Reminiscences of Steven Weinberg and his Impact on High Energy Physics

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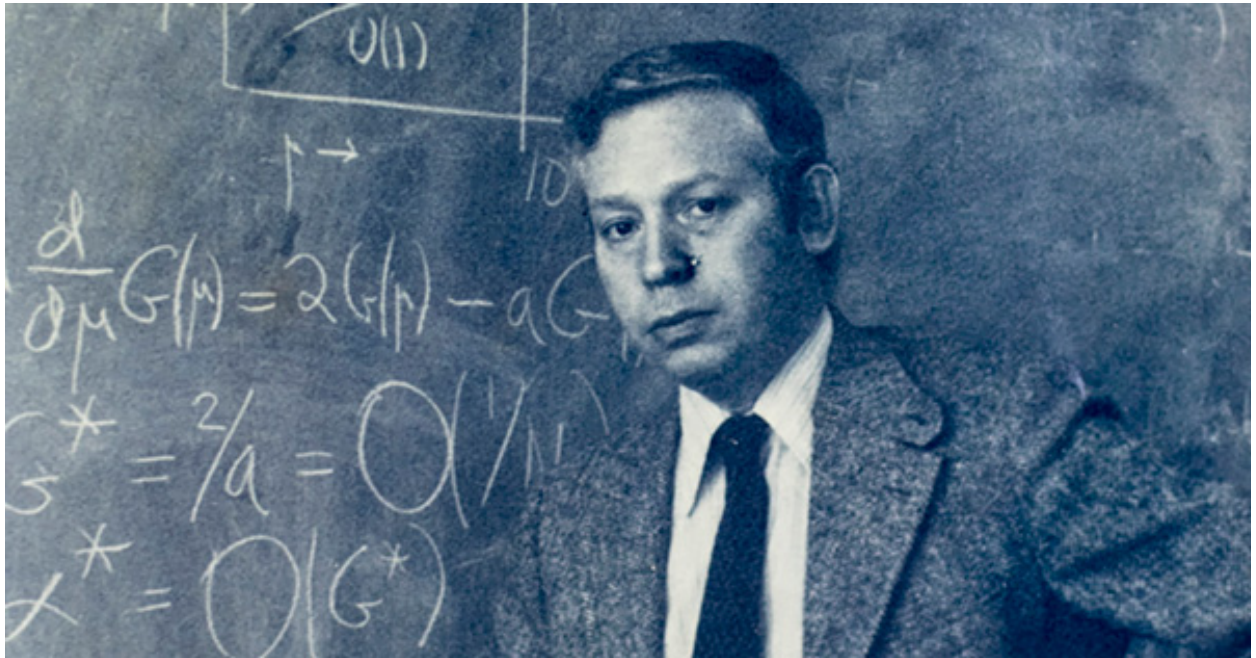
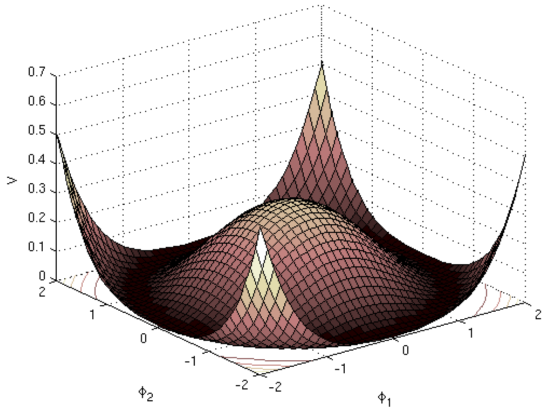
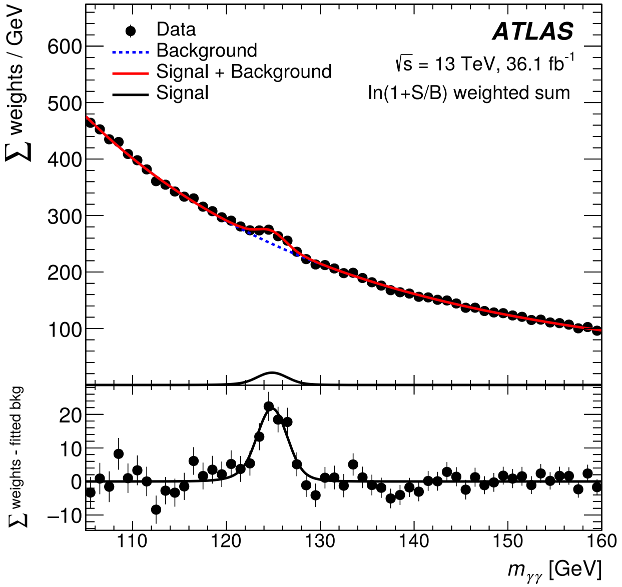


Standard Model

Name	Label	$\mathrm{SU}(3)_C,\ \mathrm{SU}(2)_L,\ \mathrm{U}(1)_Y$	Spin
Quarks	$Q_L^i = \begin{pmatrix} u_L^i \\ d_L^i \end{pmatrix}$	$(\mathbf{3}, \mathbf{2}, +\frac{1}{6})$	$\frac{1}{2}$
	u_R^i	$(\bar{\mathbf{3}}, \mathbf{1}, \frac{2}{3})$	$\frac{1}{2}$
	d_R^i	$(\bar{\mathbf{3}}, \mathbf{1}, -\frac{1}{3})$	$\frac{1}{2}$
Leptons	$L_L^i = \begin{pmatrix} \nu_L^i \\ e_L^i \end{pmatrix}$	$(\mathbf{1}, \mathbf{2}, -\frac{1}{2})$	$\frac{1}{2}$
	e_R^i	$(\mathbf{1}, \mathbf{1}, -1)$	$\frac{1}{2}$
	ν_R^{i*}	$(\mathbf{1}, \mathbf{1}, 0)$	$\frac{1}{2}$
Higgs	H	$(\mathbf{1}, \mathbf{2}, +\frac{1}{2})$	0
Gluons	g_α	$(\mathbf{8}, \mathbf{1}, 0)$	1
W/Z-Bosons	$W^\pm, Z^0$	$(\mathbf{1}, \mathbf{3}, 0)$	1
Photon	γ	$(\mathbf{1}, \mathbf{1}, 0)$	1
Graviton*	$h_{\mu\nu}$	$(\mathbf{1}, \mathbf{1}, 0)$	2

The Higgs





Weinberg's Scientific Production throughout 8 decades!

1950’s

PHYSICAL REVIEW VOLUME 118, NUMBER 3 MAY 1, 1960

High-Energy Behavior in Quantum Field Theory*

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(Received May 21, 1959)

An attack is made on the problem of determining the asymptotic behavior at high energies and momenta of the Green's functions of quantum field theory, using new mathematical methods from the theory of real variables. We define a class A_n of functions of n real variables, whose asymptotic behavior may be specified in a certain manner by means of certain "asymptotic coefficients." The Feynman integrands of perturbation theory (with energies taken imaginary) belong to such classes. We then prove that if certain conditions on the asymptotic coefficients are satisfied then an integral over k of the variables converges, and belongs to the class A_{n-k} with new asymptotic coefficients simply related to the old ones. When applied to perturbation theory this theorem validates the renormalization procedure of Dyson and Salam, proving that the renormalized integrals actually do always converge, and provides a simple rule for calculating the asymptotic behavior of any Green's function to any order of perturbation theory.

I. INTRODUCTION

IN many respects, the central formal problem of the modern quantum theory of fields is the determination of the asymptotic behavior at high energies and momenta of the Green's functions of the theory, the vacuum expectation values of time-ordered products. Complete knowledge of the asymptotic properties of these functions would allow us to test the renormalizability of a given Lagrangian, to count the number of subtractions that must be performed in dispersion theory, etc. We shall attack this problem from a rather new direction, which allows a solution in perturbation theory, and which provides an analytic tool that may prove useful in solving the problem in the exact theory.

One might hope to find a solution either kinematically, using only assumptions of covariance, causality, etc., or dynamically, by using the field equations that actually determine the Green's functions. The first method has been successfully applied to the 2-field functions, the particle propagators, and yields the result that the true propagators are asymptotically "larger" than the bare propagators.[1] However, because the theory of several complex variables is so difficult and incomplete, this approach seems unpromising for expectation values of three or more fields. For this reason, and also because we would eventually like to obtain renormalizability conditions on the Lagrangian, we propose to attack the problem on the dynamical level.

Now, what are the equations that, in principle, would determine the Green's functions. In perturbation theory we know that the Green's functions appear as multiple integrals, the integrand being constructed according to the Feynman rules. In a nonperturbative approach the Green's functions are again given by multiple integrals, but with integrands that themselves depend on the

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[1] H. Lehmann, Nuovo cimento **11**, 342 (1954).

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Charge Symmetry of Weak Interactions*

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The invariance of strong interactions under G, the product of charge symmetry and charge conjugation, has important consequences for strangeness-conserving lepton interactions. According to the G-transformation properties of the strongly interacting "currents," we may divide the primary weak interactions into two classes. The first class includes the conventional nucleon-lepton Fermi interaction, and is the only class that contributes to the β-decay coupling constants. Unambiguous tests for the existence of second-class interactions include: (a) induced scalar term in μ^- absorption, (b) inequality of certain small correction terms in B^{12} and N^{12}, or in Li^8 and B^8 β decay, (c) inequality in rates of $\Sigma^\pm \to \Lambda^0 + e^\pm + \nu$. Absence of second-class interactions would indicate a deep relation between isotopic spin and weak interactions; for example, the recent Feynman-Gell-Mann theory predicts that all vector weak interactions are first class. The presence of second-class interactions would mean that the usual Fermi interaction is insufficient, and must be supplemented by terms involving strange particles. Some general remarks are also made about the relations between $(l^-,\bar{\nu})$ and (l^+,ν) processes, and we prove the following useful theorem: no interference between V and A may occur in any experiment which treats both leptons identically and in which no parity nonconservation effects are measured, providing that we may neglect the mass and charge of the leptons.

