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The Standard Model in Its Other Phase

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Abstract

The standard model of particle physics is analyzed for the case of a Higgs potential not favoring spontaneous electroweak symmetry breaking to gain insight into the physics of the standard model. Electroweak breaking still takes place, and quarks and leptons still acquire masses but through bosonic technicolor. This "other" phase of the standard model exhibits interesting phenomena.

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I. Introduction

The Higgs potential in the standard model of particle physics is

$$V_{Higgs} = M_{\mathcal{H}}^2 \mathcal{H}^\dagger \mathcal{H} + \frac{\lambda_{\mathcal{H}}}{2} \left(\mathcal{H}^\dagger \mathcal{H} \right)^2 \quad , \tag{1.1}$$

where the broken phase, corresponding to $M_{\mathcal{H}}^2 < 0$, is chosen in nature. This mass parameter and coupling $\lambda_{\mathcal{H}}$ are arranged so that

$$\left\langle \mathcal{H}^0 \right\rangle_{broken} \approx \frac{246\, GeV}{\sqrt{2}} \approx 175\, GeV \quad , \tag{1.2}$$

in order to produce the experimental values of the masses of the W and Z.

In this work, we consider the standard model in the phase for which $M_{\mathcal{H}}^2 > 0$. This is more than an exercise. Considerable insight into the physics of the standard model can be gained by analyzing it in other regions of parameter space. Furthermore, at high temperatures such as those that occurred in the early universe, thermal effects added a positive contribution to the Higgs mass effectively reversing its sign. Thus, many of the results of our current work are relevant to particle physics at high temperatures and to early universe cosmology. In particular, it is a better approximation to use zero for the current algebra masses of leptons and quarks (including the top quark!) at temperatures above $200\, GeV$ than to use the values found in nature at zero temperature.

We refer to the usual case of $M_{\mathcal{H}}^2 < 0$ as the *broken phase*, and to $M_{\mathcal{H}}^2 > 0$ as the *non-breaking-Higgs-potential phase*. Actually, the non-breaking-Higgs-potential phase still undergoes electroweak breaking but dynamically. The breaking effects are quite small. In comparing the two phases, we fix a mass scale of QCD rather than fixing the strong interaction coupling constant g_s.

Our definition of the standard model is the $SU_c(3) \times SU_L(2) \times U_Y(1)$ gauge theory with three generations of quarks and leptons. It incorporates the Weinberg-Salam-Glashow electroweak model including a single Higgs doublet and its Yukawa couplings, even though this sector has not been experimentally confirmed. If no fundamental Higgs is discovered, then it is likely that the true electroweak symmetry breaking mechanism will be similar to that created by a Higgs field and that quark and lepton

mass generation will be mimicked by a phenomenological Higgs field and Yukawa interactions.

It turns out that the "standard model in its other phase" exhibits some interesting physics. Spontaneous dynamical symmetry breaking of the electroweak group $SU_L(2) \times U(1)$ takes place so that the W and Z still achieve masses, although the masses are much smaller than in the usual broken phase.[1] Quarks and leptons obtain tiny masses by a mechanism identical to that used in bosonic technicolor.[2] In effect, bosonic technicolor is almost realized in nature. Had the deconfinement temperature of the strong interactions been much higher, there would have been a period of the early universe dominated by bosonic technicolor!

We analyze the model in stages: We first ignore electroweak effects and Yukawa interactions. The electroweak contributions are incorporated next, and the Higgs field and its interactions are treated as a final perturbation.

Section II analyzes the global symmetries of pure QCD, how they are dynamically broken by the strong interactions and the way in which they are explicitly broken by electroweak effects. Dynamical symmetry breaking arises through quark condensates. We assume that $\langle \bar{q}_i q_i \rangle \neq 0$, so that axial currents are broken. A posteriori, we find this to be consistent: Had the wrong pattern of symmetry breaking been assumed, certain pseudo-Goldstone bosons would have acquired negative square masses.[3]

The spontaneously breaking of axial symmetries implies the existence of Goldstone bosons. The charged ones acquire small masses through electroweak interactions. In Section III, we compute the one-loop effects. One interesting result is a cancellation between photon and Z^0 contributions, so that charged Goldstone boson masses are an order of magnitude smaller than one might have expected.

The neutral Goldstone bosons obtain even smaller masses after masses for the quarks and leptons are generated by the bosonic-technicolor mechanism: When $\langle \bar{q}_i q_i \rangle \neq 0$, terms linear in the neutral Higgs field are produced, and it acquires a vacuum expectation value. This value expectation value, in turn, leads to quark and lepton masses in the same manner as in the Weinberg-Salam-Glashow model. However, because these effects are all tiny, quark and lepton masses are orders of magnitude smaller in the non-breaking-Higgs-potential phase than in the breaking phase. Section IV

calculates the masses.

In Section V, we compute the light hadron spectrum drawing on a variety of methods including lattice QCD, experimental data and the quark model. There is a rich spectrum of lightest baryons not only involving up and down quarks but the other four quarks too. As a result, there is an explosion in the number of possible nuclei. See Section VI. This section also discusses the particle physics, the atomic physics and the cosmology of the standard model in the non-breaking-Higgs-potential phase.

Here is a partial list of other results: At tree level, the Weinberg weak-mixing angle and the Cabibbo-Kobayashi-Maskawa matrix in the non-breaking-Higgs-potential phase are the same as in the breaking phase. The masses of the quarks and leptons in the non-breaking-Higgs-potential phase range from about a milli-electron Volt for the electron to several hundred electron Volts for the top quark. The masses of the lightest vector meson states and the spin 3/2 baryons remain about the same. There are, however, many more of these states. The masses of the lightest baryons are lower by about 3%, and the neutron is lighter than the proton.

II. The Pattern of Symmetry Breaking

This section analyzes the global symmetries of the standard model. Some symmetries are spontaneously broken dynamically by the strong interactions. Since the strong interactions are dominant, we shall initially ignore electroweak effects and the scalar Higgs field.

It is convenient to assemble the Dirac fields of the six flavors of quarks into one column vector Ψ as

$$\Psi = \begin{pmatrix} u \\ d \\ c \\ s \\ t \\ b \end{pmatrix} \quad . \tag{2.1}$$

The strong interactions have a global symmetry group given by $S_{strong} = SU_L(6) \times SU_R(6) \times U_V(1)$, where the $SU_L(6)$ left-handed currents are

$$J_L^{a\mu} = \bar{\Psi}\gamma^\mu \Lambda^a \frac{(1-\gamma_5)}{2}\Psi \quad , \tag{2.2}$$

the $SU_R(6)$ currents are

$$J_R^{a\mu} = \bar{\Psi}\gamma^\mu\Lambda^a\frac{(1+\gamma_5)}{2}\Psi \quad , \tag{2.3}$$

and the vector $U_V(1)$ current, which is proportional to baryon number, is

$$J_V^{U(1)} = 3J_{baryon\, number} = 3J_b = \bar{\Psi}\gamma^\mu\Psi \quad . \tag{2.4}$$

In eqs.(2.2) and (2.3), Λ^a is an element of the Lie algebra of the group $SU(6)$. The symmetry associated with the $U_A(1)$ current is absent due to the axial anomaly. The group S_{strong} is 71 dimensional.

