using our operator $\hat{a}_{\boldsymbol{p}}^\dagger$ we now excite particles that may interact with each other, causing more particles to be created or annihilated. The original particle may become lost in the havoc, and it is legitimate to ask whether it is even *possible* to have particles at all in an interacting system.

It will turn out that particles of a kind are still possible in the presence of interactions: they are called *quasiparticles*. These quasiparticles have different properties to their non-interacting relatives. The process of changing the properties of a particle (such as its mass or charge) as we turn on interactions is known as *renormalization*.[2]

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Renormalization in a metal

An influential picture of the example of renormalization of electrons in a metal was formulated by Lev Landau and is known as Landau Fermi Liquid theory (Lifshitz and Pitaevskii 1980; Anderson 1984). It has been called the standard model of condensed matter physics. The quasiparticles in Landau Fermi Liquid theory are slightly different from the field theory quasiparticles described thus far (we will call them Landau quasiparticles to avoid confusion), but the lack of mathematical machinery needed to understand Landau's picture makes it useful for gaining an insight into the subject. In Landau's picture we pay particular attention to the turning on of the interactions.

The non-interacting model of a metal is the *Fermi gas*. This is a state of matter formed by trapping non-interacting electrons in a rigid box. Although the electrons in this model do not interact with each other via the Coulomb force, they must obey Fermi statistics and the Pauli exclusion principle, which ensures that electrons (which are identical Fermi particles) may not doubly occupy momentum states. Electrons therefore stack up in energy, up to the so-called Fermi energy. As we imagine turning on Coulomb interactions, the particles of the Fermi gas change their energy, with the gas evolving into the interacting Landau Fermi liquid. This is known as the *adiabatic turning on* process (Anderson 1984) and provides us with a picture of how particle properties emerge as interactions are turned on.[3]

The vital thing here is that all of the singly occupied states in the gas become singly occupied states in the liquid. We call the object occupying the states in the gas a bare particle and in the liquid that emerges from the gas, a Landau quasiparticle. Despite the fact that the energies change as interactions are turned on, there is no ambiguity about which bare particle has evolved into which Landau quasiparticle, which is to say that the energy levels should not cross during the turning-on process, as shown in Figure 22.1(a). This will be the case since

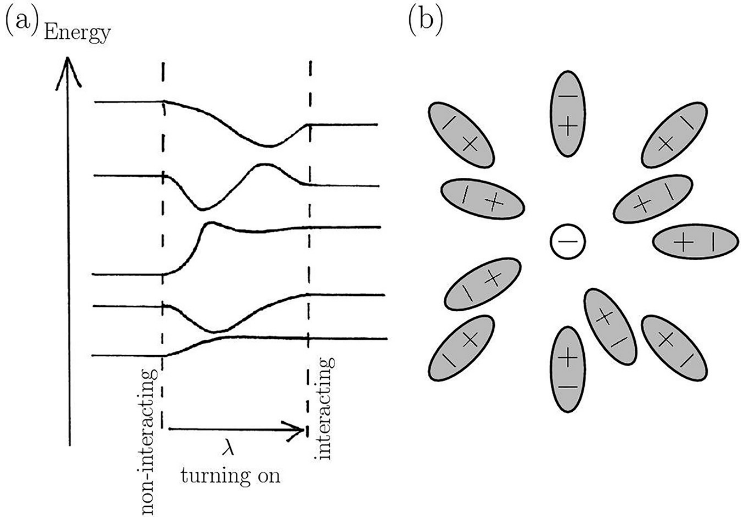

Figure 22.1 (a) The turning-on process in the Landau Fermi liquid evolves non-interacting electrons into Landau quasiparticles. The key feature is that the energy levels change, but never cross, as the interactions are turned on. (b) The change in particle charge on renormalization, viewed as a classical picture of a bare charge being screened by a vacuum made of dipoles.

[Adapted from Lancaster and Blundell 2014]

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the Fermi gas and the Landau Fermi liquid share the same symmetry (i.e. the interactions do not cause a symmetry-breaking phase transition, a topic to which we return later). In the absence of a phase transition the slow turning-on process provides a one-to-one correspondence between Landau quasiparticles and free particles. This state of affairs is known as *adiabatic continuity*.