I. INTRODUCTION

STRONG interactions are charge symmetric and charge conjugation invariant, and therefore also invariant under the product[1] G,

$$G \equiv C e^{i\pi I_2},$$
$$G\psi_N G^{-1} = i\tau_2 \psi_N, \quad G\phi_\pi G^{-1} = -\phi_\pi, \text{ etc.} \tag{1}$$

This G invariance plays a fundamental part in considering the effects of strong interactions on weak processes, and the rôle of isotopic spin in the primary weak interactions. We will show that all strangeness-conserving lepton interactions may be split into two

* This research was supported by the U. S. Atomic Energy Commission.
[1] T. D. Lee and C. N. Yang, Nuovo cimento **3**, 749 (1956). See also A. Pais and R. Jost, Phys. Rev. **87**, 871 (1952); L. Michel, Nuovo cimento **10**, 319 (1953), etc.

General theorems on Quantum Field Theory (QFT) and weak interactions

1960’s

Goldstone Bosons

PHYSICAL REVIEW VOLUME 127, NUMBER 3 AUGUST 1, 1962

Broken Symmetries*

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AND

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(Received March 16, 1962)

Some proofs are presented of Goldstone's conjecture, that if there is continuous symmetry transformation under which the Lagrangian is invariant, then either the vacuum state is also invariant under the transformation, or there must exist spinless particles of zero mass.

I. INTRODUCTION

IN the past few years several authors have developed an idea which might offer hope of understanding the broken symmetrics that seem to be characteristic of elementary particle physics. Perhaps the fundamental Lagrangian is invariant under all symmetries, but the vacuum state[1] is not. It would then be impossible to prove the usual sort of symmetry relations among S-matrix elements, but enough symmetry might remain (perhaps at high energy) to be interesting.

But whenever this idea has been applied to specific models, there has appeared an intractable difficulty. For example, Nambu suggested that the Lagrangian might be invariant under a continuous chirality transformation $\psi \rightarrow \exp(i\boldsymbol{\theta}\cdot\boldsymbol{\tau}\gamma_5)\psi$ even if the fermion physical mass M were nonzero. But then there would be a conserved current J_λ, with matrix element

$$\langle p'|J_\lambda|p\rangle = f(q^2)\bar{u}'\gamma_5[i\gamma_\lambda - (2M/q^2)q_\lambda]u,$$

where $q=p-p'$. The pole at $q^2=0$ can only arise from a spinless particle of mass zero, which almost certainly does not exist. Of course, the pole would not occur if $f(0)=0$, which might be the case if we do not insist on identifying J_λ with the axial vector current of β decay. But Nambu showed that this unwanted massless "pion" also appears as a solution of the approximate Bethe-Salpeter equation.[1]

Goldstone[2] has examined another model, in which the manifestation of "broken" symmetry was the nonzero vacuum expectation value of a boson field. (This was suggested as an explanation of the $\Delta I=\frac{1}{2}$ rule by Salam and Ward.)[3] Here again there appeared a spinless particle of zero mass. Goldstone was led to conjecture that this will always happen whenever a continuous symmetry group leaves the Lagrangian but not the vacuum invariant.

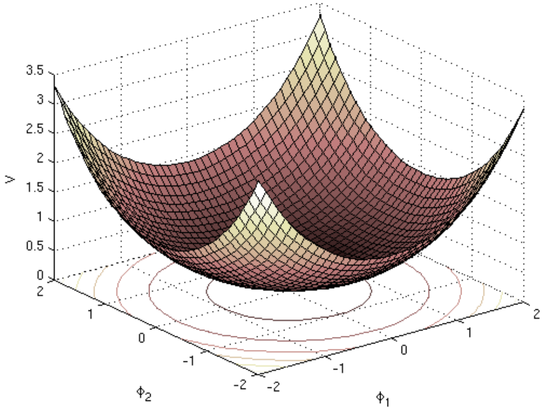
* This research was supported in part by the U. S. Air Force under a contract monitored by the Air Force Office of Scientific Research of the Air Development Command and the Office of Naval Research.

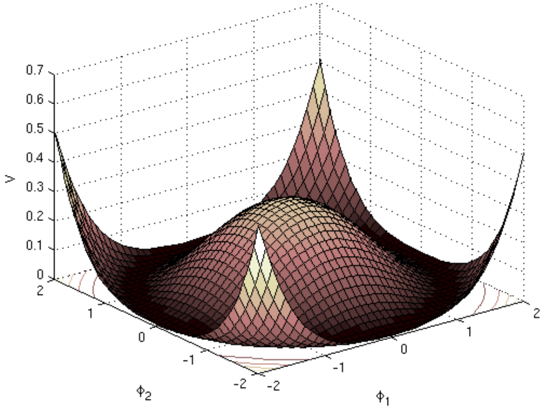
† Alfred P. Sloan Foundation Fellow; Permanent address: University of California, Berkeley, California.

[1] Y. Nambu and G. Jona-Lasinio, Phys. Rev. **122**, 345 (1961); W. Heisenberg, Z. Naturforsch. **14**, 441 (1959).

[2] J. Goldstone, Nuovo cimento **19**, 154 (1961).

[3] A. Salam and J. C. Ward, Phys. Rev. Letters **5**, 512 (1960).





PHYSICAL REVIEW VOLUME 135, NUMBER 4B 24 AUGUST 1964

Photons and Gravitons in S-Matrix Theory: Derivation of Charge Conservation and Equality of Gravitational and Inertial Mass*

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(Received 13 April 1964)

We give a purely S-matrix-theoretic proof of the conservation of charge (defined by the strength of soft photon interactions) and the equality of gravitational and inertial mass. Our only assumptions are the Lorentz invariance and pole structure of the S matrix, and the zero mass and spins 1 and 2 of the photon and graviton. We also prove that Lorentz invariance alone requires the S matrix for emission of a massless particle of arbitrary integer spin to satisfy a "mass-shell gauge invariance" condition, and we explain why there are no macroscopic fields corresponding to particles of spin 3 or higher.

I. INTRODUCTION

IT is not yet clear whether field theory will continue to play a role in particle physics, or whether it will ultimately be supplanted by a pure S-matrix theory. However, most physicists would probably agree that the place of local fields is nowhere so secure as in the theory of photons and gravitons, whose properties seem indissolubly linked with the space-time concepts of gauge invariance (of the second kind) and/or Einstein's equivalence principle.

The purpose of this article is to bring into question the need for field theory in understanding electromagnetism and gravitation. We shall show that there are no general properties of photons and gravitons, which *have* been explained by field theory, which cannot also be understood as consequences of the Lorentz invariance and pole structure of the S matrix for massless particles of spin 1 or 2.[1] We will also show why there can be no macroscopic fields whose quanta carry spin 3 or higher.

What are the special properties of the photon or graviton S matrix, which might be supposed to reflect specifically field-theoretic assumptions? Of course, the usual version of gauge invariance and the equivalence principle cannot even be stated, much less proved, in terms of the S matrix alone. (We decline to turn on external fields.) But there are two striking properties of the S matrix which *seem* to require the assumption of gauge invariance and the equivalence principle:

(1) The S matrix for emission of a photon or graviton can be written as the product of a polarization "vector" ϵ^μ or "tensor" $\epsilon^\mu\epsilon^\nu$ with a covariant vector or tensor amplitude, and it vanishes if any ϵ^μ is replaced by the photon or graviton momentum q^μ.

(2) Charge, defined dynamically by the strength of soft-photon interactions, is additively conserved in all reactions. Gravitational mass, defined by the strength of soft graviton interactions, is equal to inertial mass for all nonrelativistic particles (and is twice the total energy for relativistic or massless particles).

Property (1) is actually a straightforward consequence of the well-known[2,3] Lorentz transformation properties of massless particle states, and is proven in Sec. II for massless particles of arbitrary integer spin. (It has already been proven for photons by D. Zwanziger.[4])

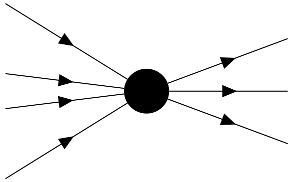
Property (2) does not at first sight appear to be derivable from property (1). Even in field theory (1) does not prove that the photon and graviton "currents" $J_\mu(x)$ and $\theta_{\mu\nu}(x)$ are conserved, but only that their matrix elements are conserved for light-like momentum transfer, so we cannot use the usual argument that $\int d^3x J^0(x)$ and $\int d^3x \theta^{0\mu}(x)$ are time-independent. And in pure S-matrix theory it is not even possible to define what we mean by the operators $J^\mu(x)$ and $\theta^{\mu\nu}(x)$.

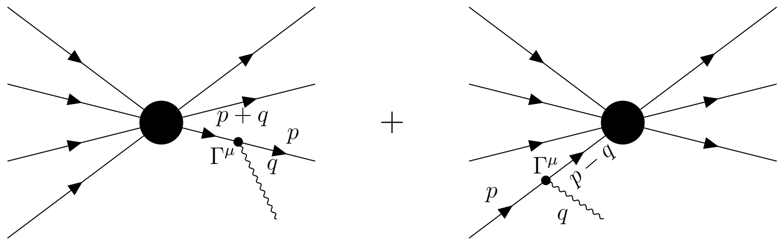
We overcome these obstacles by a trick, which replaces the operator calculus of field theory with a little simple polology. After defining charge and gravitational mass as soft photon and graviton coupling constants in Sec. III, we prove in Sec. IV that if a reaction violates charge conservation, then the same process with inner bremsstrahlung of a soft extra photon would have an S matrix which does not satisfy property (1), and hence would not be Lorentz invariant; similarly, the inner bremstrahlung of a soft graviton would violate Lorentz invariance if any particle taking part in the reaction has an anomalous ratio of gravitational to inertial mass.

Appendices A, B, and C are devoted to some technical problems: (A) the transformation properties of polarization vectors, (B) the construction of tensor amplitudes for massless particles of general integer spin, and (C) the presence of kinematic singularities in the conventional $(2j+1)$-component "M functions."

A word may be needed about our use of S-matrix theory for particles of zero mass. We do not know whether it will ever be possible to formulate S-matrix

* Research supported by the U. S. Air Force Office of Scientific Research, Grant No. AF-AFOSR-232-63.
† Alfred P. Sloan Foundation Fellow.
[1] Some of the material of this article was discussed briefly in a recent letter [S. Weinberg Phys. Letters **9**, 357 (1964)]. We will repeat a few points here, in order that the present article be completely self-contained.
[2] E. P. Wigner, in *Theoretical Physics* (International Atomic Energy Agency, Vienna, 1963), p. 59. We have repeated Wigner's work in Ref. 3.
[3] S. Weinberg, Phys. Rev. **134**, B882 (1964).
[4] D. Zwanziger, Phys. Rev. **113**, B1036 (1964). Zwanziger omits some straightforward details, which are presented here in Appendix B.