It is well known that the strong interactions spontaneously break S_{strong} down to a subgroup H

$$SU_L(6) \times SU_R(6) \times U_V(1) \to SU_V(6) \times U_V(1) \tag{2.5}$$

due to the formation of quark condensates:

$$\left\langle \bar{\Psi}\Psi \right\rangle \neq 0 \quad , \tag{2.6}$$

so that the unbroken strong interaction group is $H = SU_V(6) \times U_V(1)$. The axial generators

$$J_A^{a\mu} = \bar{\Psi}\gamma^\mu\Lambda^a\gamma_5\Psi \tag{2.7}$$

are spontaneously broken, while the vector generators

$$J_V^{a\mu} = \bar{\Psi}\gamma^\mu\Lambda^a\Psi \tag{2.8}$$

and $J_V^{U(1)}$ survive. The corresponding broken charges are

$$Q_A^a = \int d^3x\bar{\Psi}(x)\gamma^0X^a\gamma_5\Psi(x) \quad , \tag{2.9}$$

where $X^a \in SU(6)$. Since there are 35 generators for $SU(6)$, there are 35 massless Goldstone bosons at this stage.

It is convenient to introduce a compact notation for the currents. They all may be expressed as

$$J_\Sigma^\mu = \bar{\Psi}\gamma^\mu\Sigma\Psi \quad , \tag{2.10}$$

where Σ is of the form

$$\Sigma \in U_{U-D}\,(2) \times U_G\,(3) \cdot \Gamma \quad , \tag{2.11}$$

where U_{U-D} is the unitary group that acts on up/down type quarks, U_G is the group that acts on the three generations, and Γ is a linear combination of γ_5 and 1_4 (1_4 is the 4×4 unit matrix in Dirac space). For example, the baryon number current J_b is associated with

$$\Sigma_b = \frac{1}{3} I_2 \times I_3 \cdot 1_4 \quad , \tag{2.12}$$

where I_n is the $n \times n$ identity matrix.

It is convenient to partition S_{strong} into six classes, which we label A, B, C, D, E and F and which we now define.

Introduce the $SU_L\,(2) \times U_Y\,(1)$ electroweak gauge interactions. The vector gauge bosons of $SU_L\,(2)$ couple to a current associated with

$$\Sigma_{SU_L(2)} = \tau^a \times I_3 \cdot \frac{(1-\gamma_5)}{2} \quad \text{(class C)} \quad , \tag{2.13}$$

while that of $U_Y\,(1)$ couples to

$$\Sigma_{U_Y(1)} = \begin{pmatrix} \frac{1}{3} & 0 \\ 0 & \frac{1}{3} \end{pmatrix} \times I_3 \cdot \frac{(1-\gamma_5)}{2} + \begin{pmatrix} \frac{4}{3} & 0 \\ 0 & -\frac{2}{3} \end{pmatrix} \times I_3 \cdot \frac{(1+\gamma_5)}{2} \quad . \tag{2.14}$$

As is well known, the electromagnetic current is a linear combination of $U_Y\,(1)$ and the "τ^3" generator of $SU_L\,(2)$:

$$\Sigma_{U_{EM}} = \begin{pmatrix} \frac{2}{3} & 0 \\ 0 & -\frac{1}{3} \end{pmatrix} \times I_3 \cdot 1_4 \quad . \tag{2.15}$$

Some of the symmetries of S_{strong} are also symmetries of the electroweak interactions. It is straightforward to find those that preserve $SU_L\,(2) \times U_Y\,(1)$. With the exception of the gauge generators, these electroweak preserving symmetries are associated with the Σ that commute with $\Sigma_{SU_L(2)}$ and $\Sigma_{U_Y(1)}$ in eqs.(2.13) and (2.14). The electroweak global symmetry group S_W has 29 elements:

$$\tau^a \times I_3 \cdot \frac{(1-\gamma_5)}{2}\ , \ a = 1,\ 2,\ or\ 3\ , \quad I_2 \times \lambda^\alpha \cdot \frac{(1-\gamma_5)}{2} \quad ,$$

$$\tau^3 \times I_3 \cdot \frac{(1+\gamma_5)}{2}\ , \quad \tau^3 \times \lambda^\alpha \cdot \frac{(1+\gamma_5)}{2}\ , \quad I_2 \times \lambda^\alpha \cdot \frac{(1+\gamma_5)}{2} \quad , \tag{2.16}$$

$$I_2 \times I_3 \cdot 1_4 \quad ,$$

where $\lambda^\alpha \in SU_G(3)$. Of these 29 generators, four are associated with the global symmeties G_W of the electroweak gauge group and are given in eqs.(2.13) and (2.14).

There are 42 generators orthogonal to those in eq.(2.16) that correspond to currents explicitly broken by electroweak gauge interactions:

$$\begin{gathered} \tau^a \times \lambda^\alpha \cdot \frac{(1-\gamma_5)}{2} \quad , \\ \tau^1 \times \lambda^\alpha \cdot \frac{(1+\gamma_5)}{2} \; , \quad \tau^2 \times \lambda^\alpha \cdot \frac{(1+\gamma_5)}{2} \quad , \\ \tau^1 \times I_3 \cdot \frac{(1+\gamma_5)}{2} \; , \quad \tau^2 \times I_3 \cdot \frac{(1+\gamma_5)}{2} \quad . \end{gathered} \tag{2.17}$$

It is convenient to re-organize these currents according to parity. The even parity vector currents

$$\tau^a \times \lambda^\alpha \cdot 1_4 \; , \quad \tau^1 \times I_3 \cdot 1_4 \; , \quad \tau^2 \times I_3 \cdot 1_4 \quad \text{(class F)} \tag{2.18}$$

are not spontaneously broken by the strong interactions, while the odd parity axial currents

$$\tau^1 \times \lambda^\alpha \cdot \gamma_5 \; , \quad \tau^2 \times \lambda^\alpha \cdot \gamma_5 \quad \text{(class E)} \tag{2.19}$$

are spontaneously broken.

The 29 symmetries of S_W can be divided into four classes. Class A consists of electroweak gauge currents that are not spontaneously broken by the strong interactions. There is only one member and it is the electromagnetic current:

$$\begin{pmatrix} \frac{2}{3} & 0 \\ 0 & -\frac{1}{3} \end{pmatrix} \times I_3 \cdot 1_4 \leftrightarrow U_{EM} \quad \text{(class A)} \quad . \tag{2.20}$$

There are nine non-gauge currents that are not broken by either the strong interactions or the electroweak interactions. They are

$$I_2 \times I_3 \cdot 1_4 \leftrightarrow U_V(1) \; , \quad I_2 \times \lambda^\alpha \cdot 1_4 \leftrightarrow SU_G(3) \quad \text{(class B)} \quad , \tag{2.21}$$

corresponding to baryon number and $SU_G(3)$, the group that acts on the three-dimensional generation space. Particle multiplets can thus be classified according to charge, baryon number and the representation of $SU_G(3)$.

