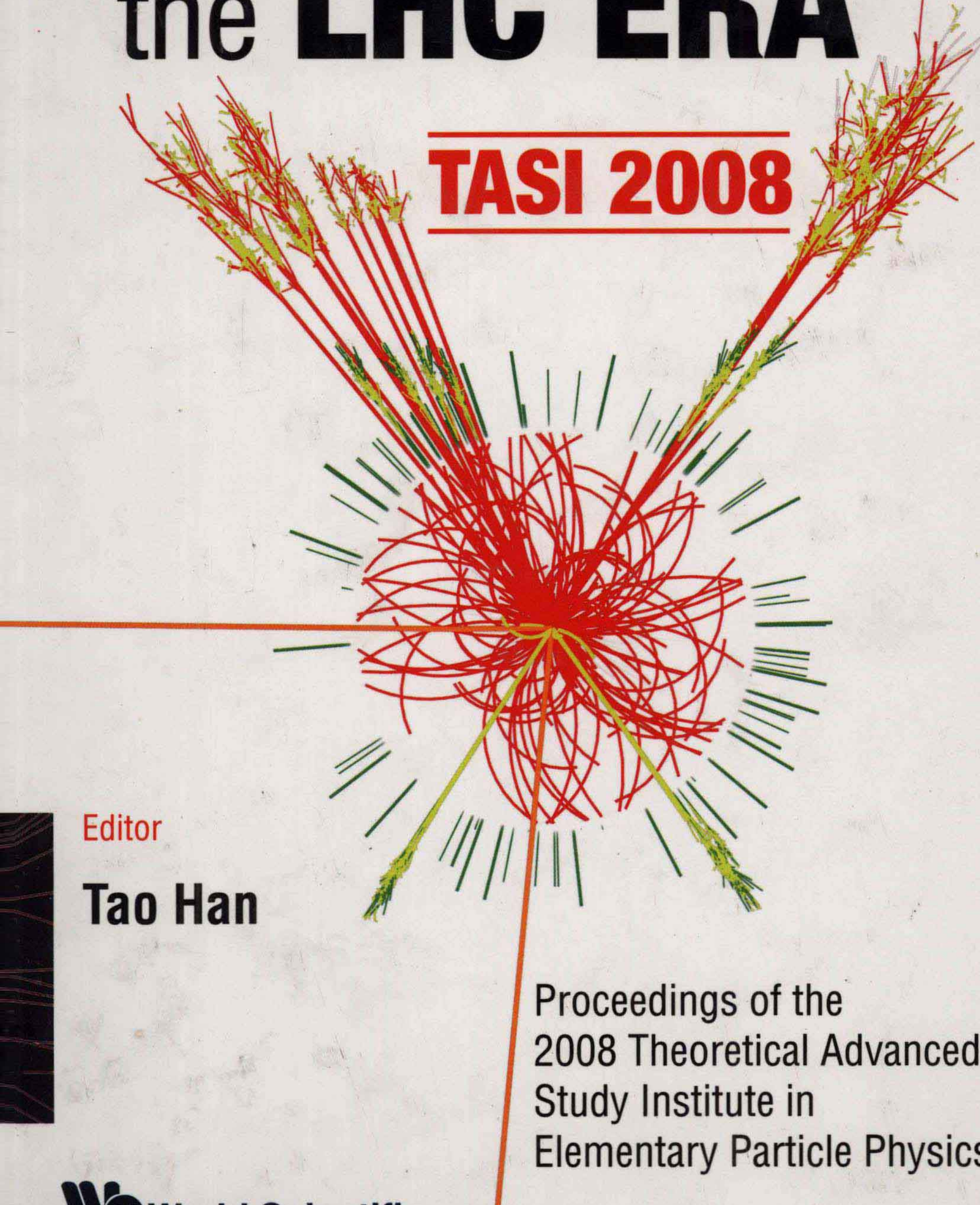
The DAWN of the LHC ERA

TASI 2008



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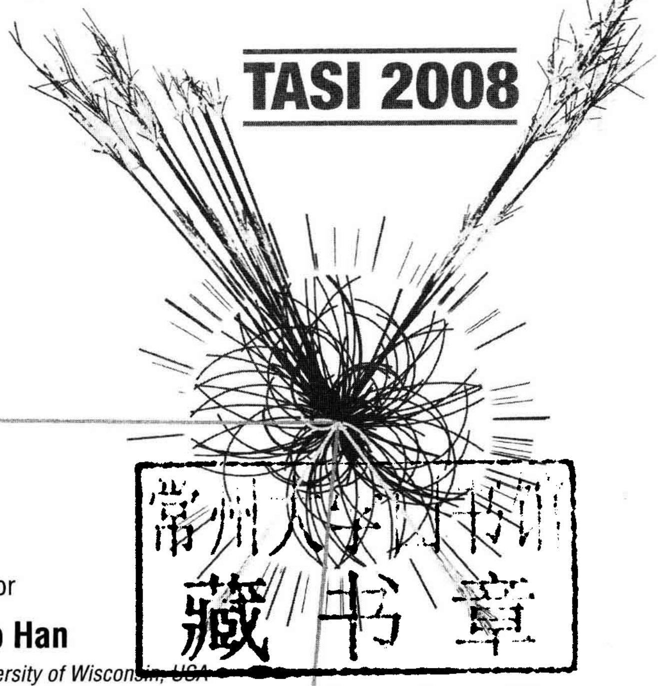

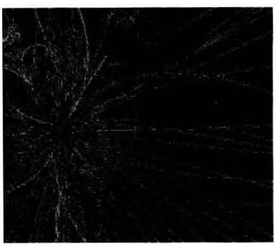

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Tao Han

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Preface

When this volume of the Proceedings is published, the CERN Large Hadron Collider (LHC) will have started taking data, and a new era in high energy physics, as well as in science overall, has since begun. The depth of the LHC program in understanding the nature at such short space-time scales, its extent in exploring fundamental questions, the scope of the involvement in the community world-wide are all unprecedented. Revolutionary scientific progress is highly anticipated from the LHC experiments. The LHC will fully explore the Tera-scale physics and possibly beyond: the mechanism of the electroweak symmetry breaking, fermion masses, their flavor mixing and CP violation, new forces and strongly interacting dynamics, the nature of dark matter, larger unification of fundamental interactions, extended symmetries such as Supersymmetry and extra spatial dimensions, to name a few exciting possibilities.

In the recent years, we have also witnessed the major discoveries in the fields of cosmology and neutrino physics. The nature of dark energy, the origin of dark matter, and the cause of the inflation are all mysteries to uncover. The tiny neutrino masses, their large flavor mixing, and the nature of Dirac or Majorana type are all pressing issues to address.

The design of the TASI-2008 lectures reflects these upcoming developments in the field. There are four parts of the lectures:

- The Standard Model And LHC Phenomenology
- LHC Experimentation