$\mathcal{M}_0 =$ 

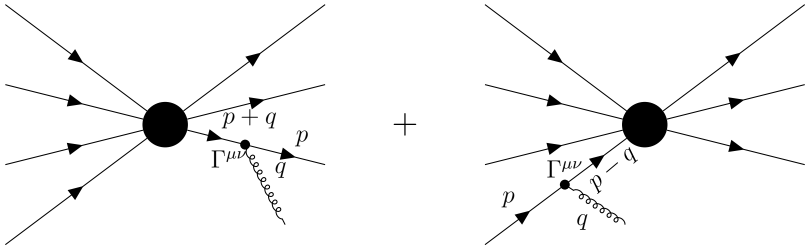
$\mathcal{M} =$ 

$$\Gamma^\mu = p^\mu F\left(\frac{p\cdot q}{m^2}\right)$$

$$\sum_{\text{incoming}} F_i(0) - \sum_{\text{outgoing}} F_i(0) = 0$$

$$\boxed{\sum_{\text{incoming}} Q_i = \sum_{\text{outgoing}} Q_i.}$$

Charge conservation!

$\mathcal{M} =$ 

$$\sum_{\text{incoming}} F_i(0)p_i^\nu - \sum_{\text{outgoing}} F_i(0)p_i^\nu = 0$$

$$\boxed{F_i(0) = \kappa \qquad \forall i}$$

Equivalence principle!

Weinberg QCD Sum rules

PRECISE RELATIONS BETWEEN THE SPECTRA OF VECTOR AND AXIAL-VECTOR MESONS

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(Received 3 January 1967)

Two sum rules are derived, relating moments of the spectral functions of the vector and axial-vector currents. If it is assumed that the ρ and $A1$ mesons dominate these moments, then their masses must be in the ratio $m_{A1}/m_\rho = \sqrt{2}$, in very good agreement with experiment.

If chiral $SU(2)\otimes SU(2)$ were an exact symmetry of the ordinary sort, we should expect the ρ meson to be accompanied with an $I=1$ axial-vector meson of the same mass. This is certainly not the case; the best candidate for the role of chiral partner of the ρ is the $A1$, which has $m_{A1}^2 \simeq 2m_\rho^2$. However, the recent successes of current algebra show that nature does obey some sort of chiral symmetry, manifested in the conservation or partial conservation of currents, and in their commutation relations. The question thus arises: What relations are imposed by current algebra upon the spectra of the 1^+ and 1^- mesons?

Our answer is contained in the following theorem: Assume that the vector and axial-vector currents obey the usual commutation relations,[1] with Schwinger terms[2] which are either c numbers or, if operators, contain no $\Delta I = 1$ terms. Neglect the pion mass altogether, so that the axial vector as well as the vector currents are conserved.[3] Then

$$\int_0^\infty [\rho_V(\mu^2) - \rho_A(\mu^2)]\mu^{-2}d\mu^2 = F_\pi^2, \tag{1}$$

where F_π is the usual pion-decay amplitude, and $\rho_{V,A}(\mu^2)$ are the spectral functions of the vector and axial-vector currents, defined by the formulas[4,5]

$$\langle V_a^\mu(x)V_b^\nu(0)\rangle_0 = (2\pi)^{-3}\delta_{ab}\int d^4p\,\theta(p^0)e^{ip\cdot x}\rho_V(-p^2)[g^{\mu\nu} - p^\mu p^\nu/p^2], \tag{2}$$

$$\langle A_a^\mu(x)A_b^\nu(0)\rangle_0 = (2\pi)^{-3}\delta_{ab}\int d^4p\,\theta(p^0)e^{ip\cdot x}\{\rho_A(-p^2)[g^{\mu\nu} - p^\mu p^\nu/p^2] + F_\pi^2\delta(p^2)p^\mu p^\nu\}. \tag{3}$$

If we further assume a very weak form of vector- and axial-vector-meson dominance, i.e., that matrix elements of the currents act at high momenta as if the currents were free $1^\pm$ fields,[6] then we also have

$$\int_0^\infty [\rho_V(\mu^2) - \rho_A(\mu^2)]d\mu^2 = 0. \tag{4}$$

Before proving these theorems, let us note some of their implications. The spectral functions $\rho_{V,A}(\mu^2)$ are measurable, in principle, from the cross sections for hadron production in electron-neutrino collisions. For the present, we can estimate $\rho_V(\mu^2)$ by using the hypothesis of ρ dominance:

$$\rho_V(\mu^2) \simeq g_\rho^2\delta(\mu^2 - m_\rho^2). \tag{5}$$

Eqs. (1) and (4) now read

$$\int_0^\infty \rho_A(\mu^2)\mu^{-2}d\mu^2 \simeq g_\rho^2 m_\rho^{-2} - F_\pi^2, \tag{6}$$

$$\int_0^\infty \rho_A(\mu^2)d\mu^2 \simeq g_\rho^2. \tag{7}$$

Hence, if $\rho_A(\mu^2)$ is sharply peaked about a point $\mu = m_A$, we must have[7]

$$m_A/m_\rho \simeq [1 - F_\pi^2 m_\rho^2/g_\rho^2]^{-1/2}. \tag{8}$$

Using ρ dominance and either current algebra[8] or the observed ρ width, we have $g_\rho^2 \simeq 2F_\pi^2 m_\rho^2$, so Eq. (8) gives

$$m_A/m_\rho \simeq \sqrt{2} \tag{9}$$

in extraordinary agreement with the observed[9] masses of the ρ and $A1$, for which $m_{A1}/m_\rho = 1.41 \pm 0.01$.

Now to the proof of Eqs. (1) and (4). Define

Chiral Perturbation Theory

Pions as pseudo Goldstone bosons

PHYSICAL REVIEW VOLUME 166, NUMBER 5 25 FEBRUARY 1968

Nonlinear Realizations of Chiral Symmetry*

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(Received 25 September 1967)

We explore possible realizations of chiral symmetry, based on isotopic multiplets of fields whose transformation rules involve only isotopic-spin matrices and the pion field. The transformation rules are unique, up to possible redefinitions of the pion field. Chiral-invariant Lagrangians can be constructed by forming isotopic-spin-conserving functions of a covariant pion derivative, plus other fields and their covariant derivatives. The resulting models are essentially equivalent to those that have been derived by treating chirality as an ordinary linear symmetry broken by the vacuum, except that we do not have to commit ourselves as to the grouping of hadrons into chiral multiplets; as a result, the unrenormalized value of g_A/g_V need not be unity. We classify the possible choices of the chiral-symmetry-breaking term in the Lagrangian according to their chiral transformation properties, and give the values of the pion-pion scattering lengths for each choice. If the symmetry-breaking term has the simplest possible transformation properties, then the scattering lengths are those previously derived from current algebra. An alternative method of constructing chiral-invariant Lagrangians, using ρ mesons to form covariant derivatives, is also presented. In this formalism, ρ dominance is automatic, and the current-algebra result from the ρ-meson coupling constant arises from the independent assumption that ρ mesons couple universally to pions and other particles. Including ρ mesons in the Lagrangian has no effect on the π-π scattering lengths, because chiral invariance requires that we also include direct pion self-couplings which cancel the ρ-exchange diagrams for pion energies near threshold.

I. INTRODUCTION

CURRENT algebra is useful because it allows us to obtain physical predictions from chiral symmetry. We have recently noted[1] that for soft-pion processes the same predictions can also be derived by a different method: Just use the lowest-order graphs generated by any chiral-invariant Lagrangian. The Lagrangian method has since been applied to pion production,[2] η decay,[3] K interactions and decay,[4] and, in various extended versions, to meson mass ratios and decay amplitudes,[5] and to the pion electromagnetic mass difference.[6] Opinions differ[7] as to whether any fundamental significance resides in the Lagrangians that have been used, but there is no doubt that they provide both a convenient method of calculation and a valuable heuristic guide to theorems that can be proved with current algebra.

There are two ways of constructing our chiral-invariant Lagrangians, which mirror two different views of the meaning of chiral symmetry. The first, *conventional* method[8] is to construct $\mathcal{L}$ to be manifestly chiral-invariant, as if chirality were an ordinary linear symmetry like isospin. For example, in the σ model[9] the $\boldsymbol{\pi}$ and σ fields form a four-vector coupled to nucleons in the combination $\sigma+i\boldsymbol{\tau}\cdot\boldsymbol{\pi}\gamma_5$, and the nucleon mass arises from the nonvanishing vacuum expectation value $\langle\sigma\rangle_0=-m_N/G$. In a closely related model[10] the Lagrangian takes the same form, but with σ replaced everywhere with $[(m_N/G)^2-\boldsymbol{\pi}^2]^{1/2}$. Such models suffer from a fundamental disadvantage: They hide the fact that soft pions are emitted in clusters by derivative cou-

* This work is supported in part through funds provided by the Atomic Energy Commission under Contract No. AT(30-1)-2098.

† On leave from the University of California, Berkeley, California.

[1] S. Weinberg, Phys. Rev. Letters **18**, 507 (1967).

[2] L.-N. Chang, Phys. Rev. **162**, 1497 (1967); Ph.D. thesis (unpublished).

[3] W. A. Bardeen, L. S. Brown, B. W. Lee, and H. T. Nieh, Phys. Rev. Letters **18**, 1170 (1967). Precisely the same calculation was done by S. Shei, but not published, because it appeared that the matrix element was too small by a factor m_π^2/m_η^2 to account for the observed decay. Bardeen *et al.* treat the η-π vertex in what seems to me a dubious manner, and thereby escape this difficulty.

[4] B. Zumino (to be published); S. Iwao (to be published).

[5] J. Schwinger, Phys. Letters **24B**, 473 (1967); S. Weinberg, Phys. Rev. Letters **18**, 507 (1967) (see footnote 7); S. Glashow, H. Schnitzer, and S. Weinberg, Phys. Rev. Letters **19**, 139 (1967) [Eq. (13) ff]; M. Lévy (to be published); J. W. Wess and B. Zumino, Phys. Rev. **163**, 1727 (1967); S. Glashow and S. Weinberg (to be published). The decay amplitudes derived using Lagrangian methods by Schwinger [and then in a somewhat more general form by Wess and Zumino] were subsequently rederived using current algebra by H. Schnitzer and S. Weinberg, Phys. Rev. **164**, 1828 (1967).

