The Process of Generation of Mass, The Higgs Boson, and Dark Matter

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A dynamical mechanism of symmetry breaking in which gauge and matter fields play an active role is proposed. It basically represents a covariant generalization of the mechanism responsible for superconductivity, and provides a natural mechanism of generation of mass which is not in conflict with the present value of the cosmological constant. When applied to SU(2)×U(1) leads to exactly the same physics (Lagrangian density) as the Standard Model but modifying only the Higgs sector. It also predicts the appearance over all space of a classical scalar field as well as the existence of density fluctuations. According to it, space would be filled with a macroscopically large number of Higgs bosons which now appear as light, stable scalar particles decoupled from ordinary matter and radiation. Therefore they would play the same role as the Cooper pairs in superconductivity and would be a natural candidate for dark matter.

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Spontaneous symmetry breaking is one of the basic underlying ideas of both Unified Gauge Theories and Inflationary Cosmology. This phenomenon occurs when the ground state of a system displays less symmetry than that of the corresponding Hamiltonian, and is related to the concept of vacuum and the process of mass generation. Its importance lies mainly on the fact that the only renormalizable gauge theories with massive vector bosons are gauge theories with spontaneous symmetry breaking \[1\]. In modern Unified Theories massive intermediate bosons are properly introduced by breaking the symmetry by means of the Higgs mechanism. It basically consist in the introduction of a complex scalar field \( \phi(x) \) subject to an effective potential of the form

\[
V(\phi) = \mu^2 \phi^2 + \lambda (\phi^2)^2
\]

with \( \mu^2 < 0 \). The necessity of introducing a negative mass squared for the scalar field is the price one has to pay in order to generate a nonzero stable configuration \( \langle \phi \rangle_0 \).

In this paper I propose a different mechanism where spontaneous symmetry breaking has its origin in the dynamics of the scalar field \( \phi \) interacting with gauge and matter fields, which now play an active role. In fact, symmetry breaks once a particular solution of the dynamical equations is chosen among a (one-parameter) family of possible solutions.

In what follows I will consider in detail the case of an abelian U(1) theory, which basically represents a relativistic generalization of the mechanism responsible for superconductivity. This analogy with superconductivity proves to be very useful in order to gain valuable insights about the process of symmetry breaking. Specific application to the Standard Model (SM) will be considered below.

The U(1) gauge invariant Lagrangian density for a complex scalar field \( \phi(x) \) reads

\[
\mathcal{L} = (D_\mu \phi)^+ (D^\mu \phi) - \mu^2 \phi^2 - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}
\]

where now \( \mu^2 > 0 \) is the mass of the scalar field and the \( F_{\mu\nu} \) tensor is given by \( F_{\mu\nu}(x) = \partial_\mu B_\nu - \partial_\nu B_\mu \).

It is convenient parametrize \( \phi(x) \) as

\[
\phi(x) = \frac{1}{\sqrt{2}} \rho(x) \exp i \xi(x)
\]

where \( \rho(x) \) and \( \xi(x) \) are real scalar fields. Then, in the unitary gauge, where \( \phi(x) \to \rho(x)/\sqrt{2} \), the Lagrangian density can be rewritten as

\[
\mathcal{L} = \frac{1}{2} (\partial_\mu \rho)(\partial^\mu \rho) + \frac{1}{2} g^2 A_\mu \rho^2 - \frac{1}{2} \mu^2 \rho^2 - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}
\]

with \( A_\mu(x) = B_\mu(x) - (1/g) \partial_\mu \xi(x) \).

The Euler-Lagrange equations of motion which describe the coupled dynamics of the fields, are

\[
\partial_\mu \partial^\mu \rho(x) = -\mu^2 \rho(x) + g^2 A^2(x) \rho(x)
\]

\[
\partial_\alpha \partial^\alpha A^\nu - \partial^\nu (\partial_\beta A^\beta) = -g^2 A^\nu \rho^2
\]

where the source of the gauge field, \( J^\mu(x) = A^\mu(x) \rho^2(x) \), is a conserved current.
In what follows I will show that the equations of motion (5)–(6) admit a family of solutions which verify

$$A^\mu(x) = \frac{J^\mu(x)}{\rho^2(x)}$$

(7)

where $J^\mu(x)$ is not an independent function of $\rho(x)$, and a subset of these solutions leads to a theory with broken symmetry and massive gauge bosons.