Landau's quasiparticles are different in some ways from QFT quasiparticles. In QFT, the ground state $|\Omega\rangle$ is the interacting vacuum which contains no quasiparticles. In Landau's theory, the liquid state contains the same number of particles as did the Fermi gas from which it evolves. This guarantees charge conservation, as long as Landau quasiparticles carry the same charge as conventional electrons. However, although Landau quasiparticles have the same charge as bare electrons, the change in energy levels that occurs leads to the particles taking on a new "effective" mass m^*. (That is, we account for the change in energy of the particle that results from the interaction by saying that its mass has changed.)

A valuable insight into the properties of the Landau Fermi liquid is found by looking at the details of the interaction of one electron with all of the others as the interactions are turned on. We therefore imagine introducing an extra electron in an unoccupied state. Interactions have the effect of causing some other electron to be scattered from its momentum state. This latter electron will therefore have a finite *lifetime*. However, the scattering process is strongly constrained by energy-momentum conservation and the Pauli exclusion principle. It may be shown that the lifetime is long (compared, say, to the timescale for turning on the interactions) only for those electrons with energies very near the Fermi energy, since in that case, there are so few available vacant states to accommodate the scattered electrons, they tend not to scatter at all. Thus, it is meaningful to speak about Landau quasiparticles *only* if they have energies near the Fermi energy, as these have a long lifetime against decay. At lower energies, well below the Fermi energy, the status of these excitations is less clear: although the energy levels are filled, the particles are scattering between states very rapidly compared to the slow timescale of the turning-on process and so enjoy only a fleeting existence. Put less technically, only the most energetic electrons qualify as quasiparticles, since the less energetic ones fail to survive for long enough to qualify as particles (where "long enough" is judged against the imagined time taken to turn on the interactions).

The formalism of making quasiparticles in field theory

We give here a brief flavour of how quasiparticles are dealt with mathematically. The process of renormalization may be summed up conceptually as

(Quasiparticle) = (bare particle) + (interactions)

As Philip Anderson points out (Anderson 1963; Anderson 1984), there is nothing guaranteeing the existence of quasiparticles in an interacting system. Returning now to QFT, in order to see where the quasiparticles arise in field theory, we may try to create a particle operationally, by using the creation operator $\hat{a}^\dagger_{\boldsymbol{p}}$ on the interacting vacuum $|\Omega\rangle$. The result is not only the expected particle with momentum $\boldsymbol{p}$ but also lots of other particles with momenta that add up to $\boldsymbol{p}$. That is

$$\hat{a}^\dagger_{\boldsymbol{p}}|\Omega\rangle = \text{(single particle part)} + \text{(multiparticle parts)}$$

It is also possible that we only approximately create the expected single particle and actually make an unstable narrow wavepacket known as a resonance. The consequence is that the particle-like

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resonance takes on a finite lifetime τ. With all of these particles flying around, we measure the likelihood of the occurrence of quasiparticles by defining the *quasiparticle weight* $Z_{\boldsymbol{p}}$ Formally speaking, $Z_{\boldsymbol{p}}$ gives us the overlap of the states of matter we have created and the particle we intended to create. We may say that if $Z_{\boldsymbol{p}} \neq 0$, then we have quasiparticles in the system. This allows us, mathematically, to write the dressing-up process as the passage from a non-interacting particle with wavefunction $e^{-i\,p.x}$ to a quasiparticle with wavefunction $Z_{\boldsymbol{p}} e^{-i\,p.x-t/\tau}$, that is, in becoming a quasiparticle, the bare particle's amplitude is reduced by a factor $Z_{\boldsymbol{p}}$ and takes on a finite lifetime τ.