[6] J. Schwinger (to be published); D. B. Fairlie and K. Yoshida (to be published). The corresponding current-algebra calculation was done by T. Das, G. S. Guralnik, V. S. Mathur, F. E. Low, and J. E. Young, Phys. Rev. Letters **18**, 759 (1967).

[7] In particular, Schwinger has argued that as long as the origin of symmetries remains obscure, the phenomenological Lagrangian provides a suitable arena for their study. [J. Schwinger, Phys. Rev. **152**, 1219 (1966); also Refs. 5 and 6, and private communication.] Others like myself remain uneasy at using a symmetry on the phenomenological level, when it is not clear how *any* fundamental Lagrangian could give rise to the supposed symmetry of phenomena. From this point of view, chirality is in good shape because we have current algebra to underwrite it, but $SU(6)$ remains obscure. Time will tell.

[8] J. Schwinger, Ann. Phys. (N. Y.) **2**, 407 (1957); M. Gell-Mann and M. Lévy, Nuovo Cimento **16**, 705 (1960); F. Gürsey, *ibid.* **16**, 230 (1960); in *Proceedings of the 1960 Rochester Conference on High-Energy Physics* (Interscience Publishers, Inc., New York, (1960), p. 572; Ann. Phys. **12**, 91 (1961); F. Gürsey and B. Zumino, (unpublished); P. Chang and F. Gürsey, Phys. Rev. **164**, 1752 (1967); H. S. Mani, Y. Tomozawa, and Y. P. Yao, Phys. Rev. Letters **18**, 1084 (1967); L. S. Brown, Phys. Rev. **163**, 1802 (1967); and J. A. Cronin, Phys. Rev. **161**, 1483 (1967); and Refs. 1–4.

[9] J. Schwinger and M. Gell-Mann and M. Lévy, Ref. 8.

[10] F. Gürsey and M. Gell-Mann and M. Lévy, Ref. 8.

Electroweak Unification!!!

Using the Higgs mechanism:

- Consistent description of weak interactions,
- Unify weak and electromagnetic interactions,
- Prediction of neutral currents
- Prediction of W^+, W^-, Z and the Higgs!!!

A MODEL OF LEPTONS*

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(Received 17 October 1967)

Leptons interact only with photons, and with the intermediate bosons that presumably mediate weak interactions. What could be more natural than to unite[1] these spin-one bosons into a multiplet of gauge fields? Standing in the way of this synthesis are the obvious differences in the masses of the photon and intermediate meson, and in their couplings. We might hope to understand these differences by imagining that the symmetries relating the weak and electromagnetic interactions are exact symmetries of the Lagrangian but are broken by the vacuum. However, this raises the specter of unwanted massless Goldstone bosons.[2] This note will describe a model in which the symmetry between the electromagnetic and weak interactions is spontaneously broken, but in which the Goldstone bosons are avoided by introducing the photon and the intermediate-boson fields as gauge fields.[3] The model may be renormalizable.

We will restrict our attention to symmetry groups that connect the observed electron-type leptons only with each other, i.e., not with muon-type leptons or other unobserved leptons or hadrons. The symmetries then act on a left-handed doublet

$$L \equiv [\tfrac{1}{2}(1+\gamma_5)]\begin{pmatrix} \nu_e \\ e \end{pmatrix} \qquad (1)$$

and on a right-handed singlet

$$R \equiv [\tfrac{1}{2}(1-\gamma_5)]e. \qquad (2)$$

The largest group that leaves invariant the kinematic terms $-\bar{L}\gamma^\mu\partial_\mu L - \bar{R}\gamma^\mu\partial_\mu R$ of the Lagrangian consists of the electronic isospin $\vec{T}$ acting on L, plus the numbers N_L, N_R of left- and right-handed electron-type leptons. As far as we know, two of these symmetries are entirely unbroken: the charge $Q = T_3 - N_R - \frac{1}{2}N_L$, and the electron number $N = N_R + N_L$. But the gauge field corresponding to an unbroken symmetry will have zero mass,[4] and there is no massless particle coupled to N,[5] so we must form our gauge group out of the electronic isospin $\vec{T}$ and the electronic hyperchange $Y \equiv N_R + \frac{1}{2}N_L$.

Therefore, we shall construct our Lagrangian out of L and R, plus gauge fields $\vec{A}_\mu$ and B_μ coupled to $\vec{T}$ and Y, plus a spin-zero doublet

$$\varphi = \begin{pmatrix} \varphi^0 \\ \varphi^- \end{pmatrix} \qquad (3)$$

whose vacuum expectation value will break $\vec{T}$ and Y and give the electron its mass. The only renormalizable Lagrangian which is invariant under $\vec{T}$ and Y gauge transformations is

$$\mathfrak{L} = -\tfrac{1}{4}(\partial_\mu \vec{A}_\nu - \partial_\nu \vec{A}_\mu + g\vec{A}_\mu \times \vec{A}_\nu)^2 - \tfrac{1}{4}(\partial_\mu B_\nu - \partial_\nu B_\mu)^2 - \bar{R}\gamma^\mu(\partial_\mu - ig'B_\mu)R - L\gamma^\mu(\partial_\mu ig\vec{t}\cdot\vec{A}_\mu - i\tfrac{1}{2}g'B_\mu)L$$

$$-\tfrac{1}{2}|\partial_\mu\varphi - ig\vec{A}_\mu\cdot\vec{t}\varphi + i\tfrac{1}{2}g'B_\mu\varphi|^2 - G_e(\bar{L}\varphi R + \bar{R}\varphi^\dagger L) - M_1{}^2\varphi^\dagger\varphi + h(\varphi^\dagger\varphi)^2. \qquad (4)$$

We have chosen the phase of the R field to make G_e real, and can also adjust the phase of the L and Q fields to make the vacuum expectation value $\lambda \equiv \langle\varphi^0\rangle$ real. The "physical" φ fields are then φ^-

1970’s

Non-Abelian Gauge Theories of the Strong Interactions*

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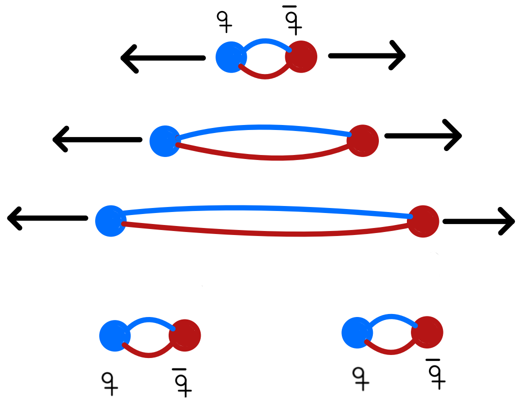
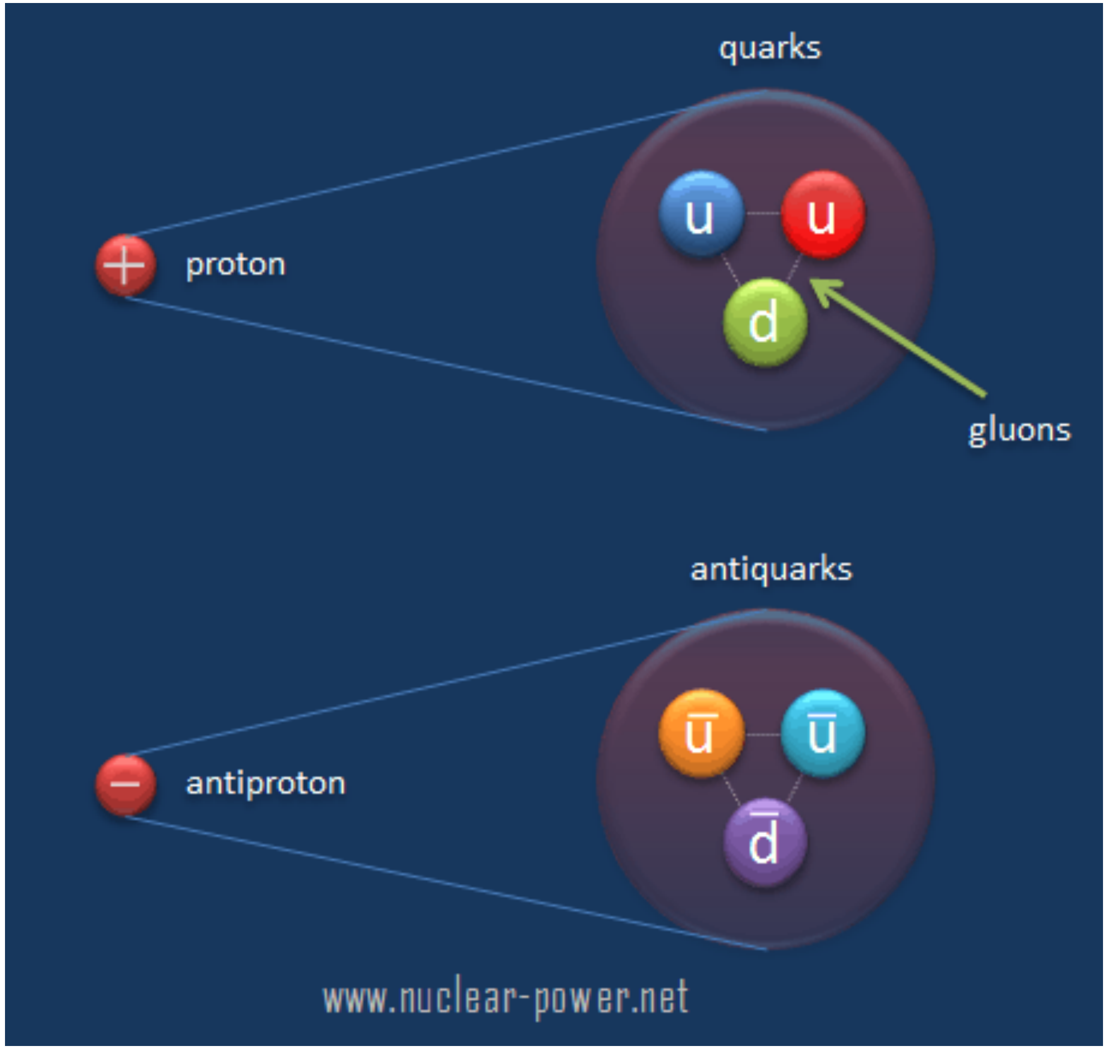
(Received 30 May 1973)

A class of non-Abelian gauge theories of strong interactions is described, for which parity and strangeness are automatically conserved, and for which the nonconservations of parity and strangeness produced by weak interactions are automatically of order α/m_W^2 rather than of order α. When such theories are "asymptotically free," the order-α weak corrections to natural zeroth-order symmetries may be calculated ignoring all effects of strong interactions. Speculations are offered on a possible theory of quarks.