- Advanced Theoretical Topics
- Neutrino Physics, Astroparticle Physics, And Cosmology

I feel deeply impressed by the balance of the contents, the coverage of the materials, the pedagogical nature of the presentation and the writing of the lecture notes, for which I am grateful to those dedicated lecturers. I believe that the students in TASI 08 must have learnt a lot from them. The Proceedings will also benefit other researchers in the related fields.

In closing, I would like to thank the TASI general Director, Prof. K.T. Mahanthappa for the enjoyable collaboration during TASI 08; Susan Spika and Elizabeth Price for their efficient secretarial help in making TASI 08 run smoothly; Tom DeGrand for organizing the mountain hikes that challenged our limitations and imagination; and my co-scientific Director Robin Erbacker for helping me putting the program together. TASI Schools thank the National Science Foundation, the Department of Energy and the University of Colorado for financial and material support. I would also like to thank the TASI Scientific Advisor Board to have invited me as one of the Scientific Program Directors, that gave me the opportunity to have worked with these leading physicists as lecturers, and also gave me the pleasure to have interacted with these motivated and talented young students, who are our future in taking high energy physics to a brand new stage.

Go LHC!

Tao Han
University of Wisconsin, Madison

Contents

PART 1
The Standard Model and LHC Phenomenology

Chapter 1

Introduction to the Standard Model and Electroweak Physics

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A concise introduction is given to the standard model, including the structure of the QCD and electroweak Lagrangians, spontaneous symmetry breaking, experimental tests, and problems.