When interactions between particles are turned on, the most obvious consequence is that particles scatter from each other. This scattering is encoded in QFT using *perturbation theory* and *Feynman diagrams*, which provide a shorthand for calculations (Lancaster and Blundell 2014). As stressed earlier, the interactions also cause renormalization of the particle properties, which can also be encoded mathematically in the following manner. For a particle to propagate in a system, we define the *propagator* G as the quantum amplitude for a particle to be created at point y and be detected at x. One possibility is that a particle might propagate from y to x without being impeded in any way. The amplitude for this is given by the non-interacting propagator G_0. If the particle suffers an interaction at some intermediate position with an amplitude V then the amplitude becomes $G = G_0 + G_0VG_0$. If there is a possibility of a second interaction then the sum becomes $G = G_0 + G_0VG_0 + G_0VG_0VG_0$, and so on. Feynman diagrams allow us to turn this infinite series of interactions into a process of cartooning, allowing the depiction of the nature of each interaction V. Since both the cartoons and the algebra represent an infinite, geometric series, these may be summed to give

$$G = G_0 + G_0VG_0 + G_0VG_0VG_0 + \ldots = \frac{G_0}{1 - VG_0} = \frac{1}{G_0^{-1} - V}$$

For the simplest quantum field theory (the scalar field theory), the non-interacting propagator G_0 may be shown to be given by $G_0 = i/\left(p^2 - m^2\right)$. This implies that the interacting propagator is given by

$$G = i/\left(p^2 - m^2 - V\right).$$

Comparing G and G_0, we see that the scattering process that the particle undergoes has had the effect of shifting its mass from m to $\left(m^2 + V\right)^{1/2}$. It is in this way in QFT that the interactions cause the mass of a particle to be altered.

The means by which interactions cause a change in the charge of a particle may be understood in a similar way to the argument presented earlier. In this case the particle of interest is a photon, which transmits the electrostatic force between charges. Quantum electrodynamics (QED) is the QFT that describes the interaction between electrons and photons. QED allows a photon to instantaneously turn into an electron–positron pair and then return to being a photon. This transformation forms the content of the interaction V discussed in the previous example, with interactions from several instances of V corresponding to the photon repeatedly turning into electron–positron pairs. In this case, however, a more intuitive picture of the charge is possible in terms of classical screening [Figure 22.1(b)], where the electron–positron pair may be thought of as a polarized dipole. The bare electron is then viewed as being surrounded by these dipoles. An attempt to measure the electronic charge from a large distance measures not only the negative charge of the electron but also the combined effect of this and the screening supplied by all of the surrounding dipoles. Viewed macroscopically, it is as if the vacuum itself has become a polarized dielectric, which attenuates the charges of electrons.

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Lessons from our case study discussion of metals can be carried over to the consideration of QED and quasielectrons more generally. In the metal, we are able to distinguish between the behaviour of the electron outside of the metal and the electron quasiparticle that exists within the metal. This allows a story to be told about the dressing up of a bare electron into a quasielectron via the electron–electron interactions. The vacuum for the metal in field theory is the Fermi gas of electrons. Quasiparticle excitations then exist as excitations out of that gas. QED is different, in that it is impossible to take an electron outside of the QED vacuum, since this would imply removing electrons from our universe. So while it is difficult to tell the same story about dressing a QED electron in QED interactions, we are left with an idea that the QED electron is itself a quasiparticle, whose properties emerge from interactions with the interacting vacuum of the universe.

Infinities

Renormalization provides an account of masses and charges being altered by interactions with the vacuum. It is especially important in the case of many field theories (including QED) where a troubling aspect of the theory arises: the presence of infinities. In fact, renormalization is often presented as a method of removing infinities from a theory. However, it should be borne in mind that renormalization is a necessary step in formulating a field theory. It is required to take interactions into account and is independent of the mathematical presence or absence of any infinities (Weinberg 1995; Lancaster and Blundell 2014).

Mathematically, infinites are found in many theories and result from attempts to sum behaviour over all length scales (or equivalently, over all momentum scales). They can be avoided if we admit that there is a minimum length scale in any practical case below which we don't have access. (See later for a discussion of this.) Technically one can then identify where this minimum length scale, known as a cut-off, occurs in the theory and then add extra terms to the Lagrangian (known as counter-terms) in order to remove them. The consequence of this strategy is to shift the values of the mass and the charge of the theory. The theory is therefore saved mathematically via the removal of infinities, with the shifts (i.e. renormalization) in particle properties a by-product.