By substituting Eq.(7) into Eq.(5) we obtain

$$\partial_\alpha \partial^\alpha \rho(x) = -\frac{\partial}{\partial \rho} V_{\text{eff}}(x)$$

(8)

where

$$V_{\text{eff}}(x) = \frac{1}{2} \mu^2 \rho^2(x) + g^2 \frac{J^2(x)}{2\rho^2(x)}$$

(9)

In field theories the vacuum $|0\rangle$ is defined as the ground state of the theory. The property of translational invariance that this state must possess requires vacuum expectation values (VEVs) of physical magnitudes to be constants independent of space-time coordinates. Then we have

$$\langle 0 | \rho(x) | 0 \rangle = \langle \rho \rangle_0 = v$$

$$\langle 0 | A_\mu(x) A^\mu(x) | 0 \rangle = \langle A_\mu A^\mu \rangle_0$$

(10)

$$\langle 0 | J^2(x) | 0 \rangle = \langle J^2 \rangle_0 = l^2$$

so that the vacuum $|0\rangle$ must verify

$$\left. \frac{\partial}{\partial \rho} V_{\text{eff}}(x) \right|_0 = \mu^2 \langle \rho \rangle_0 - g^2 \langle J^2 \rangle_0 \langle \rho^3 \rangle_0 = 0$$

(11)

which is nothing but the equation obtained by projecting Eq.(8) onto the vacuum. Different values of $l^2$ characterize different possible solutions and therefore different possible vacuum states. This means that the properties of the vacuum depend on the current squared $J^2(x)$. Thus gauge (and matter) fields play an active role in selecting the ground state $|0\rangle$. Note, however, that due to the fact that $l^2$ is a constant it is not possible to smoothly pass from one to another solution. We have indeed a one-parameter family of stable solutions, and the U(1) symmetry breaks once a particular one is taken.

The subset of solutions we are interested in are those with $l^2 > 0$, because as can be seen from Eq.(11) for $l^2 \neq 0$ it follows that $\langle \rho \rangle_0 = v \neq 0$. In this case we have

$$\mu^2 \langle \rho \rangle_0 \langle \rho^3 \rangle_0 = g^2 l^2$$

(12)

Furthermore, Eq.(7) leads to

$$l^2 = \langle \rho^4 \rangle_0 \langle A_\mu A^\mu \rangle_0$$

(13)
Therefore the solutions we are interested in satisfy
\[ \mu^2 \langle \rho \rangle^3_0 = g^2 \langle \rho^4 \rangle_0 \langle A_{\mu} A^\mu \rangle_0 \quad (14) \]

Notice that a similar equation, which relates the mass of the scalar field \( \rho(x) \) to the vacuum fluctuations of the gauge field, can be directly obtained from the initial equations of motion. Indeed by projecting Eq.(4) onto the vacuum one finds that solutions with \( \langle \rho \rangle_0 \neq 0 \) verify
\[ \mu^2 = g^2 \langle A_{\nu} A^{\nu} \rangle_0 \quad (15) \]

This equation represents a necessary condition for a solution with \( \langle \rho \rangle_0 \neq 0 \) can be taken. It basically states that in order for the process to take place, \( \mu \) particles should be created from energy fluctuations of gauge fields in the ground state.

By comparing Eqs.(14) and (15) one expects
\[ \langle \rho^n \rangle_0 \simeq \langle \rho \rangle^n_0 = v^n \quad (16) \]

so that we can rewrite Eq.(12) as
\[ v^4 = \frac{g^2 \mu^2}{m^2} \quad (17) \]