1.1. The Standard Model Lagrangian

1.1.1. *QCD*

The standard model (SM) is a gauge theory[1,2] of the microscopic interactions. The strong interaction part, quantum chromodynamics (QCD)* is an $SU(3)$ gauge theory described by the Lagrangian density

$$\mathcal{L}_{SU(3)} = -\frac{1}{4} F^i_{\mu\nu} F^{i\mu\nu} + \sum_r \bar{q}_{r\alpha} i \not{D}^\alpha_\beta \, q^\beta_r, \tag{1.1}$$

where g_s is the QCD gauge coupling constant,

$$F^i_{\mu\nu} = \partial_\mu G^i_\nu - \partial_\nu G^i_\mu - g_s f_{ijk} \, G^j_\mu \, G^k_\nu \tag{1.2}$$

is the field strength tensor for the gluon fields G^i_μ, $i = 1, \cdots, 8$, and the structure constants f_{ijk} $(i, j, k = 1, \cdots, 8)$ are defined by

$$[\lambda^i, \lambda^j] = 2i f_{ijk} \lambda^k, \tag{1.3}$$

where the $SU(3)$ λ matrices are defined in Table 1.1. The λ's are normalized by $\mathrm{Tr}\, \lambda^i \lambda^j = 2\delta^{ij}$, so that $\mathrm{Tr}\, [\lambda^i, \lambda^j] \lambda^k = 4i f_{ijk}$.

*See Ref. [3] for a historical overview. Some recent reviews include Ref. [4] and the QCD review in Ref. [5].

Table 1.1. The $SU(3)$ matrices.

$\lambda^i = \begin{pmatrix} \tau^i & 0 \\ 0 & 0 \end{pmatrix}, \quad i = 1,2,3$	
$\lambda^4 = \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix}$	$\lambda^5 = \begin{pmatrix} 0 & 0 & -i \\ 0 & 0 & 0 \\ i & 0 & 0 \end{pmatrix}$
$\lambda^6 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}$	$\lambda^7 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -i \\ 0 & i & 0 \end{pmatrix}$
$\lambda^8 = \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}$	

The F^2 term leads to three and four-point gluon self-interactions, shown schematically in Figure 1.1. The second term in $\mathcal{L}_{SU(3)}$ is the gauge covariant derivative for the quarks: q_r is the r^{th} quark flavor, $\alpha, \beta = 1, 2, 3$ are color indices, and

$$D^{\alpha}_{\mu\beta} = (D_\mu)_{\alpha\beta} = \partial_\mu \delta_{\alpha\beta} + i g_s \, G^i_\mu \, L^i_{\alpha\beta}, \tag{1.4}$$

where the quarks transform according to the triplet representation matrices $L^i = \lambda^i/2$. The color interactions are diagonal in the flavor indices, but in general change the quark colors. They are purely vector (parity conserving). There are no bare mass terms for the quarks in (1.1). These would be allowed by QCD alone, but are forbidden by the chiral symmetry of the electroweak part of the theory. The quark masses will be generated later by spontaneous symmetry breaking. There are in addition effective ghost and gauge-fixing terms which enter into the quantization of both the $SU(3)$ and electroweak Lagrangians, and there is the possibility of adding an (unwanted) term which violates CP invariance.

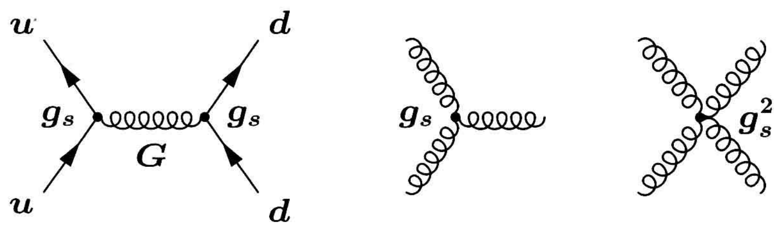

Fig. 1.1. Interactions in QCD.

QCD has the property of asymptotic freedom,[6,7] i.e., the running coupling becomes weak at high energies or short distances. It has been extensively tested in this regime, as is illustrated in Figure 1.2. At low energies or large distances it becomes strongly coupled (infrared slavery),[8] presumably leading to the confinement of quarks and gluons. QCD incorporates the observed global symmetries of the strong interactions, especially the spontaneously broken global $SU(3) \times SU(3)$ (see, e.g., Ref. [9]).