Recently Gross and Wilczek and Politzer have made the exciting observation that non-Abelian gauge theories can exhibit free-field asymptotic behavior at large Euclidean momenta.[1] However, the physical application of this discovery raises serious problems: (1) Why don't the weak interactions produce parity and stangeness nonconservations of order α? (This problem finds a natural solution when the strong interactions are described by *Abelian* gauge models,[2] but not, to the best of my knowledge, in non-Abelian models of the "Berkeley" type.[3]) (2) Even with asymptotic freedom, when can the strong interactions actually be neglected? (3) Even if asymptotic freedom explains the success of naive quark-model calculations, why don't we see physical quarks? This

QCD: Massless and confined gluons!





ULTIMATE TEMPERATURE AND THE EARLY UNIVERSE*

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(Received 5 August 1970)

The early history of the universe is discussed in the context of an exponentially rising density of particle states.

There are now plausibe theoretical models[1] for the thermal history of the universe back to the time of helium synthesis, when the temperature was 0.1 to 1 MeV. Our present theoretical apparatus is really inadequate to deal with much earlier times, say when $T \gtrsim 100$ MeV, and in lieu of any better ideas it is usual to treat the matter of the very early universe as consisting of a number of species of essentially free particles. But how many species?

At one extreme, it might be assumed that the number of particle species stays fixed (perhaps just quarks, antiquarks, leptons, antileptons, photons, and gravitons). In this case, the temperature T will vary with the cosmic scale factor[2] $R(t)$ according to the relation $T \propto 1/R$. The present universe should then contain various relics of the early inferno: There sould be a 1°K blackbody gravitational radiation,[3] if TR stayed roughly constant between the times that the gravitons and the photons decoupled from the rest of the universe; also, according to Zeldovich,[4] the leftover quarks should be about as common as gold atoms. The gravitational radiation would not have been seen, but the quarks would have been, unless, of course, quarks do not exist.

At the other extreme, one might assume that the number of species of particles with mass between m and $m+dm$ increases as $m\to\infty$ as fast as possible:

$$N(m)dm \to Am^{-B}e^{\beta_0 m}dm. \quad (1)$$

If $N(m)$ increased any faster, the partition function would not converge. With the increase (1), the partition function converges only if the temperature[5] is less than $1/\beta_0$. The quantity $T_0 \equiv 1/\beta_0$ is thus a maximum temperature for any system in thermal equilibrium.

Support for this latter sort of model comes from two quite different directions:

(1) The transverse momentum distribution of secondaries in very high energy collisions is observed to be roughly $\exp(-|p_\perp|/160 \text{ MeV})$. Hagedorn[6] interprets this distribution in terms of a statistical model with $T_0 \simeq 160$ MeV and $B=\frac{5}{2}$.

(2) If particles fall on families of parallel linearly rising Regge trajectories, their masses take discrete values $m_1, m_2, \cdots$, where

$$\alpha' m_n^2 + \alpha_0 = n. \quad (2)$$

Here $\alpha' \approx 1 \text{ GeV}^{-2}$ is the universal Regge slope and α_0 is a number, of order unity, characterizing the family. The extension of the Veneziano model[7] to multiparticle reactions requires[8] that the number of particle states with mass m_n equals the degeneracy of the eigenvalue n of the operator

$$N = \sum_{\mu=1}^{D}\sum_{k=1}^{\infty} k a_{\mu k}^\dagger a_{\mu k}, \quad (3)$$

where $a_{\mu k}$ and $a_{\mu k}^\dagger$ are an infinite set of annihilation and creation operators. For $n\to\infty$, this number is[9]

$$P_{nD} \to 2^{-1/2}(D/24)^{(D+1)/4} n^{-(D+3)/4} \times \exp\{2\pi(\tfrac{1}{6}Dn)^{1/2}\}. \quad (4)$$

Equations (2) and (4) lead to an asymptotic level density of form (1), with

$$\beta_0 = 2\pi(\tfrac{1}{6}D\alpha')^{1/2}, \quad B = \tfrac{1}{2}(D+1). \quad (5)$$

The value of D is not certain—originally Fubini and Veneziano[8] had $D=4$, but Lovelace[10] argues that D is larger, possibly $D=5$.

Table I summarizes the values of T_0 and B for these various models. Lovelace[10] has emphasized the striking agreement between the values of T_0 derived in such different ways. We now see that

Table I. Possible values of the parameters in the level-density formula (1).

Model	$T_0 \equiv 1/\beta_0$	B
(1) Hagedorn[a]	~160 MeV	$\frac{5}{2}$
(2) Veneziano[b] (with $\alpha'=1$ GeV^{-2})		
$D=4$	180 MeV	$\frac{5}{2}$
$D=5$	174 MeV	3
$D=6$	159 MeV	$\frac{7}{2}$
$D=7$	147 MeV	4

[a]Ref. 6. [b]Ref. 8.

Hagedorn Temperature (String theory)

Quantum Field Theory at finite temperature

PHYSICAL REVIEW D VOLUME 9, NUMBER 12 15 JUNE 1974

Gauge and global symmetries at high temperature*

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It is shown how finite-temperature effects in a renormalizable quantum field theory can restore a symmetry which is broken at zero temperature. In general, for both gauge symmetries and ordinary symmetries, such effects occur only through a temperature-dependent change in the effective bare mass of the scalar bosons. The change in the boson bare mass is calculated for general field theories, and the results are used to derive the critical temperatures for a few special cases, including gauge and nongauge theories. In one case, it is found that a symmetry which is unbroken at low temperature can be broken by raising the temperature above a critical value. An appendix presents a general operator formalism for dealing with higher-order effects, and it is observed that the one-loop diagrams of field theory simply represent the contribution of zero-point energies to the free energy density. The cosmological implications of this work are briefly discussed.

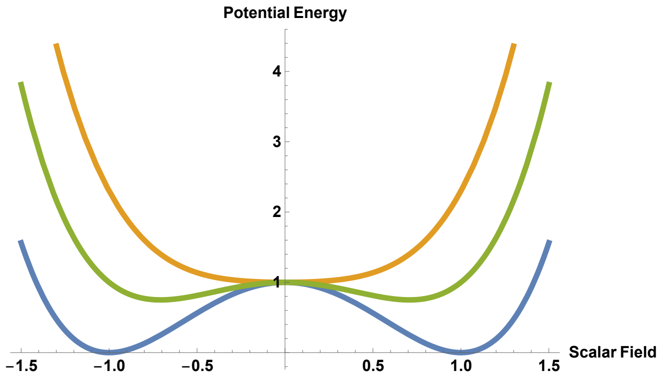
I. INTRODUCTION

The idea of broken symmetry was originally brought into elementary-particle physics on the basis of experience with many-body systems.[1] Just as a piece of iron, although described by a rotationally invariant Hamiltonian, may spontaneously develop a magnetic moment pointing in any given direction, so also a quantum field theory may imply physical states and *S* matrix elements which do not exhibit the symmetries of the Lagrangian.

It is natural then to ask whether the broken symmetries of elementary-particle physics would be restored by heating the system to a sufficiently high temperature, in the same way as the rotational invariance of a ferromagnet is restored by raising its temperature. A recent paper by Kirzhnits and Linde[2] suggests that this is indeed the case. However, although their title refers to a gauge theory, their analysis deals only with ordinary theories with broken global symmetries. Also, they estimate but do not actually calculate the critical temperature at which a broken symmetry is restored.

The purpose of this article is to extend the analysis of Kirzhnits and Linde to gauge theories,[3] and to show how to calculate the critical temperature for general renormalizable field theories, with either gauge or global symmetries. Our results completely confirm the more qualitative conclusions of Kirzhnits and Linde.[2]

The diagrammatic formalism[4] used here is described in Sec. II. Any finite-temperature Green's function is given by a sum of Feynman diagrams, just as in field theory, except that en-



Symmetry restoration at finite temperature

Hierarchy of Interactions in Unified Gauge Theories*

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(Received 15 May 1974)

We present a general formalism for calculating the renormalization effects which make strong interactions strong in simple gauge theories of strong, electromagnetic, and weak interactions. In an SU(5) model the superheavy gauge bosons arising in the spontaneous breakdown to observed interactions have mass perhaps as large as 10^{17} GeV, almost the Planck mass. Mixing-angle predictions are substantially modified.

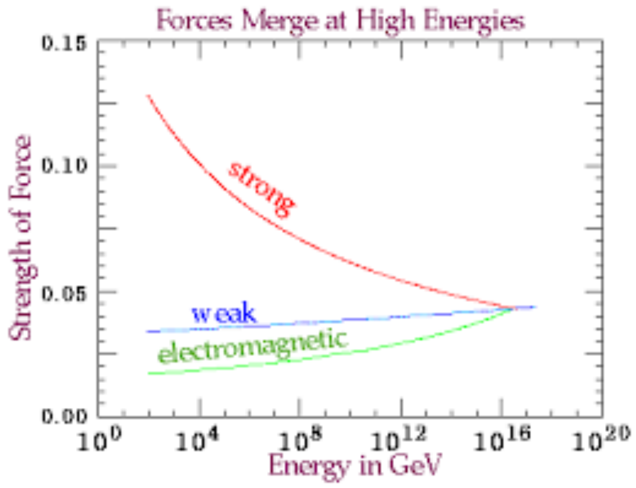
The scaling observed in deep inelastic electron scattering suggests that what are usually called the strong interactions are not so strong at high energies. Asymptotically free gauge theories of the strong interactions[1] provide a possible explanation: The gluon coupling constant $g(\mu)$ (defined as the value of a three-gluon or gluon-fermion-fermion vertex with momenta characterized by a mass μ) is small when μ is several GeV or larger, but becomes large when μ is small, through the piling up of the logarithms encountered in perturbation theory. In one recent calculation[2] a fit was found for a gauge coupling [in a color SU(3) model][3] with $g^2(\mu)/4\pi \simeq 0.1$ when $\mu \simeq 2$ GeV.