Of the remaining 19 global symmetries of the electroweak interactions, three correspond to gauge symmetries that are spontaneously broken by the strong interactions. They constitute Class C and are given in eq.(2.13). The broken currents $\tau^a \times I_3 \cdot \gamma_5$ can be expressed as a linear combination of these three gauge currents (eq.(2.13)) and currents conserved by the strong interactions. In such a case, it is well-known[4] that the Goldstone bosons corresponding to the broken charges $Q_A^a = \int d^3x \bar{\Psi}(x)\gamma^0 \tau^a \times I_3 \gamma_5 \Psi(x)$ are "eaten" by electroweak gauge bosons to become massive vector particles via the Higgs mechanism. Weinberg calls these types of Goldstone bosons fictitious because they do not appear as scalar particles in the physical spectrum. Of course, this is dynamical gauge symmetry breaking,[5, 6, 7, 4] which is the basis for technicolor.[1, 8]

It turns out that the same linear combinations of neutral electroweak gauge bosons acquire masses as in the case of the usual broken standard model, and that the Weinberg mixing angle θ_w is the same: $\sin(\theta_W) = \frac{g'}{\sqrt{g_2^2+g'^2}}$.[1, 8] Thus, the pattern of electroweak breaking is identical for non-breaking and breaking Higgs-potential phases. The main difference is that the masses of the W and Z are smaller, the scale being set by the strong interactions:

$$M_W = \sqrt{n_d}\,\frac{g_2}{2} f_\pi \approx 50\,MeV\ , \quad M_Z = \frac{M_W}{\cos(\theta_W)} \approx 57\,MeV \quad , \tag{2.22}$$

where $n_d = 3$ is the number of electroweak doublet condensates. Here, g_2 is the $SU_L(2)$ gauge coupling (see eq.(3.1) below) and the pion decay constant is normalized as $f_\pi \approx 93\,MeV$. The electroweak ρ_{EW} parameter $\frac{M_W^2}{M_Z^2\cos^2(\theta_W)}$ is thus 1 at tree level. When loop corrections are taken into account, ρ_{EW} in the non-breaking-Higgs-potential phase will differ slightly from ρ_{EW} in the breaking phase.

We have accounted for all the currents except for 16 corresponding to symmetries of S_W that are not global versions of electroweak gauge interactions and that are broken by the strong interactions. They are

$$I_2 \times \lambda^\alpha \cdot \gamma_5\ , \quad \tau^3 \times \lambda^\alpha \cdot \frac{(1+\gamma_5)}{2} \quad . \tag{2.23}$$

The currents

$$I_2 \times \lambda^\alpha \cdot \gamma_5\ , \quad \tau^3 \times \lambda^\alpha \cdot \gamma_5 \quad \text{(class D)} \tag{2.24}$$

can be written as a linear combination of the currents in eq.(2.23) and currents not broken by the strong interactions. It follows that the electrically neutral bosons associated with the spontaneously broken charges $\int d^3x\bar{\Psi}(x)\gamma^0 I_2 \times \lambda^\alpha \gamma_5 \Psi(x)$ and $\int d^3x\bar{\Psi}(x)\gamma^0\tau^3 \times \lambda^\alpha \cdot \gamma_5 \Psi(x)$ do not achieve masses through electroweak interactions. At this stage, they are massless Goldstone bosons, or true Goldstone bosons in the language of Weinberg.[4]

Figure 1 summarizes the partitioning of S_{strong} into the six classes. Classes A and B are the global versions of the electroweak gauge symmetries with class B being spontaneously broken by the strong interactions. The union of classes A, B, C and D constitutes the global symmetry group S_W of the electroweak interactions, classes C and D being non-electroweak gauge symmetries. Class C differs from class D in that it is not spontaneously broken by the strong interactions. Classes A, B and F compose H, the unbroken strong interaction group. Members of F are not symmetries of the electroweak gauge interactions. The complement of H, that is, the union of classes C, D and E are symmetries spontaneously broken by the strong interactions. There are potentially Goldstone bosons associated with these three classes. Class E differs from classes C and D in that these symmetries are not respected by the electroweak gauge interactions. The bosons associated with class C are fictitious in that they are eaten by the electroweak vector gauge bosons through the Higgs mechanism. The bosons of class E are pseudo Goldstone because the achieve masses through the electroweak gauge interactions. Finally, class D is associated with true Goldstone bosons, particles with zero mass when Higgs-Yukawa interactions are neglected.

III. The Goldstone Boson Spectrum

Of the 35 Goldstone bosons, three are eaten by the Z and W, the 16 electrical neural ones are massless and the 16 charged ones acquire masses from electroweak interactions. The purpose of this section is to compute the common mass of this charged octet of $SU_G(3)$ using current algebra. A useful formalism for this calculation has been developed in refs.[9, 10]. We shall review just enough of this formalism to set notation and to render the text readable.

The electroweak gauge bosons interact with quarks through the following la-

grangian

$$\mathcal{L}_{EW} = \frac{g_2}{2}\sum_{a=1}^{3} A^a_\mu J^\mu_{\Sigma^a_{SU_L(2)}} + \frac{g'}{2} B_\mu J^\mu_{\Sigma_{U_Y(1)}} \quad , \tag{3.1}$$

where A^a_μ and B_μ are respectively the gauge bosons for the $SU_L(2)$ and $U_Y\,(1)$ gauge groups, which couple to their corresponding currents $J^\mu_{\Sigma^a_{SU_L(2)}}$ and $J^\mu_{\Sigma_{U_Y(1)}}$ (given in eqs.(2.10), (2.13) and (2.14)) with strengths $g_2/2$ and $g'/2$. Following ref.[9], we rewrite $\mathcal{L}_{EW}$ in terms of left-handed Dirac fields as

$$\mathcal{L}_{EW} = \sum_{\Xi} g_\Xi A^\Xi_\mu J^\mu_\Xi \quad , \tag{3.2}$$

by using the charge conjugates of right-handed fields. Let ψ be a column vector of 12 entries with the first six given by the left-handed components of Ψ and the second six given by the conjugates of the right-handed components of Ψ. The vector fields A^Ξ_μ in eq.(3.2) couple to the currents

$$J^\mu_\Xi = \bar{\psi}\gamma^\mu \Xi \frac{(1-\gamma_5)}{2}\psi \quad , \tag{3.3}$$

where Ξ is a matrix in flavor space. There are four terms in eq.(3.2) corresponding to the four vector gauge bosons of the electroweak interactions. It is convenient to choose these gauge bosons to be in a mass diagonal basis, that is, to be γ, $W^\pm$ and Z^0.