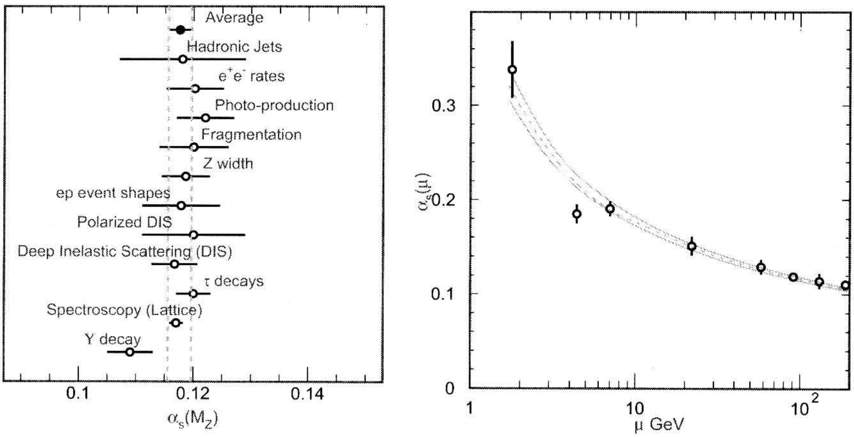

Fig. 1.2. Running of the QCD coupling $\alpha_s(\mu) = g_s(\mu)^2/4\pi$. Left: various experimental determinations extrapolated to $\mu = M_Z$ using QCD. Right: experimental values plotted at the μ at which they are measured. The band is the best fit QCD prediction. Plot courtesy of the Particle Data Group,[5] `http://pdg.lbl.gov/`.

1.1.2. *The Electroweak Theory*

The electroweak theory[10–12] is based on the $SU(2) \times U(1)$ Lagrangian†

$$\mathcal{L}_{SU(2)\times U(1)} = \mathcal{L}_{gauge} + \mathcal{L}_\phi + \mathcal{L}_f + \mathcal{L}_{Yuk}. \tag{1.5}$$

The gauge part is

$$\mathcal{L}_{gauge} = -\frac{1}{4} W^i_{\mu\nu} W^{\mu\nu i} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu}, \tag{1.6}$$

where W^i_μ, $i = 1,\ 2,\ 3$ and B_μ are respectively the $SU(2)$ and $U(1)$ gauge fields, with field strength tensors

$$\begin{aligned} B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu \\ W^i_{\mu\nu} &= \partial_\mu W^i_\nu - \partial_\nu W^i_\mu - g\epsilon_{ijk} W^j_\mu W^k_\nu, \end{aligned} \tag{1.7}$$

†For a recent discussion, see the electroweak review in Ref. [5].

where $g(g')$ is the $SU(2)$ ($U(1)$) gauge coupling and ϵ_{ijk} is the totally antisymmetric symbol. The $SU(2)$ fields have three and four-point self-interactions. B is a $U(1)$ field associated with the weak hypercharge $Y = Q - T^3$, where Q and T^3 are respectively the electric charge operator and the third component of weak $SU(2)$. (Their eigenvalues will be denoted by y, q, and t^3, respectively.) It has no self-interactions. The B and W_3 fields will eventually mix to form the photon and Z boson.

The scalar part of the Lagrangian is

$$\mathcal{L}_\phi = (D^\mu \phi)^\dagger D_\mu \phi - V(\phi), \tag{1.8}$$

where $\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}$ is a complex Higgs scalar, which is a doublet under $SU(2)$ with $U(1)$ charge $y_\phi = +\frac{1}{2}$. The gauge covariant derivative is

$$D_\mu \phi = \left(\partial_\mu + ig\frac{\tau^i}{2} W_\mu^i + \frac{ig'}{2} B_\mu \right) \phi, \tag{1.9}$$

where the τ^i are the Pauli matrices. The square of the covariant derivative leads to three and four-point interactions between the gauge and scalar fields.

$V(\phi)$ is the Higgs potential. The combination of $SU(2) \times U(1)$ invariance and renormalizability restricts V to the form

$$V(\phi) = +\mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2. \tag{1.10}$$

For $\mu^2 < 0$ there will be spontaneous symmetry breaking. The λ term describes a quartic self-interaction between the scalar fields. Vacuum stability requires $\lambda > 0$.