On the other hand, by expanding \( V_{\text{eff}}(x) \) about the vacuum one obtains
\[ V_{\text{eff}}(x) = V_{\text{eff}} \bigg|_{0} + \frac{1}{2!} \frac{\partial^2}{\partial \rho^2} V_{\text{eff}} \bigg|_{0} (\rho(x) - v)^2 + \frac{1}{3!} \frac{\partial^3}{\partial \rho^3} V_{\text{eff}} \bigg|_{0} (\rho(x) - v)^3 + \ldots \quad (18) \]

where terms proportional to \((J^2(x) - l^2)\), which have no influence on the dynamics of \( \rho(x) \), have been omitted. Substitution of Eq.(18) into Eq.(8) leads to
\[ \partial_\alpha \partial^\alpha \rho(x) + m^2 (\rho(x) - v) = \lambda (\rho(x) - v)^2 + \ldots \quad (19) \]

where
\[ m^2 = \frac{\partial^2}{\partial \rho^2} V_{\text{eff}} \bigg|_{0} = \mu^2 + 3 \frac{g^2 l^2}{v^4} = 4 \mu^2 \quad (20) \]
\[ \lambda = -\frac{1}{2} \frac{\partial^3}{\partial \rho^3} V_{\text{eff}} \bigg|_{0} = 6 \frac{g^2 l^2}{v^5} = 6 \mu^2 \quad (21) \]

Then defining a real scalar field with vanishing VEV
\[ \eta(x) = \rho(x) - v \quad (22) \]

the equations of motion finally read
\[ (\partial_\alpha \partial^\alpha + m^2) \eta(x) = \lambda \eta^2(x) + \ldots \quad (23) \]
\[ (\partial_\alpha \partial^\alpha + m^2 A^\nu) - \partial^\nu (\partial_\beta A^\beta) = j^\nu (x) \quad (24) \]
where
\[ j^\nu(x) = -(2gm\eta(x) + g^2\eta^2(x))A^\nu(x) \] (25)

From these equations we see that for a solution with \( l^2 > 0 \), the scalar field develops a nonvanishing VEV
\[ \phi(x) = \frac{1}{\sqrt{2}}(v + \eta(x)) \exp i\xi(x) \] (26)

and the gauge field \( A_\mu(x) \) acquires a mass
\[ m_A^2 = g^2v^2 \] (27)

Furthermore, according to Eq. (23) the field \( \eta(x) \) (as well as \( \rho(x) = v + \eta(x) \)) decouples. Thus, they become no affected by the process of interaction with gauge fields. Indeed its dynamics becomes governed by the effective Lagrangian density
\[ \mathcal{L}_{\text{eff}} = \frac{1}{2} (\partial_\mu \eta)(\partial^\mu \eta) - \left( V_{\text{eff}}|_0 + \frac{1}{2} m^2 \eta^2 - \frac{\lambda}{3} \eta^3 + \ldots \right) \] (28)

Quantum fields exhibit this kind of behaviour in the limiting case where a macroscopically large number of quanta appear in a coherent state, and in such a situation they behave as ordinary classical fields.

In particular, given a solution with \( l^2 > 0 \), the field \( \rho(x) \), which describes \( \mu \) particles, behaves as a classical scalar field which fluctuates about \( \rho(x) = v \) with a mass \( m = 2\mu \). Then one expects that a macroscopic physical meaning should be possible to be given to \( \rho(x) \). Indeed, it can be interpreted as describing a Bose condensate of \( \mu \) particles with vanishing momenta and a number density given by
\[ n(x) = \frac{1}{2} m\rho^2(x) \] (29)

With this interpretation the minimum of the potential energy density corresponds, as expected, to the internal energy of a Bose condensate. In fact, it can be written as
\[ V_{\text{eff}}|_0 = \mu^2v^2 = \mu n_0 \] (30)

with \( n_0 = \langle n(x) \rangle_0 \) being the mean particle density. On the other hand, according to Eqs. (27) and (29), the gauge boson acquires a mass
\[ m_A^2 = \frac{g^2n_0}{\mu} \] (31)

which also agrees with what one would expect. In fact, \( m_A^{-1} \) coincides with the London penetration depth for a condensate of \( \mu \) particles with number density \( n_0 \). Furthermore, for small fluctuations about the equilibrium configuration, \( \eta(x) \ll v \), Eq. (29) leads to
\[ \eta(x) \simeq \frac{1}{mv}(n(x) - n_0) \] (32)
which shows that the classical scalar field $\eta(x)$ can be regarded in turn as describing density fluctuations about the mean value $n_0$. Incidentally, note that the mass of this field corresponds to the energy necessary to create a pair of $\mu$ particles.