The next step in the computation is to write

$$\Xi = \Xi_T + \Xi_X \quad , \tag{3.4}$$

where $\Xi_T \in H$ is a generator not broken by the strong interactions, and $\Xi_X \in S_{strong}/H$ is a broken generator.

Define the Goldstone decay constant f^a_Π by

$$\langle\Omega|\; J^{a\mu}_A |\Pi^b\,\rangle = i p^\mu f^a_\Pi \delta^{ab} \quad , \tag{3.5}$$

where Ω is the vacuum state and $|\Pi^b\,\rangle$ is the Goldstone boson of momentum p^μ associated with the broken charge Q^b_A. The current algebra $SU(6)$ matrices in $J^{a\mu}_A$ and Q^b_A of eqs.(2.7) and (2.9) are normalized so that

$$Tr\left(\Lambda^a\Lambda^b\right) = \delta^{ab}\ , \quad Tr\left(X^a X^b\right) = \delta^{ab} \quad . \tag{3.6}$$

The contribution to the Goldstone boson mass matrix m_{ab} from a symmetry breaking perturbation $\delta\mathcal{H}$ is given by Dashen's theorem[11]

$$m_{ab}^2 = -\frac{1}{f_\Pi^a f_\Pi^b} \langle\Omega| \left[Q_A^a, \left[Q_A^b, \delta\mathcal{H}\right]\right] |\Omega\rangle \quad . \tag{3.7}$$

In the present case, $\delta\mathcal{H}$ arises in second order perturbation theory through the one-loop exchange of electroweak vector gauge bosons. Figure 2 shows the diagram. The shaded region represents all possible QCD interactions. These involve exchanges of colored gluons and sea quark loops. Thus, the computation is non-perturbative in the strong interactions.

Neglecting electroweak interactions, one has

$$f_\Pi^a = f_\Pi \quad , \tag{3.8}$$

and eq.(3.8) is expected to hold quite well even when non-strong-interaction perturbations are included.

Using the above formalism, one is able to separate out the group theory factors in eq.(3.7) via [9, 10]

$$m_{ab}^2 = \frac{1}{4\pi} \sum_\Xi g_\Xi^2 \Delta^\Xi Tr \left(\left[\Xi_T, \left[\Xi_T, \Lambda^a\right]\right] \Lambda^b - \left[\Xi_X, \left[\Xi_X, \Lambda^a\right]\right] \Lambda^b \right) \quad , \tag{3.9}$$

where Δ^Ξ, a parameter of dimension mass squared, involves the effects of the propagator $\Delta_{\mu\nu}^\Xi(x)$ of the vector boson A_μ^Ξ in the following time-ordered expectation:

$$\Delta^\Xi = \frac{4\pi}{f_\Pi^2} \int d^4x \Delta_{\mu\nu}^\Xi(x) \langle\Omega| T\left(J_V^\mu(x) J_V^\nu(0) - J_A^\mu(x) J_A^\nu(0)\right) |\Omega\rangle \quad . \tag{3.10}$$

Here, J_V^μ and J_A^μ are any normalized vector and axial currents. For example, one can take them to be the ones appearing in the third generator of $SU_L(2)$:

$$J_V^\mu = \frac{1}{\sqrt{6}} \bar\Psi \gamma^\mu \tau^3 \times I_3 \Psi \ , \quad J_A^\mu = \frac{1}{\sqrt{6}} \bar\Psi \gamma^\mu \tau^3 \times I_3 \gamma_5 \Psi \quad . \tag{3.11}$$

It turns out that, of the four contributions in the sum over Ξ in eq.(3.9), the two associated with $W^\pm$ vanish. The contribution from the photon to a positively charge Goldstone boson Π^+ is

$$m_{\Pi^+}^2 = \alpha_{EM} \Delta^\gamma \quad . \tag{3.12}$$

It remains to determine the non-perturbative parameter Δ^γ. It has been accurately estimated analytically in ref.[12], but it may be computed using the experimental spectrum of the pion:

$$m_{\pi^+}^2 - m_{\pi^0}^2 = \alpha_{EM}\Delta^\gamma \quad \Rightarrow \quad \Delta^\gamma \approx (415\, MeV)^2 \quad . \tag{3.13}$$

Using eqs.(3.12) and (3.13), one finds that the photon contribution to m_{Π^+} is approximately 35 MeV. This is what one would expect: m_{Π^+} should be of order of $\sqrt{\alpha}\Lambda_{QCD}$, where Λ_{QCD} is a QCD scale, which we take to be 300 MeV.

However, it turns out that the contribution of Z^0 cancels most of that of the photon. The final result from electroweak interactions is

$$m_{\Pi^+}^2 = \alpha_{EM}\left(\Delta^\gamma - \Delta^Z\right) \quad . \tag{3.14}$$

Since the mass of Z^0 is small compared to Λ_{QCD}, Δ^γ and Δ^Z of eq.(3.14) are almost equal. In fact, the difference vanishes as M_Z^2. Thus, m_{Π^+} should be of order $\sqrt{\alpha}M_Z$, or about 5 MeV, an order of magnitude smaller than without the cancellation.

A similar cancellation[9, 10] occurs in technicolor models for the pseudo Goldstone boson often refered to as P.[13] Unfortunately, there is no known non-perturbative way to compute Δ^Z. However, eq.(3.14) can be calculated in perturbation theory. One finds

$$m_{\Pi^+}^2 \approx \frac{3\alpha_{EM}}{4\pi} M_Z^2 \ln\left(\frac{\Lambda_{QCD}^2}{M_Z^2}\right) \sim (4\, MeV)^2 \quad . \tag{3.15}$$

In summary, the eight charged pseudo Goldstone bosons, the analog of the pions of the standard model, obtain masses of about 4 MeV through the electroweak interactions.

IV. The Effects of the Higgs Sector

No bare quark masses have been generated by the electroweak gauge interactions, so that quarks are massless at this stage. When the Higgs sector along with its Yukawa couplings are included, then tiny quark and lepton masses arise through a mechanism identical to that of bosonic technicolor.[2]

The coupling of the $SU_L(2)$ Higgs double $\mathcal{H} = \begin{pmatrix} \mathcal{H}^+ \\ \mathcal{H}^0 \end{pmatrix}$ to quarks and leptons in

the standard model is

$$\mathcal{L}_{Yukawa} = \sum_{i,j=1}^{3} \left(\lambda_{ij}^{D} \bar{Q}_L^i \mathcal{H} D_R^j + \lambda_{ij}^{U} \bar{Q}_L^i \tilde{\mathcal{H}} U_R^j + \lambda_{ij}^{E} \bar{L}_L^i \mathcal{H} E_R^j \right) + \text{c.c.} \quad , \qquad (4.1)$$

where i and j denote different generations; λ_{ij}^D, λ_{ij}^U and λ_{ij}^E are Yukawa coupling constants; Q_L^i and L_L^i are the left-handed $SU_L(2)$ doublets of quarks and leptons for the ith generation; D_R^j, U_R^j and E_R^j are right-handed $SU_L(2)$ singlets of down quarks, up quarks and charged leptons for the jth generation; and $\tilde{\mathcal{H}} = \begin{pmatrix} \mathcal{H}^{0*} \\ -\mathcal{H}^- \end{pmatrix}$ is the $SU_L(2)$ conjugate Higgs field.