The fermion term is

$$\begin{aligned} \mathcal{L}_f = \sum_{m=1}^{F} \big(& \bar{q}_{mL}^0 i \not{D} q_{mL}^0 + \bar{l}_{mL}^0 i \not{D} l_{mL}^0 + \bar{u}_{mR}^0 i \not{D} u_{mR}^0 \\ & + \bar{d}_{mR}^0 i \not{D} d_{mR}^0 + \bar{e}_{mR}^0 i \not{D} e_{mR}^0 + \bar{\nu}_{mR}^0 i \not{D} \nu_{mR}^0 \big) . \end{aligned} \tag{1.11}$$

In (1.11) m is the family index, $F \geq 3$ is the number of families, and $L(R)$ refer to the left (right) chiral projections $\psi_{L(R)} \equiv (1 \mp \gamma_5)\psi/2$. The left-handed quarks and leptons

$$q_{mL}^0 = \begin{pmatrix} u_m^0 \\ d_m^0 \end{pmatrix}_L \qquad l_{mL}^0 = \begin{pmatrix} \nu_m^0 \\ e_m^{-0} \end{pmatrix}_L \tag{1.12}$$

transform as $SU(2)$ doublets, while the right-handed fields u_{mR}^0, d_{mR}^0, e_{mR}^{-0}, and ν_{mR}^0 are singlets. Their $U(1)$ charges are $y_{q_L} = \frac{1}{6}$, $y_{l_L} = -\frac{1}{2}$, $y_{\psi_R} = q_\psi$.

The superscript 0 refers to the weak eigenstates, i.e., fields transforming according to definite $SU(2)$ representations. They may be mixtures of mass eigenstates (flavors). The quark color indices $\alpha = r,\ g,\ b$ have been suppressed. The gauge covariant derivatives are

$$\begin{aligned} D_\mu q^0_{mL} &= \left(\partial_\mu + \tfrac{ig}{2}\vec{\tau}\cdot\vec{W}_\mu + \tfrac{ig'}{6}B_\mu\right) q^0_{mL} & D_\mu u^0_{mR} &= \left(\partial_\mu + \tfrac{2ig'}{3}B_\mu\right) u^0_{mR} \\ D_\mu l^0_{mL} &= \left(\partial_\mu + \tfrac{ig}{2}\vec{\tau}\cdot\vec{W}_\mu - \tfrac{ig'}{2}B_\mu\right) l^0_{mL} & D_\mu d^0_{mR} &= \left(\partial_\mu - \tfrac{ig'}{3}B_\mu\right) d^0_{mR} \\ & & D_\mu e^0_{mR} &= \left(\partial_\mu - ig'B_\mu\right) e^0_{mR} \\ & & D_\mu \nu^0_{mR} &= \partial_\mu \nu^0_{mR}, \end{aligned} \tag{1.13}$$

from which one can read off the gauge interactions between the W and B and the fermion fields. The different transformations of the L and R fields (i.e., the symmetry is chiral) is the origin of parity violation in the electroweak sector. The chiral symmetry also forbids any bare mass terms for the fermions. We have tentatively included $SU(2)$-singlet right-handed neutrinos ν^0_{mR} in (1.11), because they are required in many models for neutrino mass. However, they are not necessary for the consistency of the theory or for some models of neutrino mass, and it is not certain whether they exist or are part of the low-energy theory.

The standard model is anomaly free for the assumed fermion content. There are no $SU(3)^3$ anomalies because the quark assignment is non-chiral, and no $SU(2)^3$ anomalies because the representations are real. The $SU(2)^2Y$ and Y^3 anomalies cancel between the quarks and leptons in each family, by what appears to be an accident. The $SU(3)^2Y$ and Y anomalies cancel between the L and R fields, ultimately because the hypercharge assignments are made in such a way that $U(1)_Q$ will be non-chiral.

The last term in (1.5) is

$$\begin{aligned} \mathcal{L}_{Yuk} = - \sum_{m,n=1}^{F} & \left[\Gamma^u_{mn} \bar{q}^0_{mL} \tilde{\phi} u^0_{nR} + \Gamma^d_{mn} \bar{q}^0_{mL} \phi d^0_{nR} \right. \\ & \left. + \Gamma^e_{mn} \bar{l}^0_{mn} \phi e^0_{nR} + \Gamma^\nu_{mn} \bar{l}^0_{mL} \tilde{\phi} \nu^0_{nR} \right] + h.c., \end{aligned} \tag{1.14}$$