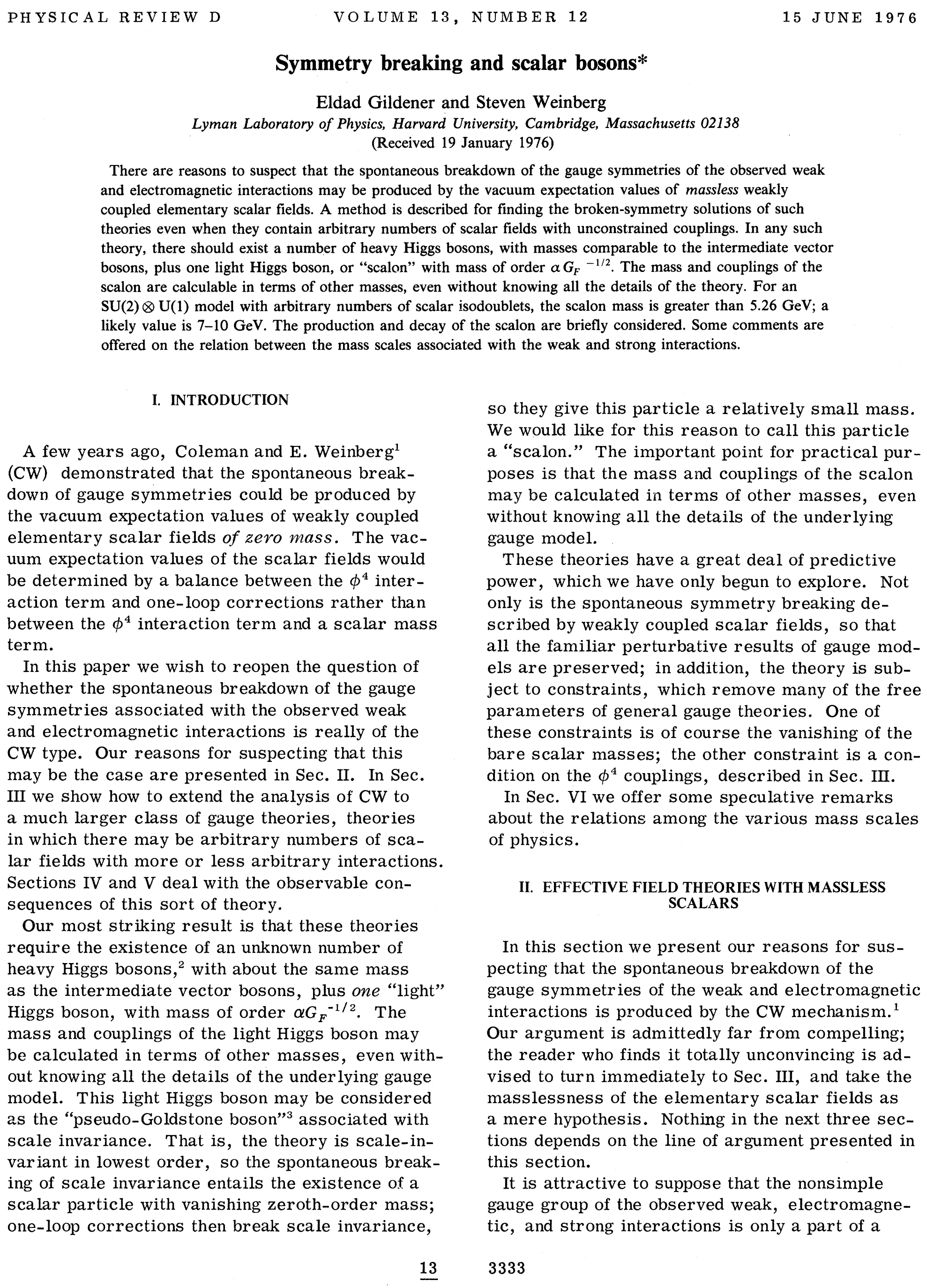
If $g(\mu)$ is small when μ is large, then perhaps the strong gauge coupling at some large fundamental mass is of the same order as the couplings in gauge theories of the weak and electromagnetic interactions.[4] Georgi and Glashow[5] have recently gone one step farther, and proposed a model based on the *simple* gauge group SU(5), in which there naturally appears only one free gauge coupling. In their model, SU(5) suffers a spontaneous breakdown to the gauge subgroups SU(3) and SU(2)⊗U(1), which are associated respectively with the strong[3] and the weak and electromagnetic[6] interactions. In order to suppress unobserved interactions, Georgi and Glashow made the necessary assumption[7] that some vector bosons are superheavy.

We find the notion of a simple gauge group uniting strong, weak, and electromagnetic interactions extraordinarily attractive. However, as emphasized by Georgi and Glashow, the success of any such scheme hinges on an understanding of the effects which produce the obvious disparity in strength between the strong and the weak and electromagnetic interactions at ordinary energies. We therefore wish to present in this paper a general formalism for the calculation of such effects. This will lead us to an estimate of the mass of the superheavy gauge bosons. Where a specific model of the gauge groups of the observed interactions is needed as an example, we shall assume that the strong and the weak and electromagnetic interactions are described by color SU(3)[3] and by SU(2)⊗U(1), respectively, and where a specific example of a unifying simple gauge group is needed, we shall use SU(5).

If we neglect all renormalization effects, the embedding of the gauge groups G_i of the observed interactions in a larger simple group G imposes a relation among their coupling constants. We

Grand Unification !





PHYSICAL REVIEW D VOLUME 13, NUMBER 12 15 JUNE 1976

Symmetry breaking and scalar bosons*

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(Received 19 January 1976)

There are reasons to suspect that the spontaneous breakdown of the gauge symmetries of the observed weak and electromagnetic interactions may be produced by the vacuum expectation values of *massless* weakly coupled elementary scalar fields. A method is described for finding the broken-symmetry solutions of such theories even when they contain arbitrary numbers of scalar fields with unconstrained couplings. In any such theory, there should exist a number of heavy Higgs bosons, with masses comparable to the intermediate vector bosons, plus one light Higgs boson, or "scalon" with mass of order $\alpha G_F^{-1/2}$. The mass and couplings of the scalon are calculable in terms of other masses, even without knowing all the details of the theory. For an SU(2) $\otimes$ U(1) model with arbitrary numbers of scalar isodoublets, the scalon mass is greater than 5.26 GeV; a likely value is 7–10 GeV. The production and decay of the scalon are briefly considered. Some comments are offered on the relation between the mass scales associated with the weak and strong interactions.

I. INTRODUCTION

A few years ago, Coleman and E. Weinberg[1] (CW) demonstrated that the spontaneous breakdown of gauge symmetries could be produced by the vacuum expectation values of weakly coupled elementary scalar fields *of zero mass*. The vacuum expectation values of the scalar fields would be determined by a balance between the ϕ^4 interaction term and one-loop corrections rather than between the ϕ^4 interaction term and a scalar mass term.

In this paper we wish to reopen the question of whether the spontaneous breakdown of the gauge symmetries associated with the observed weak and electromagnetic interactions is really of the CW type. Our reasons for suspecting that this may be the case are presented in Sec. II. In Sec. III we show how to extend the analysis of CW to a much larger class of gauge theories, theories in which there may be arbitrary numbers of scalar fields with more or less arbitrary interactions. Sections IV and V deal with the observable consequences of this sort of theory.

Our most striking result is that these theories require the existence of an unknown number of heavy Higgs bosons,[2] with about the same mass as the intermediate vector bosons, plus *one* "light" Higgs boson, with mass of order $\alpha G_F^{-1/2}$. The mass and couplings of the light Higgs boson may be calculated in terms of other masses, even without knowing all the details of the underlying gauge model. This light Higgs boson may be considered as the "pseudo-Goldstone boson"[3] associated with scale invariance. That is, the theory is scale-invariant in lowest order, so the spontaneous breaking of scale invariance entails the existence of a scalar particle with vanishing zeroth-order mass; one-loop corrections then break scale invariance, so they give this particle a relatively small mass. We would like for this reason to call this particle a "scalon." The important point for practical purposes is that the mass and couplings of the scalon may be calculated in terms of other masses, even without knowing all the details of the underlying gauge model.

These theories have a great deal of predictive power, which we have only begun to explore. Not only is the spontaneous symmetry breaking described by weakly coupled scalar fields, so that all the familiar perturbative results of gauge models are preserved; in addition, the theory is subject to constraints, which remove many of the free parameters of general gauge theories. One of these constraints is of course the vanishing of the bare scalar masses; the other constraint is a condition on the ϕ^4 couplings, described in Sec. III.

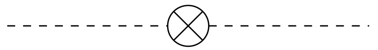
In Sec. VI we offer some speculative remarks about the relations among the various mass scales of physics.

II. EFFECTIVE FIELD THEORIES WITH MASSLESS SCALARS

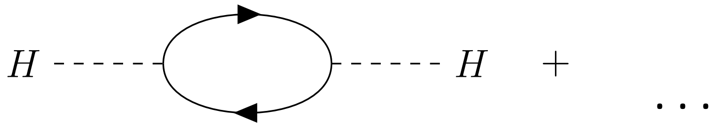
In this section we present our reasons for suspecting that the spontaneous breakdown of the gauge symmetries of the weak and electromagnetic interactions is produced by the CW mechanism.[1] Our argument is admittedly far from compelling; the reader who finds it totally unconvincing is advised to turn immediately to Sec. III, and take the masslessness of the elementary scalar fields as a mere hypothesis. Nothing in the next three sections depends on the line of argument presented in this section.

It is attractive to suppose that the nonsimple gauge group of the observed weak, electromagnetic, and strong interactions is only a part of a

Hierarchy Problem



$$\frac{m_h}{M_P} \sim 10^{-15}$$



Why the Higgs mass is not much bigger (fine tuning)??