When $\bar{q}q$ quark condenstates form, the Yukawa interactions generate linear terms for $\mathcal{H}$, thereby causing the Higgs field to acquire a vacuum expectation value: The neutral component of $\mathcal{H}$ must be shifted to eliminate the linear terms.

Although the bosons absorbed by the Z^0 and $W^\pm$ now involve a tiny add-mixture of Higgs, the Higgs survives, and there are two neutral bosons and one charged boson of mass of $\sim M_{\mathcal{H}}$.

When the vacuum expectation of $\mathcal{H}^0$ is substituted into $\mathcal{L}_{Yukawa}$, the quarks and charged leptons obtain masses. The same unitary transformations on matter fields as in the broken phase diagonalize mass matrices. It follows that, at tree-level, the Cabibbo-Kobayashi-Maskawa matrix in the non-breaking-Higgs-potential phase is the same as in the broken phase. The mixing-angle effects enter the charged gauge weak interactions but do not affect the pseudo Goldstone boson computations in Section III, because the $W^\pm$ do not contribute, and even if they did, the mixing angles are small.

The calculation of $\langle \mathcal{H}^0 \rangle$ is straightfoward. The shift is quite small so that the quartic term in V_{Higgs} may be neglected. One finds

$$M_{\mathcal{H}}^2 \left\langle \mathcal{H}^0 \right\rangle \approx \sum_{k=1}^{6} \lambda_{q_i} \left\langle \bar{q}_i q_i \right\rangle \quad , \qquad (4.2)$$

where λ_{q_i} are the diagonalized Yukawa couplings for quarks ($m_{q_i} = \lambda_{q_i} \langle \mathcal{H}^0 \rangle$). Although $\langle \bar{q}_i q_i \rangle$ is known, $\langle \bar{q}_i q_i \rangle \approx (225\, MeV)^3$, the mass of the Higgs field is not. Therefore, only an order of magnitude estimate for $\langle \mathcal{H}^0 \rangle$ and quark and lepton masses can be made. Assuming $100\, GeV \leq M_{\mathcal{H}} \leq 300\, GeV$, one finds

$$\left\langle \mathcal{H}^0 \right\rangle \sim 100\, eV\, to\, 1\, KeV \quad . \qquad (4.3)$$

If follows that

$$\frac{\langle \mathcal{H}^0 \rangle}{\langle \mathcal{H}^0 \rangle_{broken}} \sim 10^{-9} \quad , \tag{4.4}$$

so that quark and lepton masses are about 10^9 times smaller in the non-breaking-Higgs-potential phase as in the broken phase.

For example, taking $M_{\mathcal{H}} = 175\, GeV$, which gives $\langle \mathcal{H}^0 \rangle / \langle \mathcal{H}^0 \rangle_{broken} \sim 2 \times 10^{-9}$, one obtains

$$\begin{gathered} m_t \sim 400\, eV \ , \quad m_c \sim 3\, eV \ , \quad m_u \sim 0.01\, eV \quad , \\ m_b \sim 10\, eV \ , \quad m_s \sim 0.4\, eV \ , \quad m_d \sim 0.02\, eV \quad , \\ m_\tau \sim 4\, eV \ , \quad m_\mu \sim 0.25\, eV \ , \quad m_e \sim 0.001\, eV \quad . \end{gathered} \tag{4.5}$$

The neutral pseudo Goldstone boson spectrum is the order of $\sqrt{\langle \mathcal{H}^0 \rangle / \langle \mathcal{H}^0 \rangle_{broken}}$ ($\sim 5 \times 10^{-5}$) times the spectrum of the unbroken model. With $M_{\mathcal{H}} = 175\, GeV$, one finds

$$\begin{gathered} \text{Those bosons with a t quark}: \ \sim 0.7\, MeV \quad (5 \text{ states}) \quad , \\ \text{Those bosons with a c quark but no t quark}: \ \sim 0.07\, MeV \quad (3 \text{ states}) \quad , \\ \text{Those bosons with a b quark}: \ \sim 0.14\, MeV \quad (5 \text{ states}) \quad , \\ \text{Those bosons with an s quark but no b quark}: \ \sim 0.025\, MeV \quad (3 \text{ states}) \quad . \end{gathered} \tag{4.6}$$

The neutral state involving only u and d quarks is mostly eaten by the Z^0, leaving no light pion;

V. The Rest of the Light Hadron Spectrum

In this section, we determine the spectrum of the lighter non-Goldstone-boson hadrons using a combination of methods that include the quark model, lattice QCD and experimental data.

There is one remaining pseudo scalar corresponding to the would-be Goldstone boson associated with $\bar{\Psi}\gamma^\mu\gamma_5\Psi$. It receives its mass m_A from topological fluctuations:[14]

$$m_A^2 = 2L\frac{\left\langle Q_{top}^2 \right\rangle}{f_\pi^2} \quad , \tag{5.1}$$

where L is the number of light quarks and $\left\langle Q_{top}^2 \right\rangle$ is the $SU_c(3)$ topological susceptibility, which can be numerically determined from experimental data as $\left\langle Q_{top}^2 \right\rangle \approx (180\ MeV)^2$. There is also a contribution to m_A^2 from topological fluctuations in the $SU_L(2)$ sector but it is considerably smaller than the $SU_c(3)$ term. Using eq.(5.1), one finds

$$m_A^2 \approx 1210\ MeV \quad . \tag{5.2}$$

This mass is considerably larger than that of the η' in the broken phase because L is 6 rather than 2 to 3.

The hadron spectrum has an approximate $SU_f(6)$ flavor symmetry due to the 6 quarks with masses much less than the scale of QCD. The vector mesons consist of 36 states transforming as an adjoint $\underline{35}$ and singlet of $SU_f(6)$. They all have approximately the same mass m_V, which can be determined as follows. In the quark model, the splitting between scalar and vector states is due to the color hyperfine interaction between pairs of quarks:

$$\Delta H_{color\ hyperfine} = -\frac{2\pi\alpha_s}{3m_1 m_2}\delta^3(r)\,\lambda_1 \cdot \lambda_2 S_1 \cdot S_2 \quad , \tag{5.3}$$

where m_i, q_i and $S_i = \sigma_i/2$ are the mass, charge and spin of the ith quark. A spin independent meson mass M_0 can be determined by removing these spin-spin interaction effects:

$$M_0 = \frac{3}{4}m_V + \frac{1}{4}m_S \quad , \tag{5.4}$$

where m_S is the mass of the pseudo scalar mesons. Using experimental data, one finds $M_0 \approx 610\,MeV$. In the non-breaking-Higgs-potential phase, the quarks are lighter and M_0 should be smaller by about $15\,MeV$. On the other hand, the color hyperfine interaction is slightly enhanced since the masses m_1 and m_2 in eq.(5.3) are smaller by about $10\,MeV$. Taking these two effects into account, we find, for the vector meson mass m_V,

$$m_V \approx 790\ MeV \quad , \tag{5.5}$$

a value slightly larger than the mass of the ρ and ω.