where the matrices Γ_{mn} describe the Yukawa couplings between the single Higgs doublet, ϕ, and the various flavors m and n of quarks and leptons. One needs representations of Higgs fields with $y = +\frac{1}{2}$ and $-\frac{1}{2}$ to give masses to the down quarks and electrons $(+\frac{1}{2})$, and to the up quarks and neutrinos $(-\frac{1}{2})$. The representation $\phi^\dagger$ has $y = -\frac{1}{2}$, but transforms as the 2^* rather than the 2. However, in $SU(2)$ the 2^* representation is related to the 2 by a similarity transformation, and $\tilde{\phi} \equiv i\tau^2\phi^\dagger = \begin{pmatrix} \phi^{0\dagger} \\ -\phi^- \end{pmatrix}$ transforms as a 2 with $y_{\tilde{\phi}} = -\frac{1}{2}$. All of the masses can therefore be generated with a single Higgs doublet if one makes use of both ϕ and $\tilde{\phi}$. The fact that the fundamental and its conjugate are equivalent does not generalize to higher unitary groups. Furthermore, in supersymmetric extensions of the standard model the supersymmetry forbids the use of a single Higgs doublet in both ways in the Lagrangian, and one must add a second Higgs doublet. Similar statements apply to most theories with an additional $U(1)'$ gauge factor, i.e., a heavy Z' boson.

1.2. Spontaneous Symmetry Breaking

Gauge invariance (and therefore renormalizability) does not allow mass terms in the Lagrangian for the gauge bosons or for chiral fermions. Massless gauge bosons are not acceptable for the weak interactions, which are known to be short-ranged. Hence, the gauge invariance must be broken spontaneously,[13–18] which preserves the renormalizability.[19–22] The idea is that the lowest energy (vacuum) state does not respect the gauge symmetry and induces effective masses for particles propagating through it.

Let us introduce the complex vector

$$v = \langle 0|\phi|0\rangle = \text{ constant}, \tag{1.15}$$

which has components that are the vacuum expectation values of the various complex scalar fields. v is determined by rewriting the Higgs potential as a function of v, $V(\phi) \to V(v)$, and choosing v such that V is minimized. That is, we interpret v as the lowest energy solution of the classical equation of motion[‡]. The quantum theory is obtained by considering fluctuations around this classical minimum, $\phi = v + \phi'$.

The single complex Higgs doublet in the standard model can be rewritten in a Hermitian basis as

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \begin{pmatrix} \frac{1}{\sqrt{2}}(\phi_1 - i\phi_2) \\ \frac{1}{\sqrt{2}}(\phi_3 - i\phi_4 \end{pmatrix}, \tag{1.16}$$

‡It suffices to consider constant v because any space or time dependence $\partial_\mu v$ would increase the energy of the solution. Also, one can take $\langle 0|A_\mu|0\rangle = 0$, because any non-zero vacuum value for a higher-spin field would violate Lorentz invariance. However, these extensions are involved in higher energy classical solutions (topological defects), such as monopoles, strings, domain walls, and textures.[23,24]

where $\phi_i = \phi_i^\dagger$ represent four Hermitian fields. In this new basis the Higgs potential becomes

$$V(\phi) = \frac{1}{2}\mu^2 \left(\sum_{i=1}^{4} \phi_i^2\right) + \frac{1}{4}\lambda \left(\sum_{i=1}^{4} \phi_i^2\right)^2, \tag{1.17}$$

which is clearly $O(4)$ invariant. Without loss of generality we can choose the axis in this four-dimensional space so that $\langle 0|\phi_i|0\rangle = 0, \quad i = 1, 2, 4$ and $\langle 0|\phi_3|0\rangle = \nu$. Thus,

$$V(\phi) \to V(v) = \frac{1}{2}\mu^2\nu^2 + \frac{1}{4}\lambda\nu^4, \tag{1.18}$$

which must be minimized with respect to ν. Two important cases are illustrated in Figure 1.3. For $\mu^2 > 0$ the minimum occurs at $\nu = 0$. That is, the vacuum is empty space and $SU(2) \times U(1)$ is unbroken at the minimum. On the other hand, for $\mu^2 < 0$ the $\nu = 0$ symmetric point is unstable, and the minimum occurs at some nonzero value of ν which breaks the $SU(2) \times U(1)$ symmetry. The point is found by requiring

$$V'(\nu) = \nu(\mu^2 + \lambda\nu^2) = 0, \tag{1.19}$$

which has the solution $\nu = \left(-\mu^2/\lambda\right)^{1/2}$ at the minimum. (The solution for $-\nu$ can also be transformed into this standard form by an appropriate $O(4)$ transformation.) The dividing point $\mu^2 = 0$ cannot be treated classically. It is necessary to consider the one loop corrections to the potential, in which case it is found that the symmetry is again spontaneously broken.[25]