We can trace the origin of the shifts in particle properties as follows. We start doing QFT by writing down a Lagrangian that (we assume) accounts for the properties of some area of physics. This includes parameters such as the masses and charges of particles. If the theory has interactions, we usually can't solve it, and so we rely on perturbation theory to account for the properties of a system. This involves expanding a series about the values of the parameters m and g that we believe happen to encode the particle properties. However, it turns out we were wrong! We were doing the wrong perturbation theory as we were expanding about the wrong masses and wrong couplings. We had asked a nonsensical question and ended up with an infinite (and therefore meaningless) answer. The couplings we should have been using are those found experimentally in nature. Renormalization is not, therefore, a formal, mathematical exercise in hiding infinities; it is an exercise in shifting the parameters that describe particle properties so that they describe nature as we experience it. For many theories, the magnitudes of the shifts in particle properties are infinite. This implies that the bare values of mass m and coupling g are infinite, and we shift from them by an infinite amount to obtain a sensible (finite) value. This is certainly a troubling consequence of renormalization that has caused practitioners of QFT much pain and uncertainty. However, the instructive use of a set of techniques known as renormalization group analysis has provided a great deal of insight into the origin and explanation of these infinities.

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Renormalization group

A valuable insight into the meaning of renormalization in QFT is found in renormalization group (RG) analysis, discussed elsewhere in this volume (Blundell 2019). RG has been the subject of much discussion in the philosophical emergence literature in the context of its usefulness in describing phase transitions (see e.g. Batterman 2001; Morrison 2012), where its use in identifying critical points and analysing universality classes of models allows an insight into multiple realizability in condensed matter physics. For our purposes, RG allows us an insight into how the dressing-up process of renormalization affects particle properties through the notion of *scale*.

In examining quasiparticles, rather than correct for infinities as is done in renormalization, in RG the cut-off is used as a handle that allows us to vary the length scale of interest in a problem. Specifically, the physical content of the cut-off is that it provides a length below which all details of the physics are irrelevant to the problem at hand. Consider, for example, some atoms in a box. If we are interested in sound waves, then the length scale of interest is the wavelength of sound (centimetres), and the cut-off might be a few microns. If we're interested in the physics of the electrons in the atoms of the gas, then the length scale of interest will be the size of an atom, and the cut-off should be set to the size of a nucleus. The renormalization group analysis consists of asking how the coupling constants (masses and charges again) vary as we change the value of this cut-off.

For theories of condensed matter physics, the motivation for varying the cut-off is to examine the limit of behaviour measured by experiment. Solid materials are generally macroscopic, and so the length scale of interest in condensed matter physics is that in which, on the scale of the underlying atoms, length tends to infinity. Examining the variation of the coupling constants in this limit then allows insights into the qualitative and quantitative details of behaviour.[4] The size of a coupling constant will change as we interrogate the system on different length scales. The result is that certain interactions will be important in a particular limit, while others will die away.

Although this all sounds rather artificial, it is a real possibility in experiments involving particle collisions, where the momentum involved in the collision, and hence the length scale probed, can be varied. This, in turn, provides an insight into the particle "dressing-up" scenario that we have described earlier. High energy and momentum corresponds to small length scales, so high-momentum collisions allow us to probe electronic charges at very low length scales. It is found that the charge on an electron appears to increase at these small probing distances. This corresponds to our measuring the electronic charge inside the screening cloud of polarized electron–positron dipoles [described earlier and shown in Figure 22.1(b)]. As the probe momentum decreases (i.e. length scale increases), the probing distance is greater, there is more screening between the bare particle and the probe and we correspondingly measure a smaller charge. One consequence of this is that many of the "fundamental" physical constants of nature, such as electronic charge, should not be regarded as constants. They depend on the length scale at which we measure them, and this is a consequence of the quasiparticle view of particles.

Finally, it is worth noting that a further consequence of the RG view of nature is that all theories (described by a Lagrangian in the sense noted earlier) are *effective theories*. They may contain terms that operate at momentum scales so large that we never detect them. As a result, the physics we have is a low-energy approximation to a more complete field theory of nature. This feature has led to philosophers such as Jonathan Bain to suggest that effective theories provide an explanation of emergent phenomena in systems described by field theories (Bain 2013).