The ligher baryons consist of a spin-3/2 $\underline{56}$ of $SU_f(6)$ containing the Δ and its partners, and a spin-1/2 $\underline{70}$ of $SU_f(6)$ containing the nucleon and its partners. These

two states are again split by the hyperfine interaction in eq.(5.3). Using similar methods as in the vector meson case, we find

$$M_{J=3/2} \approx 1217\ MeV\ , \quad M_{J=1/2} \approx 909\ MeV \quad . \tag{5.6}$$

The mass of the $J = 3/2$ baryons is slightly less than the experimental mass of the Δ, while the mass of the $J = 1/2$ baryons is somewhat less than the mass of the proton.

Of particular interest is the mass of the lightest baryon since it will be stable. Small mass differences among the $J = 1/2$ baryons are generated by the electroweak interactions. These interactions have an $SU_G(3)$ symmetry. Under $SU_G(3)$, the $\underline{70}$ of $SU_f(6)$ decomposes into an octet with electric charge $+2$; a singlet, two octets and a decaplet with charge $+1$; a singlet, two octets and a decaplet with charge 0; and an octet with electric charge -1. The two $+1$ charged octets are distinguished by their properties under interchange of the two up-type quarks: one is symmetric while the other is antisymmetric. Likewise, the two neutral octets are distinguished by their properties under the interchange of the two down-type quarks.

In the quark model in its broken phase, there are three contributions to baryonic mass differences: (a) the up/down quark mass difference, (b) the electrostatic potential energy, which may be estimated from the pairwise interaction of quarks via

$$\Delta H_{static\ potential} = q_1 q_2 \left\langle \frac{1}{r_{12}} \right\rangle \quad , \tag{5.7}$$

where q_i is the charge on the ith quark and where $\left\langle \frac{1}{r_{12}} \right\rangle$ is the average inverse distance between the two quarks, and (c) the electric hyperfine interaction of

$$\Delta H_{hf} = -\frac{q_1 q_2}{3 m_1 m_2} \delta^3 \left(r_{12}\right) S_1 \cdot S_2 \quad . \tag{5.8}$$

The baryonic mass differences can be fairly reliably computed in the quark model. For example, ref.[15] obtained the above three contributions using lattice QCD to compute $\left\langle \frac{1}{r_{12}} \right\rangle$ and $\psi\left(r_{12} = 0\right)$, which enters in eq.(5.8) when baryonic wave functions ψ are used. Reference [15] finds the following results

$Particles$	$Static\,Term$	$u-d\,Mass\,Term$	$Total$	$Experiment$
$p-n$	0.51	-2.21	-1.69	-1.29
$\Sigma^{+}-\Sigma^{0}$	0.07	-2.89	-2.82	-3.23
$\Sigma^{0}-\Sigma^{-}$	-1.62	-2.89	-4.51	-4.81
$\Xi^{0}-\Xi^{-}$	-1.76	-3.47	-5.23	-6.4 ± 0.6
$\Sigma^{*+}-\Sigma^{*0}$	0.41	-2.05	-1.64	-0.9 ± 1.1
$\Sigma^{*+}-\Sigma^{*-}$	-0.86	-4.11	-4.97	-4.4 ± 0.7
$\Xi^{*0}-\Xi^{*-}$	-1.29	-2.32	-3.60	-3.2 ± 0.7

Table 1 Baryonic Mass Splittings for the Broken Phase (in MeV)

It should be noted that the above calculation involved no adjustable parameters because input parameters had been fixed in an earlier computation of the meson spectrum. By comparing theory with experiment, Table 1 provides an estimate of the accuracy of our methods. In particular, the sign of all mass differences is correctly reproduced.

The baryonic mass splittings for the non-breaking-Higgs-potential phase proceed as in the broken phase with the following differences: there is no significant contribution from item (a) above because up-type and down-type quarks have almost identical masses, and there are contributions from the Z^0 and $W^{\pm}$. Because the baryons involve equal admixtures of left and right quarks, formulas (5.7) and (5.8) hold for Z^0 exchange with $q \to (q_L + q_R)/2$. The $W^{\pm}$ only contributes to baryons with both up-type and down-type quarks and interchanges the two types, which must be taken into account. We have computed the effects of the exchange of the Z^0 and $W^{\pm}$ using the baryonic wave functions of ref.[15] and incorporating a suppression factor of $exp(-M_V/\Lambda_{QCD}) \approx 0.83$ due to the non-zero value of the mass M_V of these vector gauge bosons.

Here are the tabulated contributions to the $J = 1/2$ baryonic mass splittings.

$State$	EM_{sp}	EM_{hf}	Z_{sp}	Z_{hf}	W_{sp}	W_{hf}	$Total$
uuu 8	2.4	0.4	0.24	0.04	0	0	3.07
uud 10	0	−0.4	−0.20	−0.11	−0.81	−0.27	−1.79
uud 8_+	0	−0.4	−0.20	−0.11	0.40	0.13	−0.17
uud 8_-	0	0.4	−0.20	0.04	0.40	0	0.64
uud 1	0	0.4	−0.20	0.04	−0.81	0	−0.57
ddu 1	−0.6	0.1	−0.03	0.13	−0.81	0	−1.20
ddu 8_+	−0.6	0.1	−0.03	0.13	0.40	0	0.00
ddu 8_-	−0.6	−0.3	−0.03	−0.14	0.40	0.13	−0.53
ddu 10	−0.6	−0.3	−0.03	−0.14	−0.81	−0.27	−2.14
ddd 8	0.6	0.1	0.76	0.13	0	0	1.58

Table 2 Electroweak Contributions to Baryonic Masses (in MeV)

The octet containing the neutron is the lightest state weighing about one-third of an MeV less than the octet containing the proton. This result is important for the discusion of the nuclear and atomic physics of the non-breaking-Higgs-potential phase.

VI. The Physics of the Standard Model in Its Other Phase

A. Nuclear Physics: A Rich World with Lots of Isotopes

The nucleons of the non-breaking-Higgs-potential phase can have four charges: +2, +1, 0 and −1. The +2 and −1 octets probably do not form stable nucleonic bound states because they are heavier than the lightest nucleon by more than $3.5\,MeV$. Of the +1-charged and neutral nucleons, the decaplets are the lowest-mass states. We call +1-charged and neutral nucleons *protons* and *neutrons* even though there are now many types of each, some of which are comprised of strange, charm, bottom and top quarks.