Thus Eqs. (23)–(24) would describe a gauge field which evolves in a classical scalar medium, acquiring mass and producing in turn density fluctuations in this medium. These results are consistent with the interpretation that for a solution with $l^2 > 0$, a phase transition takes place and a Bose condensate of $\mu$ particles is formed.

Once the phase transition takes place, the dynamics of $A_\mu(x)$ becomes governed by the Lagrangian density

$$L' = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \frac{1}{2} m^2 A_\mu A^\mu + L'_1$$

where

$$L'_1 = g m_{\nu} A_\mu A^\mu + \frac{1}{2} g^2 \eta^2 A_\mu A^\mu$$

(34)

describes the interaction of the quantum field $A_\mu(x)$ with the external classical field $\eta(x)$ governed by the Lagrangian density $L_{\text{eff}}$ given by Eq. (28).

In order to get a better understanding on the mechanism which causes the decoupling of the scalar field, let us rewrite Eq. (23) retaining now the term in $(J^2(x) - l^2)$ which was neglected before

$$(\partial_\alpha \partial^\alpha + m^2)\eta(x) = \frac{3m^2}{2v} \eta^2(x) + \frac{g^2}{v^3} (J^2(x) - l^2) + \ldots$$

(35)

This equation explicitly shows that the decoupling occurs when $v \to \infty$, which according to Eq. (29) means $n_0 \to \infty$. Thus the decoupling is caused by the presence of a macroscopically large number of quanta of the scalar field, which then behaves as a classical field.

Note finally that the classical external medium seems to play the role of a regulator. In fact, if one tries to decouple the massive gauge field from the scalar medium by taking $g \to 0$ in Eqs. (24)–(25), then Eq. (15) would imply $\langle A_\mu A^\mu \rangle_0 \to \infty$.

The above mechanism can also be applied to the SU(2) $\times$ U(1) Standard Model in a straightforward manner. However, in this case it is also possible to break the symmetry starting from a massless scalar field, which seems a more natural alternative. In what follows I will briefly consider this case. A detailed derivation will be given elsewhere [2].

We start from the Lagrangian density of the SM with the only modification that now the Higgs potential vanishes. Then, in terms of eigenstates of the mass operator the Euler-Lagrange equation of motion for the massless scalar field, in the unitary gauge has the form

$$\partial_\mu \partial^\mu \rho(x) = \frac{1}{4} \left( 2g^2 W^- W^+ + (g^2 + g'^2) Z_\mu Z^\mu \right) \rho(x)$$

$$- \sum_f \frac{\lambda_f}{\sqrt{2}} \bar{\psi}_f \psi_f$$

(36)

where $W^+_\mu$ and $Z_\mu$ are the mass eigenstates of the weak vector bosons, $\psi_f$ denotes a Dirac spinor, and the sum runs over massive fermions corresponding to the three families. As for
the equations of motion of the gauge and matter fields, they are exactly the same as those of the Standard Model and are not explicitly shown.

In order to get a theory with broken symmetry and massive gauge and fermion fields, we look for solutions of the equations of motion with the requirement that the sources of the ‘free’ gauge and matter fields are not independent functions of $\rho(x)$. In this case Eq.(36) takes the form of Eq.(8), where now $V_{\text{eff}}(x)$ is given by

$$V_{\text{eff}}(x) = \frac{1}{8} \frac{J_G^2(x)}{\rho^2(x)} - \frac{J_F^2(x)}{\rho(x)}$$  \hspace{1cm} (37)$$

and here $J_G^2(x)$ and $J_F^2(x)$ play the role of gauge and fermion currents squared, respectively.

On the other hand, by projecting onto the vacuum $|0\rangle$ we find that this state must verify

$$\left. \frac{\partial}{\partial \rho} V_{\text{eff}}(x) \right|_0 = -\frac{1}{4} \langle J_G^2 \rangle_0 + \langle J_F^2 \rangle_0 = 0$$  \hspace{1cm} (38)$$

It is not difficult to see (by projecting Eq.(36) onto $|0\rangle$ and comparing with Eq.(38)), that these solutions lead to a nonvanishing VEV $\langle \rho \rangle_0 = v \neq 0$ only if $\langle \rho^4 \rangle_0 = \langle \rho^3 \rangle_0 \langle \rho \rangle_0$, so that we are again led to look for solutions satisfying Eq.(16). Then, from Eq.(38) we have

$$v = \frac{1}{4} \langle J_G^2 \rangle_0 \langle J_F^2 \rangle_0$$  \hspace{1cm} (39)$$

which shows that a particular solution with $\langle J_G^2 \rangle_0, \langle J_F^2 \rangle_0 > 0$ breaks the symmetry.