We are interested in the case $\mu^2 < 0$, for which the Higgs doublet is replaced, in first approximation, by its classical value $\phi \to \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ \nu \end{pmatrix} \equiv v$. The generators L^1, L^2, and $L^3 - Y$ are spontaneously broken (e.g., $L^1 v \neq 0$). On the other hand, the vacuum carries no electric charge ($Qv = (L^3 + Y)v = 0$), so the $U(1)_Q$ of electromagnetism is not broken. Thus, the electroweak $SU(2) \times U(1)$ group is spontaneously broken to the $U(1)_Q$ subgroup, $SU(2) \times U(1)_Y \to U(1)_Q$.

To quantize around the classical vacuum, write $\phi = v + \phi'$, where ϕ' are quantum fields with zero vacuum expectation value. To display the physical particle content it is useful to rewrite the four Hermitian components of ϕ' in terms of a new set of variables using the Kibble transformation:[26]

$$\phi = \frac{1}{\sqrt{2}} e^{i \sum \xi^i L^i} \begin{pmatrix} 0 \\ \nu + H \end{pmatrix}. \tag{1.20}$$

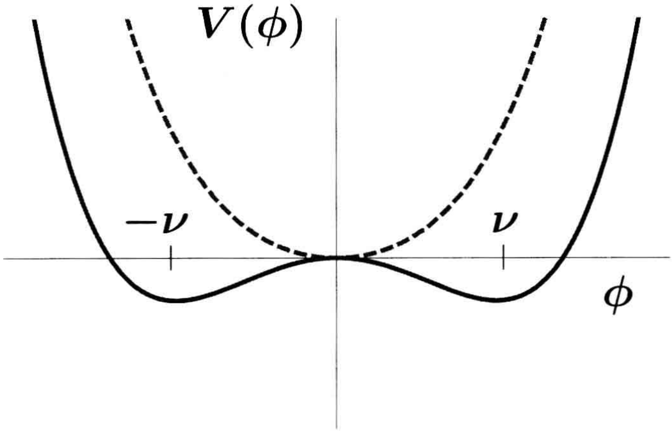

Fig. 1.3. The Higgs potential $V(\phi)$ for $\mu^2 > 0$ (dashed line) and $\mu^2 < 0$ (solid line).

H is a Hermitian field which will turn out to be the physical Higgs scalar. If we had been dealing with a spontaneously broken global symmetry the three Hermitian fields ξ^i would be the massless pseudoscalar Nambu-Goldstone bosons[27–30] that are necessarily associated with broken symmetry generators. However, in a gauge theory they disappear from the physical spectrum. To see this it is useful to go to the unitary gauge

$$\phi \to \phi' = e^{-i\sum \xi^i L^i}\phi = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ \nu + H \end{pmatrix}, \tag{1.21}$$

in which the Goldstone bosons disappear. In this gauge, the scalar covariant kinetic energy term takes the simple form

$$\begin{aligned}(D_\mu\phi)^\dagger D^\mu\phi &= \frac{1}{2}(0\ \nu)\left[\frac{g}{2}\tau^i W_\mu^i + \frac{g'}{2}B_\mu\right]^2\begin{pmatrix} 0 \\ \nu \end{pmatrix} + H \text{ terms} \\ &\to M_W^2 W^{+\mu}W_\mu^- + \frac{M_Z^2}{2}Z^\mu Z_\mu + H \text{ terms},\end{aligned} \tag{1.22}$$

where the kinetic energy and gauge interaction terms of the physical H particle have been omitted. Thus, spontaneous symmetry breaking generates mass terms for the W and Z gauge bosons

$$\begin{aligned}W^\pm &= \frac{1}{\sqrt{2}}(W^1 \mp iW^2) \\ Z &= -\sin\theta_W B + \cos\theta_W W^3.\end{aligned} \tag{1.23}$$