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Summary: emergence of particle properties in QFT

The quasiparticles concept implies that the properties of particles (their masses and other couplings) take on values by virtue of the interaction of the fields in a theory. These interactions are themselves mediated via particles, so we might say that the many-particle nature of QFT is what gives rise to the couplings we observe. It is perhaps in this sense that particles' properties should be thought of as emergent. Moreover the ability to describe and predict the details of the dressing-up process makes this a rather weak form of emergence when compared to some of the more dramatic phenomena outlined later, which resist such a description. Whether we regard the bare particles of a theory realistically depends, according to Sidney Coleman (Coleman 2018), on "how weird you're willing to believe the world is". That is, one can imagine an extreme realist position where the electronic charge is infinite, but we never observe this infinity due to the spontaneous existence of particles that screen the infinite charge.

From the point of view of renormalization group analysis, the emergence of particle properties is bound up in notions of scale. Specifically, different properties emerge at different scales. We implicitly accept that a theory will have a realm of applicable scales, and insight is provided through considering how the theory must change to remain consistent as the scale of interest is varied. The RG picture gives us a clue as to the reality of the dressing-up process that turns particles into quasiparticles. It tells us that the electronic charge depends on the momentum used to investigate it. If one attempts to investigate the electronic charge at sufficiently small distances (using very high momentum probes), then the measured charge of the electron is actually larger. This makes sense in the context of the quasiparticle view of the world as a particle being dressed in interactions.

Broken symmetry and topological objects

A very clear link between particle properties and emergence is found in the much-discussed topic of broken symmetry (Blundell 2019). In fact, Philip Anderson discusses this topic in terms of a breakdown of adiabatic continuity (Anderson 1984). In the case of the metal, discussed earlier, turning on interactions causes particle properties to vary smoothly. This might be characterized in terms of a weak form of emergence, where we follow the gradual change in particle properties as interactions are dialled up. However, in the case where turning on interactions causes the particle energy levels to cross, there is a breaking of symmetry. The result is that the particles effectively lose their identity in the resulting reconstruction of a system. The properties of the new excitations realized in the broken symmetry state might then have a claim to be realized by a stronger form of emergence.

In QFT we think of the production of a particle as involving the excitation of the ground state. Breaking a symmetry in QFT is characterized in terms of a dramatic change in the potential energy density (which is part of the Lagrangian describing the system) as shown in Figure 22.2. When a symmetry is broken, the nature of the vacuum state itself is reconstructed, with the two-fold consequence that (i) the properties of the particles change and (ii) qualitatively new varieties of particles are possible.

The first of these features is due to the details of the change in the shape around the minima of the potential being reflected in a change in the Lagrangian parameters, from which the particle properties follow. The second of these is more dramatic and reflects the fact that the new potential landscape allows new forms of excitation within it.

Qualitatively new particles arise, for example, when a *continuous symmetry* is broken. In this case, new forms of particle excitation are possible, whose very existence could not be supported

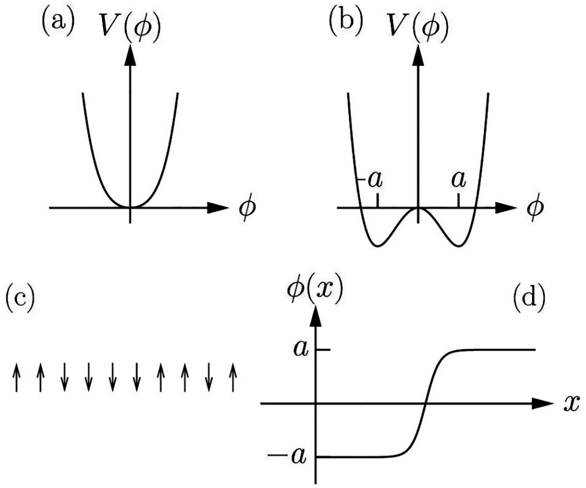

Figure 22.2 The dramatic change in potential energy density caused by the breaking of a symmetry. (a) Before symmetry breaking; (b) after symmetry breaking. (c) A model of a one-dimensional magnet with spins (arrows) arranged randomly. (d) The kink in a field theory can be thought of as an ordered state where arrows [represented here by a field $\phi(x)$] go through a transition as a function of distance x, from being aligned downwards (on the left) to upwards (on the right).