By expanding $V_{\text{eff}}(x)$ about the vacuum, as before, one finally obtains

$$(\partial_\alpha \partial^\alpha + m_H^2) \eta(x) = \lambda \eta^2(x) + \ldots$$  \hspace{1cm} (40)$$

where $\eta(x) = \rho(x) - v$, and

$$m_H^2 = \frac{2}{4} \langle W_{\mu}^+ W_{\mu} \rangle_0 + \frac{1}{4} \langle g^2 + g'^2 \rangle \langle Z_{\mu} Z^\mu \rangle_0$$

$$= \sum_f \frac{\lambda_f}{\sqrt{2} v} \langle \bar{\psi}_f \psi_f \rangle_0$$  \hspace{1cm} (41)$$

$$\lambda = \frac{3 m_H^2}{v}$$  \hspace{1cm} (42)$$

Substitution of $\rho(x) = v + \eta(x)$ into the equations of motion of the gauge and matter fields then leads to exactly the same equations for the corresponding massive quantum fields as the SM. This means that the dynamics of the gauge and matter fields would be governed by an effective Lagrangian density which coincides with that of the SM but interpreting now the field $\eta(x)$ as an external classical field. Therefore the present mechanism would lead to exactly the same physics as the SM except for the Higgs sector.

On the other hand, when $\rho(x)$ evolves towards its stable configuration in $V_{\text{eff}}(x)$, it acquires a nonvanishing mass $m_H$ from vacuum fluctuations of quantum fields. In this process a phase transition takes place and it develops a nonzero VEV $\langle \rho \rangle_0 = v$, reflecting the appearance of a macroscopically large number of quanta of the field. These quanta are
Higgs bosons created from the vacuum energy. The process gives rise to the formation of a Bose condensate of Higgs particles, and as a consequence $\rho(x) = v + \eta(x)$ behaves as a decoupled classical field. Then, as long as the Standard Model remains valid everywhere, there must exist a macroscopic classical scalar field over all space. This can give us a relation between microphysics and the macroscopic world. According to this mechanism starting from a microscopic system interacting with a scalar field, as a result of a phase transition a macroscopic classical medium can arise where quantum fields evolve. Gauge and fermion fields would acquire mass as a consequence of the inertia they manifest evolving through this classical medium. In fact, this is a natural mechanism for particles to get mass. Motion of charged particles inside a crystal lattice, through a plasma, or through a strong electromagnetic field are some examples where particles respond with a larger inertia.

According to the present mechanism space would be filled with a macroscopically large number of Higgs bosons created in the early universe from the vacuum energy. Being stable, scalar particles decoupled from ordinary matter and radiation, the Higgs bosons only would be detectable by gravitation and they would be a natural candidate for dark matter.

On the other hand, it can be shown that the vacuum energy density $\rho_V$, given by the minimum of $V_{\text{eff}}(x)$, takes the form

$$\rho_V = V_{\text{eff}}|_\langle \rangle = -\frac{1}{2} m_H^2 v^2$$

(43)

Thus, the experimental value of $\rho_V$ can be used in order to estimate an upper bound to the mass of the Higgs boson. From Eq.(43), one has

$$m_H^2 = \frac{2|\rho_V|}{v^2}$$

(44)

so that in order to be consistent with the present value of the cosmological constant, $\Lambda = 8\pi G|\rho_V| \leq 10^{-84}\text{GeV}^2$, it is only required

$$m_H \leq 10^{-26}\text{GeV}$$

(45)

It should be notice that this bound is not an specific feature of the mechanism of symmetry breaking proposed in this paper. In fact, if one demands the Higgs mechanism to be consistent with the present value of the cosmological constant then one obtains the same result [3]. However, the Higgs mechanism cannot consistently account for both the scale of electroweak symmetry breaking (which requires $v \simeq 250$ GeV) and such a small mass for the Higgs boson (or what is the same, such a small value for the cosmological constant). The reason is that the Higgs mechanism leads to $m_H^2 = 2\lambda_c v^2$ ($\lambda_c$ being the quartic coupling constant) so that one expects $m_H$ to be of order $\mathcal{O}(v)$ (unless $\lambda_c$ takes an unnatural small value). According to the present mechanism, however, the scale of the electroweak symmetry breaking and the mass of the Higgs boson are related only through the value of the vacuum energy density, Eq.(43), and the surprisingly small value of $m_H$ would simply reflect the small value of the cosmological constant.