PHYSICAL REVIEW D VOLUME 13, NUMBER 4 15 FEBRUARY 1976

Implications of dynamical symmetry breaking

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(Received 8 September 1975)

An analysis is presented of the physical implications of theories in which the masses of the intermediate vector bosons arise from a dynamical symmetry breaking. In the absence of elementary spin-zero fields or bare fermion masses, such theories are necessarily invariant to zeroth order in the weak and electromagnetic gauge interactions under a global $U(N)\otimes U(N)$ symmetry, where N is the number of fermion types, not counting color. This symmetry is broken both intrinsically by the weak and electromagnetic interactions and spontaneously by dynamical effects of the strong interactions. An effective Lagrangian is constructed which allows the calculation of leading terms in matrix elements at low energy; this effective Lagrangian is used to analyze the relative direction of the intrinsic and spontaneous symmetry breakdown and to construct a unitarity gauge. Spontaneously broken symmetries which belong to the gauge group of the weak and electromagnetic interactions correspond to fictitious Goldstone bosons which are removed by the Higgs mechanism. Spontaneously broken symmetries of the weak and electromagnetic interactions which are not members of the gauge group correspond to true Goldstone bosons of zero mass; their presence makes it difficult to construct realistic models of this sort. Spontaneously broken elements of $U(N)\otimes U(N)$ which are not symmetries of the weak and electromagnetic interactions correspond to pseudo-Goldstone bosons, with mass comparable to that of the intermediate vector bosons and weak couplings at ordinary energies. Quark masses in these theories are typically less than 300 GeV by factors of order α. These theories require the existence of "extra-strong" gauge interactions which are not felt at energies below 300 GeV.

I. INTRODUCTION

When unified gauge theories of the weak and electromagnetic interactions were first proposed, it was assumed[1] that the spontaneous symmetry breakdown responsible for the intermediate-vector-boson masses is due to the vacuum expectation values of a set of spin-zero fields. For a variety of reasons, the attention of theorists has since been increasingly drawn to the possibility that this symmetry breaking is of a purely dynamical nature.[2] That is, it is supposed that there may be no elementary spin-zero fields in the Lagrangian, and that the Goldstone bosons associated with the spontaneous symmetry breakdown are bound states.

Almost all the effort that has been put into analyses of dynamical symmetry breaking has been directed to the difficult mathematical problem, of whether and how this phenomenon can occur in a variety of field-theoretic models. In this article I would like to address quite a different question: Assuming that dynamical symmetry breaking is a mathematical possibility in gauge field theories, what are the consequences for the real world?

Why should we believe that the masses of the intermediate vector bosons arise from dynamical symmetry breaking? The absence of *strongly* interacting elementary spin-zero fields is indicated by a number of requirements: asymptotic freedom,[3] electroproduction sum rules,[4] and the naturalness of order-α parity and strangeness conservation.[5] On the other hand, the absence of *weakly* interacting elementary spin-zero fields is much less certain. Apart from simplicity, the best reason for this assumption comes from the requirement for a natural hierarchy of gauge symmetry breaking.[6] In order to put together the observed weak and electromagnetic interactions into a simple gauge group, it is necessary to suppose[7] that in the spontaneous breakdown of this simple group to the nonsimple gauge group of the observed interactions, vector-boson masses are generated that are much larger than the masses expected for the W and Z; this conclusion is even stronger if we try to include the strong interactions as well.[8] This superstrong symmetry breakdown may well be due to the vacuum expectation values of elementary spin-zero fields. However, at ordinary energies, far below the superheavy vector-boson masses, physics is described by an effective field theory[9] involving those fermions and vector bosons that did not get masses from the superstrong spontaneous symmetry breakdown, but no spin-zero fields. Likewise, the gauge group of this effective field theory consists of a direct product of those simple and U(1) subgroups of the simple gauge group that were not broken at the superstrong level. The only way that the non-superheavy fermions and vector bosons can then

Technicolor

Is the Higgs composite? (No!)

Neutrino cosmology

PHYSICAL REVIEW LETTERS

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Cosmological Lower Bound on Heavy-Neutrino Masses

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The present cosmic mass density of possible stable neutral heavy leptons is calculated in a standard cosmological model. In order for this density not to exceed the upper limit of 2×10^{-29} g/cm^3, the lepton mass would have to be *greater* than a lower bound of the order of 2 GeV.

There is a well-known cosmological argument[1] against the existence of neutrino masses greater than about 40 eV. In the "standard" big-bang cosmology,[2] the present number density of each kind of neutrino is expected[3] to be $\frac{6}{11}$ the number density of photons in the 3°K black-body background radiation, or about 300 cm^{-3}; hence if the neutrino mass were above 40 eV, their mass density would be greater than 2×10^{-29} g/cm^3, which is roughly the upper limit allowed by present estimates[4] of the Hubble constant and the deceleration parameter.

However, this argument would not apply if the neutrino mass were much larger than 1 MeV. Neutrinos are generally expected[2] to go out of thermal equilibrium when the temperature drops to about 10^{10} °K, the temperature at which neutrino collision rates become comparable to the expansion rate of the universe. If neutrinos were much heavier than 1 MeV, then they would already be much rarer than photons at the time when they go out of thermal equilibrium, and hence their number density would now be much less than 300 cm^{-3}.

Of course, the familiar electronic and muonic neutrinos are known to be lighter than 1 MeV. However, heavier stable neutral leptons could easily have escaped detection, and are even required in some gauge models.[5] In this Letter, we suppose that there exists a neutral lepton L^0 (the "heavy neutrino") with mass well above 1 MeV, and we assume that L^0 carries some additive or multiplicative quantum number which keeps it absolutely stable. We will present arguments based on the standard big-bang cosmology to show that the mass of such a particle must be above a lower bound of order 2 GeV.

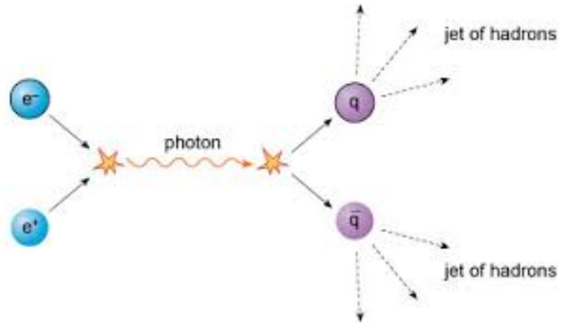
At first glance, it might be thought that the present number density of heavy neutrinos would simply be less than the above estimate of 300 cm^{-3} by the value $\exp[-m_L/(1\text{ MeV})]$ of the Boltzmann factor at the time the heavy neutrinos go out of thermal equilibrium. If this were the case, then an upper limit of 2×10^{-29} g/cm^{-3} on the present cosmic mass density would require that $m_L\exp[-m_L/(1\text{ MeV})]$ should be less than 40 eV, and hence that m_L should either be less than 40 eV or greater than 13 MeV.

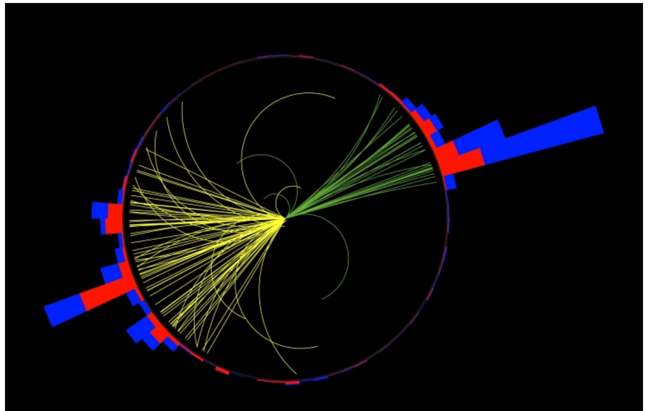
However, the true lower bound on the heavy-neutrino mass is considerably more stringent.

Jets to detect quarks

Quark jets $e^+e^- \to q\bar{q} \to 2\,\text{jets}$

Gluon jets $e^+e^- \to q\bar{q}q \to 3\,\text{jets}$





VOLUME 39, NUMBER 23 PHYSICAL REVIEW LETTERS 5 DECEMBER 1977

Jets from Quantum Chromodynamics

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The properties of hadronic jets in e^+e^- annihilation are examined in quantum chromodynamics, without using the assumptions of the parton model. We find that two-jet events dominate the cross section at high energy, and have the experimentally observed angular distribution. Estimates are given for the jet angular radius and its energy dependence. We argue that the detailed results of perturbation theory for production of arbitrary numbers of quarks and gluons can be reinterpreted in quantum chromodynamics as predictions for the production of jets.

The observation[1] of hadronic jets in e^+e^- annihilation provides one of the most striking confirmations of the parton picture.[2] In particular, the distribution of events in the angle θ between the jet axis and the e^+-e^- beam line is observed to be very close to the form $1+\cos^2\theta$ that would be expected for the production of a pair of relativistic charged pointlike particles of spin $\frac{1}{2}$. We shall argue here that the existence, angular distribution, and some aspects of the structure of these jets follow as consequences of the perturbation expansion[3] of quantum chromodynamics[4] (QCD), without assuming the parton picture (in particular, the transverse-momentum cutoff) in advance. Thus, the observed features of jets provide evidence for an underlying asymptotically free gauge field theory with elementary spin-$\frac{1}{2}$ quarks. We also wish here to demonstrate a general approach, which may be applicable to a wide range of high-energy phenomena.

Our procedure is to define a partial cross section for jet production, which in asymptotically free theories like QCD can be calculated perturbatively at high energy. By ordinary dimensional analysis, any sort of total or partial cross section in QCD can be written in the form

$$\sigma = E^{-2} f(m/E, g_E, x), \tag{1}$$

where E is the energy; x stands for all other dimensionless variables characterizing the final state; m stands for all mass variables; and g_E is the gauge coupling constant, defined at a renormalization point with four-momenta of order E. [We express the cross section in terms of g_E, rather than a coupling g_κ defined at a renormalization point with momenta of arbitrary scale κ, in order to avoid factors of $\ln(E/\kappa)$. Physical quantities are of course independent of the choice of renormalization point.] Even in asymptotically free theories, where $g_E \to 0$ as $E \to \infty$, it is generally not possible to calculate the cross section perturbatively for large E, because the cross section will exhibit singularities for $m/E \to 0$. It is of course for this reason that asymptotic freedom has as a rule been used directly to justify perturbative calculations of Green's functions and Wilson coefficient functions, rather than cross sections themselves.

However, by performing various sums over states, it is possible to define a wide range of cross sections which are free of $m \to 0$ singularities. To learn what they are, we observe that quantum field theories of massless particles have always been found (in the absence of superrenormalizable couplings) to be physically sensible, i.e., that any cross section which would actually be measurable in such a massless theory is free of infrared divergences in each order of perturbation theory.[5] Hence in the real world with $m \neq 0$, any sort of partial cross section which would be measurable for $m = 0$ is expected to be free of singularities in m as $m \to 0$, and can therefore be calculated perturbatively[3] in QCD for $E \to \infty$.