by the non-broken symmetry state. An example is found by examining the freezing of a liquid into a crystalline solid. The solid is a state of broken *translational* symmetry. On average, the density of atoms is the same at every point in a liquid. In a solid, atoms are arranged into a periodic array and are only found at regular distances from other atoms. The solid supports the existence of a new particle excitation: the phonon. Phonons are massless excitations, which means they can be created with arbitrarily low energy. In addition, the number of phonons is not a conserved quantity. In general, the emergence of massless particles on symmetry breaking is not too mysterious – it simply reflects the fact that the broken symmetry potential can support a different manner of excitation from the non-broken symmetry potential. At an atomic level, the phonon quasiparticle owes its existence to all of the coordinates of the underlying crystal lattice changing. Phonons are therefore sometimes called quanta of lattice vibrations. As a result of this rather extreme reliance on the existence of the ordered crystal, the phonon can't exist outside of the solid.

There is still another class of particle-like objects that exist in theories with broken symmetry: these are *topological objects*. The key point to these objects is that they owe their existence to a system-breaking symmetry in different ways in different spatial locations. The simplest example is a magnet with spins (i.e. magnetic moments, or arrows) arranged along a line [Figure 22.2(c)]. The non-broken symmetry system is disordered, with the spins arranged randomly pointing up or down. (In the language of potential energy [Figure 22.2(a)], the system sits in the minimum of the potential well.) In breaking symmetry in this system, the spins must collectively align up or down. [This corresponds to the system choosing to sit in one of the two minima of the double well potential in Figure 22.2(b).] However, a more interesting configuration is possible if half of the spins on the left of the chain order down, while the other half align up. This leads to a transition region in the middle of the chain known as a *kink* [Figure 22.2(d)]. The kink is yet another form of particle-like excitation. This object may be shown to exist with a finite energy above

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the ground state. It can be moved around, rotated and translated just like any other particle. The main difference is that it is a spatially extended object, in a sense quite unlike the spatial extension of the quantum particle localized in a momentum state.

Kinks certainly exist in nature, most obviously in magnets where they are known as domain walls and are used extensively in magnetic memory applications. The kink is a good example of the general concept of a *topological object*. Topological objects are (roughly speaking) a sort of untieable knot, in that it is impossible to remove them from a system without incurring an enormous energetic cost. These objects are therefore stable. It is impossible to deform the kink shape without costing the system a (semi-) infinite amount of energy, since this would involve lifting half of the kink shown in Figure 22.2(d) over the potential barrier shown in Figure 22.2(b). Like many particles, kinks carry a conserved quantity. However, this is not the usual conserved charge (guaranteed by symmetry due to Noether's theorem). Rather, it is a topological charge, independent of the geometry of space-time. [Technically, the quantity does not involve the metric $g_{\mu\nu}$ (Lancaster and Blundell 2014).]

Kinks exist as a result of the one-dimensional physics of a chain of spins [Figure 22.2(c)]. In two dimensions we have a new topological object: the vortex, which looks rather like a whirlpool in a liquid. Again, this is formed by considering configurations that exist in different broken symmetry ground states in different regions of space. In three dimensions the analogous particle-like object is the monopole (also known as a hedgehog, due to its resemblance to that mammal, when the field is drawn as arrows). In addition, there exist many more examples of objects whose stability is due to topological considerations, which emerge on symmetry breaking and which deserve serious consideration as particles.[5] In summary, the topological objects we have described are a species of excitation since they have a finite energy above the system's ground state. They can exist only by virtue of the system breaking a symmetry. They are necessarily extended over space, so are rather different from classical point-like particles, and so one might be tempted to classify them as non-fundamental. They do, however, have many of the fundamental symmetry properties of particles (they can be translated and transformed in the same way that particles can), so bear a very strong resemblance to simple particles. Finally, just like the new particles that exist only by virtue of symmetry breaking, the topological objects lose their meaning in a system without broken symmetry and therefore have the same emergent character.

Conclusions

We conclude with some observations about the emergent nature of quasiparticles. There are many ways to characterize weak and strong forms of emergence (Lancaster and Pexton 2015), invoking, for example, (i) a failure of explanatory reduction; (ii) the appearance of new entities; (iii) novel systematic properties; and (iv) a failure of mereological supervenience, leading to the necessity of considering whole-system properties.