Compared to the usual standard model, the non-breaking-Higgs-potential phase has four main differences as far as nuclear physics is concerned: (i) the neutron is stable and the proton is slightly heavier than the neutron, (ii) the Yukawa forces between nucleons is stronger and longer ranged because the pseudo-Goldstone bosons are lighter, (iii) there are a decaplet of lightest protons p_{10} and a decaplet of lightest neutrons n_{10} (which we denote simply by p and n), and (iv) besides the lightest

protons and neutrons, there are other nucleons, the least massive of which is the $SU_G(3)$ singlet neutron n_1.

These four differences lead to an extraodinarily rich spectrum of nuclei. Using charge to distinguish nucleonic elements, there are an infinite number of elements, each of which has an infinite number of isotopes! The lightest $Q = 0$ nucleus is the spin-1/2 decaplet of neutrons. Two neutrons also form a bound state, the lowest-energy state being a spin-1 $\underline{45}$ of $SU_G(3)$; the spin-orbit interaction favors the formation of a spin-1 state over a spin-0 state. The lightest state of three neutrons is a spin-3/2 $\underline{120}$ of $SU_G(3)$. It is guaranteed to be stable because ^3He in nature is stable, intra-nucleonic forces are stronger, and the *n*-*n*-*n* does not have the repulsive Coulombic energy of ^3He. Although slightly heavier than its spin-3/2 counterpart, the spin-1/2 *n*-*n*-*n* transforming as a $\underline{330}$ of $SU_G(3)$ is stable. It is clear that any number of neutrons can join to form a nucleus. Thus, there are an infinite number of $Q = 0$ isotopes. The same is true for $Q > 0$ nuclei: any number of neutrons can be added to a stable nucleus to produce stable isotopes. It is therefore possible to form giant nucleonic nuggets constisting of countless neutrons. Roughly speaking, they would be similar to neutron stars but would have a greater range of sizes, anything varying from the microscopic to the macroscopic: nuclei of hundreds of neutrons, neutron nuggets, neutron balls and neutron planets.

The $Q = 1$ proton is not stable; it decays into a neutron, a neutrino and a positively charged lepton. Its lifetime can be estimated to be about 10^{-7} seconds. However, the "deuteron" consisting of a proton and a neutron is stable. It is a spin-1 $\underline{45}$ of $SU_G(3)$. Among the "tritium" *n*-*n*-*p* states, the $\underline{450}$ spin-3/2 nuclei are lightest. There are also stable spin-1/2 *n*-*n*-*p* states transforming as $\underline{450}$ and $\underline{550}$ of $SU_G(3)$. In the $\underline{450}$, the neutron spins are "aligned," whereas in the $\underline{550}$ they are "oppositely" paired. The spin-1 *n*-*n* unit is a particularly stable building block that can be added to any nucleus to form an isotope.

The $\underline{450}$ spin-3/2 "^{3}He" *n*-*p*-*p* states are probably stable: On one hand, the intra-nucleonic forces are stronger; On the other hand, the protons are heavier thereby leading to less stability. However, "^{4}He" is guaranteed to exist. The lightests states are spin 2. It is a $45 \times 45 = \underline{2025}$ of $SU_G(3)$. In addition, there are spin-0 and

spin-1 multiplets of $\underline{2025}$ in which both the two protons have aligned spins and the two neutrons have aligned spins. There are also two spin-1 $\underline{2475}$'s in which two like-nucleons have aligned spins while the other two like-nucleons have oppositely aligned spins. Finally, there are 3025 spin-0 states in which the two protons have oppositely paired spins and the two neutrons have oppositely paired spins.

As in the usual standard model, the "^{4}He" units form particularly stable building blocks for nucleons. Thus, the most stable lighter nuclei are formed from "^{4}He" and n-n units. Examples are the nuclei ^6He, ^8He, ... , ^{12}C, ^{14}C, ^{16}C, ... , ^{16}O, ^{18}O, ^{20}O, ... , and ^{20}Ne, ^{22}Ne, ^{24}Ne, ... Needless to say, these nuclei appear in huge multiplets.

It is also likely that the singlet neutron n_1 whose mass is about $0.94\,MeV$ heavier than the decaplet neutron can participate in nuclear bound states. In isolation, it decays to a charged lepton, an antineutrino and a proton. However, it may be that n_1 and p_{10} bind to form 10 heavy deuteron states. If the usual deuteron is considered to be "heavy hydrogen," then n_1-p_{10} would be heavy "heavy hydrogen."

B. Stronger Weak Interactions

In terms of particle physics, the most pronounced difference[1] between the breaking and the non-breaking Higgs-potential phases is the strength of the weak interactions: G_F is about 2.5×10^6 times bigger due to the smaller value of the W mass. This leads to an enhancement of about 6.7×10^{12} in matrix elements squared. This is the reason why the proton is so short-lived, decaying in about 10^{-7} seconds. On the other hand, the phase space for many weak processes is greatly reduced or smaller mass parameters enter in decay rates; This is actually the dominant effective for many light states. For example, the muon can decay into an electron, a muon neutrino and an anti-electron neutrino. However, the muon's lifetime is about 2×10^{17} years! The tau is also long-lived, lasting 1.5×10^{11} years and decaying into a muon, a tau neutrino and an anti-muon neutrino. For the charged pseudo scalar bosons, the above two

[1]There are many smaller differences. For example, the strong interaction coupling constant runs more slowly in the non-breaking-Higgs-potential phase because there are more light quarks. Also, CP violation is likely to be less in this phase because the magnitude of the entries in the qaurk mass matrix are smaller, thereby implying that the phases in the Cabibbo-Kobayashi-Maskawa matrix have less effect.

effects roughly cancel. The lifetime of these particles is about 5×10^{-8} seconds. They decay into a tau and a tau neutrino.

Because the weak interactions are stronger, parity-violating effects will be more pronounced in the non-breaking-Higgs-potential phase. Such effects show up as small admixtures of opposite parity in nuclei and atoms.

C. Atomic Physics: Giant Atoms

Because of the electron's low mass of a milli-electron Volt, it would be very difficult for charged nuclei to capture electrons to form atoms. If this were to happen, the binding energy would be about $10^{-8}\,eV$ and the atom would be a dozen centimeters large.

However, since tau's and muons are long-lived, they can bind to nuclei. Atoms constructed with tau's replacing electrons would be about a dozen microns big and have binding energies of about a milli-electron Volt. Atoms constructed with muons would be about ten times larger and have binding energies of about $10^{-4}\,eV$. It would take a long time for the universe to sufficiently cool to allow these types of atoms.

D. Cosmology: Plasma Domination, Exotic Leptons and Anti-Leptons

There are a lot of interesting details concerning the cosmology of the standard model in its other phase. However, in this subsection, we focus on some general features.