In order to try to understand how such a small mass could be understood in the context of the Standard Model let us take a closer look at the way in which masses generate. Gauge and matter fields acquire mass evolving inside a dense medium described by a field $\rho(x) = v + \eta(x)$ with a nonvanishing mean value. This is basically a classical mechanism
(apart from the well known examples of superconductivity or motion of classical particles inside dense media, consider, for instance, plasma oscillations or the screening mechanism of Debye-Hückel) and the corresponding masses are of order $O(v)$. However, the way in which the Higgs boson acquires mass is completely different. As Eq. (41) shows its mass appears as a purely quantum effect. In fact, taking into account that the dynamics of the Higgs boson, Eq. (40), becomes governed by the Lagrangian density

$$
\mathcal{L}_{eff} = \frac{1}{2}(\partial_\mu \eta)(\partial^\mu \eta) - \frac{1}{2}m_H^2 \eta^2 + \frac{m_H^2}{v} \eta^3 + \ldots (46)
$$

we see that $m_H$ explicitly breaks the invariance of $\mathcal{L}_{eff}$ under dilatations (scale invariance). Then we would have a partially conserved dilatation current and the mass of the Higgs would reflect the explicit breaking of this symmetry.

On the other hand, as stated before, once the phase transition has taken place the classical field $\rho(x)$ can be regarded as describing a Bose condensate of Higgs bosons with a number density

$$
n_H(x) = \frac{1}{2}m_H \rho^2(x) (47)
$$

Thus the density of dark matter $\rho_{dm}$ would be given by

$$
\rho_{dm} = m_H \langle n_H \rangle_0 = \frac{1}{2}m_H^2 v^2 (48)
$$

so that by using $\rho_{dm} \approx 10^{-29} \text{gr/cm}^3$ one obtains a number density of Higgs bosons

$$
\langle n_H \rangle_0 \approx 10^{21} \text{part/cm}^3 (49)
$$

Notice that Eq. (48) imposes a non trivial relationship among the experimental quantities $v$, $\Lambda$, and $\rho_{dm}$, so that, even if one does not identify the Higgs bosons with the dark matter one would obtain a matter density of Higgs bosons comparable with the density of (dark) matter $\rho_{dm}$.

On the other hand, Eq. (43) can be rewritten in terms of $\rho_{dm}$ in the form

$$
\rho_V = -\rho_{dm} (50)
$$

This equation basically states that the dark matter (Higgs bosons) has been created from the vacuum energy density when the scalar field evolved to the minimum of the effective potential. Therefore this mechanism predicts $|\rho_V| = \rho_{dm}$. This could explain why apparently the vacuum energy density and the matter density are of the same order of magnitude [4].

In the ground state of the theory (i.e. the vacuum) the Higgs bosons are uniformly distributed with a number density $\langle n_H \rangle_0$. According to Eq. (10), in this state they would shield or screen the (negative) energy density of the vacuum

$$
\rho_V + \rho_{dm} = 0 (51)
$$

This situation is quite similar to that encountered in an ionic plasma. In fact, from Eq. (17) follows that small deviations from the ground state give rise to the appearance of density fluctuations
\[ \eta(x) \simeq \frac{1}{m_H v} (n_H(x) - \langle n_H \rangle_0) \] (52)

Then, the classical scalar field \( \eta(x) \), governed by Eq. (40), would describe small fluctuations about a uniform distribution of (dark) matter, and the universe would be quite similar to a superconductor.
REFERENCES

[1] G’tHooft, Nucl. Phys. B35, 167 (1971).
[2] V. Delgado, in preparation.
[3] J. Dreitlein, Phys. Rev. Lett. 33, 1243 (1974).
[4] S. Weinberg, Rev. Mod. Phys. 61, 1 (1989).