Perhaps the key factor in all of the examples presented has been interactions. A system without interactions might support particle excitations, but these will not be detectable as they pass through each other (and everything else). As soon as interactions are added to a system, they change the properties of the particles. In fact, the emergence, existence and demergence of a particle are quite dramatic in Landau's quasiparticle picture. Assuming that no symmetry is broken, the properties of the particles can be shown to evolve continuously from the non-interacting theory upon turning on the interactions. We might characterize this as the emergence of the particle properties. If we introduce a new quasiparticle into the metal, then the length of time it will retain its identity as a particle is a continuous function of its energy. Too far from the Fermi energy, and it ceases to exist as a particle over the lifetime of an observation, instead being

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consumed by the liquid-like system. Its existence becomes longer the closer it gets in energy to the Fermi energy. The short lifetimes of those excitations with energies well below the Fermi energy cause them to lose their meaning as particles and instead exist as slightly ill-defined parts of a whole metallic system. It is worth noting in this context that there are also more exotic cases of quasiparticles, such as a phenomenon known as the fractional quantum Hall effect (Lancaster and Pexton 2015), where if we introduce an extra electron to the system, it will fall apart into the allowed quasiparticle excitations of the quantum Hall fluid. This is particularly dramatic in this case since the fractional quantum Hall quasiparticles carry fractions of electronic charge. There is a sense in which the electron gives up its properties to the fractional quantum Hall fluid, where they reappear as fractionally charged quasiparticles.

Interactions can also lead to the spontaneous breaking of symmetry. Symmetry breaking leads to the possibility of new forms of particles, such as the massless particles like the photon, and the massive W and Z particles in the standard model of particle physics. In the broken symmetry case we also have a new class of excitations that have a topological character such as kinks or vortices. Both of these types of broken symmetry particles might be considered under each of the characterizations of emergence given earlier: (i) A failure of explanatory reduction occurs due to a singularity in the free energy on symmetry breaking. (More specifically, if you use perturbation theory to perturb around the symmetric ground state, you won't derive the existence of the broken symmetry particles, nor the topological particles, which are fundamentally non-perturbative. If we were to heat the system up and regain the symmetry, then these particles would cease to exist.) (ii) These quasiparticles are qualitatively new excitations that could not exist before symmetry breaking. (iii) They exist by virtue of a systematic restructuring of a system on symmetry breaking.

However, what all of the quasiparticles we have considered have in common (i.e. Landau/field theory quasiparticles and broken symmetry quasiparticles) is that their emergence follows from (iv): the necessity of considering whole-system interactions. This is most obvious for the collective excitations, such as the phonon, for example, which represent a change in coordinates of all of the atoms making up a crystal. However, there is a sense in which it is true for all quasiparticles since, in all cases, the renormalization of particles involves them dressing themselves in interactions, which are collective expressions of the excitations of the vacuum of the whole system. Finally, the theme of internal interactions leading to constraints on the whole of a system seems to be common in describing many of the candidates for emergent phenomena examined within the physical sciences, both in this volume and elsewhere (e.g. Clark and Lancaster 2017) and has become an increasingly pervasive theme in modern condensed matter physics research, where notions of fields, topology and emergence are now integral concepts.

Notes

1 The Hamiltonian and related concepts are discussed in the chapter by Clark and Thomas in this volume (Clark and Thomas 2019).

2 A distinction may be made between two sorts of elementary excitation of a system. The first are called *collective excitations* and correspond to a change in coordinates of all particles in a system (and we discuss these later in this chapter). The second are the quasiparticles described here. Examples of the latter include the quasielectrons in a metal.

3 This is not as artificial as it might sound. In some materials, interactions can be tuned by applying hydrostatic pressure to the solid system, for example.

4 It is also possible to ask about the very low length scale regions of the theory, which corresponds to high-energy measurements carried out by particle physicists, for example.

5 Interestingly, we can also go further and investigate these structures in time rather than space. The kink in time is known as an *instanton* and has found uses mathematically in solving tunnelling problems.

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