Up to a trillionth of a second, there are no essential changes because high temperatures prevent electroweak breaking. In the non-breaking-Higgs-potential phase, the $SU_L(2) \times U_Y(1)$ breaking occurs when confinement sets in at about a millionth of a second when the temperature T of the universe is around $200\,MeV$. Neutrino decoupling, which usually occurs around 1 second ($T \sim 1\,MeV$), now takes place when $T \sim 50\,eV$ because the weak interactions are so much stronger. Thus, the generation of light nuclei (of the type described in subsection A above) takes place under equilibrium conditions. The ratio of the number of protons to the number of neutrons in nuclei will therefore be slightly less than one (in the usual standard model it is 7 to 1). Big Bang Nucleosynethesis also takes place in the presence of tau's, muons and electrons and their anti-particles and in the presence of many neutral mesons because these states are all light. It also happens a little earlier: from a fraction of a second

to one minute because the universe expands more quickly and cools more rapdily due to the presence of these additional light particles.

Recombination only takes place after the universe has become frigid. When $T \sim 10^{-7}\, eV$, corresponding to when the universe is 10^{15} years old, muons bind to charged nuclei. The tau leptons actually decay before having a chance to form atoms. Electrons bind to atoms even later.

When the universe reaches one billion years old, galaxy and star formation should begin when clouds of nuclei and charged leptons collapse. The main difference compared to the usual standard model is that the universe remains in a plasma for billions of years. Also noteworthy is the presence of positions, muons, anti-muons, tau's and anti-tau's – particles that are absent in the case of the usual standard model. It would be interesting to investigate the detailed evolution of such a universe, but this topic is beyond the goals of the present work.

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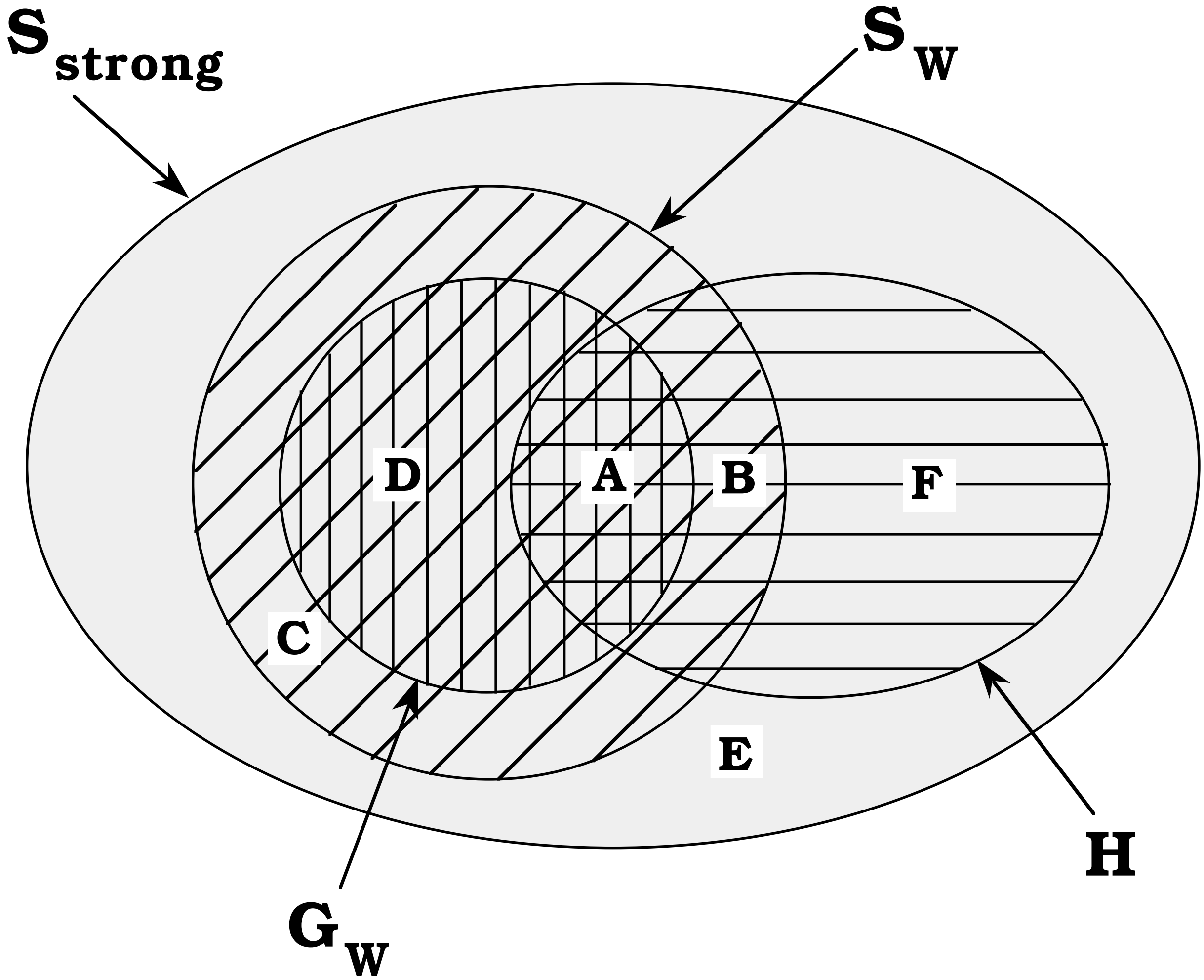

Figure 1. The Partition of Global Symmetries into Different Classes

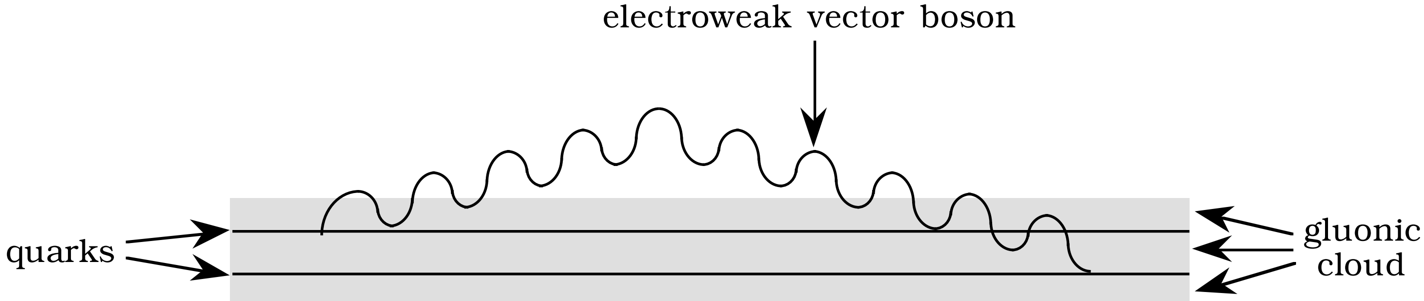

Figure 2. The One-Loop Electroweak Contribution to Pseudo-Goldstone Boson Masses