For instance, the cross section for production of a definite number of particles does have singularities for $m \to 0$, because for $m = 0$ we could not expect to be able to tell the difference between one particle or several particles moving in the same direction. At the opposite extreme, the total cross section for $e^+e^- \to$ hadrons would clearly be measurable even for zero quark mass, and hence must be free of singularities in m (to

Axions!!!

$$\mathcal{L}_\theta = \theta_3 \frac{g_s^2}{64\pi^2} \epsilon^{\mu\nu\rho\sigma} G^A_{\mu\nu} G^A_{\rho\sigma}$$

$$\theta_3 \leq 10^{-10}$$???

θ_3 Promoted to a field $a(x)$ Peccei-Quinn

$$a \to a + \alpha.$$

$$V(a) = E(a(x), \bar\theta) \sim -F_\pi^2 m_\pi^2 \cos\left(\frac{a}{f_a} + \bar\theta\right)$$

Minimum of energy at a=0

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A New Light Boson?

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It is pointed out that a global U(1) symmetry, that has been introduced in order to preserve the parity and time-reversal invariance of strong interactions despite the effects of instantons, would lead to a neutral pseudoscalar boson, the "axion," with mass roughly of order 100 keV to 1 MeV. Experimental implications are discussed.

One of the attractive features of quantum chromodynamics[1] (QCD) is that it offers an explanation of why C, P, T, and all quark flavors are conserved by strong interactions, and by order-α effects of weak interactions.[2] However, the discovery of quantum effects[3] associated with the "instanton" solution of QCD has raised a puzzle with regard to P and T conservation. Because of Adler-Bell-Jackiw anomalies, the chiral transformation which is needed in QCD to bring the quark-mass matrix to a real, diagonal, γ_5-free form will in general change the phase angle θ associated[3] with instanton effects, leaving $\bar\theta \equiv \theta + \arg\det m$ invariant. [Here m is the coefficient of $\frac{1}{2}(1+\gamma_5)$ in a decomposition of the quark-mass matrix into $\frac{1}{2}(1\pm\gamma_5)$.] The condition for P and T conservation is that $\theta = 0$ when the quark fields are defined so that m is real, or more generally, that $\bar\theta = 0$. But θ is a free parameter, and in QCD there is no reason why it should take the value $-\arg\det m$. Furthermore, even if we simply demanded that the strong interactions in isolation conserve P and T, so that $\bar\theta = 0$, there would still be a danger that the weak interactions would introduce P- and T-nonconserving phases of order $10^{-3}\alpha$ in m, leading to an unacceptable neutron electric dipole moment, of order 10^{-18} $e\cdot$cm.

An attractive resolution of this problem has been proposed by Peccei and Quinn.[5] They note that the quark-mass matrix is a function $m(\langle\varphi\rangle)$ of the vacuum expectation values of a set of weakly coupled scalar fields φ_i. Although θ is arbitrary, $\langle\varphi\rangle$ is not; it is determined by the minimization of a potential $V(\varphi)$ which depends on θ. Peccei and Quinn assume that the Lagrangian has a global U(1) chiral symmetry [which I will call $\mathrm{U(1)_{PQ}}$], under which $\det m(\varphi)$ changes by a phase. The phase of $\det m(\varphi)$ at the minimum of $V(\varphi)$ is then undetermined in any finite order of perturbation theory, and is fixed only by instanton effects which break the $\mathrm{U(1)_{PQ}}$ symmetry. However, the potential will then depend on $\bar\theta$, but not separately on θ and $\arg\det m$, so that it is not a miracle if the phase of $\det m(\varphi)$ at the minimum of $V(\varphi)$ happens to have the P- and T-conserving value $-\theta$. Peccei and Quinn[5] show in a number of examples that this is just what happens.

Now, the $\mathrm{U(1)_{PQ}}$ symmetry of the Lagrangian is intrinsically broken by instantons, and so at first sight one might not expect that it would have any further physical consequences. Certainly it does not lead to the strongly interacting isoscalar pseudoscalar meson below $\sqrt{3}m_\pi$,[6] that was the bugbear of the old U(1) problem. However, the scalar fields φ do not know about instantons, except through a semiweak ($\propto G_F^{1/2}$) coupling to quarks. Hence the spontaneous breakdown of the chiral $\mathrm{U(1)_{PQ}}$ symmetry associated with the appearance of nonzero vacuum expectation values $\langle\varphi\rangle$ leads[7] to a very light pseudoscalar pseudo-Goldstone boson,[8] the "axion," with m_a^2 proportional to the Fermi coupling G_F.

For insight in to the properties of the axion, it is useful to examine how they appear in the simplest realistic model that admits a $\mathrm{U(1)_{PQ}}$ symmetry. We assume an SU(2)⊗ U(1) gauge group, with quarks in $N/2$ left-handed doublets and N right-handed singlets, and just two scalar doublets $\{\varphi_i^+, \varphi_i^0\}$, carrying $\mathrm{U(1)_{PQ}}$ quantum numbers such that φ_1 (φ_2) couples right-handed quarks of charge $-\frac{1}{3}$ ($+\frac{2}{3}$) to left-handed quarks. By writing the Yukawa interaction in terms of quark fields of definite mass, we easily see that the interaction of neutral scalar fields with quarks is[9]

$$\mathfrak{L}_N = -[m_d\bar d_R d_L + m_s\bar s_R s_L + m_b\bar b_R b_L + \cdots]\varphi_1^{0*}\langle\varphi_1^0\rangle^{*-1} - [m_u\bar u_R u_L + m_c\bar c_R c_L + m_t\bar t_R t_L + \cdots]\varphi_2^0\langle\varphi_2^0\rangle^{-1} + \mathrm{H.c.}, \quad (1)$$

where L and R indicate multiplication with $\frac{1}{2}(1\pm\gamma_5)$. The part of $\mathfrak{L}_N$ involving the light quarks u, d, and s may be treated as a perturbation $\mathfrak{L}_{uds}$, while terms in $\mathfrak{L}_N$ involving c, t, b,... must be included in the

Effective Field Theories (EFT)

$$\mathcal{L} = \underbrace{\underbrace{\partial^\mu \phi \partial_\mu \phi - m^2\phi^2 - g\phi^3 - \lambda\phi^4}_{\text{Renormalisable}} + \frac{\alpha}{\Lambda}\phi^5 + \frac{\beta}{\Lambda^2}\phi^6 + \cdots}_{\text{Non-Renormalisable}}$$

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PHENOMENOLOGICAL LAGRANGIANS*

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1. Introduction: A reminiscence

Julian Schwinger's ideas have strongly influenced my understanding of phenomenological Lagrangians since 1966, when I made a visit to Harvard. At that time, I was trying to construct a phenomenological Lagrangian which would allow one to obtain the predictions of current algebra for soft pion matrix elements with less work, and with more insight into possible corrections. It was necessary to arrange that the pion couplings in the Lagrangian would all be derivative interactions, to suppress the incalculable graphs in which soft pions would be emitted from internal lines of a hard-particle process. The mathematical approach I followed[1]) at first was quite clumsy; I started with the old σ-model[2]), in which the pion is in a chiral quartet with a 0+ isoscalar σ; then performed a space-time dependent chiral rotation which transformed $\{\boldsymbol{\pi}, \sigma\}$ everywhere into $\{0, \sigma'\}$ with $\sigma' \equiv (\sigma^2 + \pi)^{1/2}$; and then re-introduced the pion field as the chiral rotation "angle". The Lagrangian obtained in this way had a complicated and unfamiliar non-linear structure, but it did have the desired property of derivative coupling, because any space-time independent part of the rotation "angle" would correspond to a symmetry of the theory, and so would not contribute to the Lagrangian.

Schwinger suggested to me that one might be able to construct a suitable phenomenological Lagrangian directly, by introducing a pion field which from the beginning would have the non-linear transformation property of chiral rotation angles, and then just obeying the dictates of chiral symmetry for such a pion field[3]). Following this suggestion, I worked out a general theory[4]) of non-linear realizations of chiral $SU(2) \times SU(2)$, which was soon after generalized to arbitrary groups in elegant papers of Callan, Coleman, Wess, and Zumino[5]), and has since been applied by many authors[6]). The importance of the approach suggested by Schwinger has been not only that it saves the work

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Violation of Lepton and Baryon numbers

$$\mathcal{L} = \mathcal{L}_{SM} + \frac{1}{M}\mathcal{L}_5 + \frac{1}{M^2}\mathcal{L}_6 + \mathcal{O}\left(\frac{1}{M^3}\right)$$

$$\mathcal{L}_5 = \left(\frac{\lambda_\nu}{M}\right) HH\nu_L\nu_L\,, \quad M \gg m_W \qquad M \sim 10^{14}\mathrm{GeV}$$

Neutrino masses

$$\mathcal{L}_6 = \left(\frac{\beta}{M^2}\right) qqql \qquad M \geq 10^{15}\mathrm{GeV}$$

Proton decay

Baryon- and Lepton-Nonconserving Processes

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A number of properties of possible baryon- and lepton-nonconserving processes are shown to follow under very general assumptions. Attention is drawn to the importance of measuring μ^+ polarizations and $\overline{\nu}_e/e^+$ ratios in nucleon decay as a means of discriminating among specific models.

Of the supposedly exact conservation laws of physics, two are especially questionable: the conservation of baryon number and lepton number. As far as we know, there is no necessity for an *a priori* principle of baryon and lepton conservation. As we shall see, even without such a principle, the fact that the weak, electromagnetic, and strong interactions of ordinary quarks and leptons conserve baryon and lepton number can be understood as simply a consequence of the SU(2)⊗U(1) and SU(3) gauge symmetries. Also, in contrast with the conservation of charge, color, and energy and momentum, the conservation of baryon number and lepton number are almost certainly not unbroken local symmetries.[1] Not only is baryon conservation unnecessary as a fundamental principle, the apparent excess of baryons over antibaryons in our universe provides a positive clue that some sort of physical processes have actually violated baryon-number conservation.[2] Violations of baryon and lepton conservation are likely to occur in grand unified theories that combine the gauge theory of weak and electromagnetic interactions with that of strong interactions and have leptons and quarks in the same gauge multiplets, and such violations have been found in various of these models.[3]

The purpose of this paper is to point out those features of baryon- or lepton-nonconserving processes that are to be expected on very general grounds. Other features will be indicated that may be used to discriminate among specific models.

No grand unified model or other specific gauge model of baryon- and lepton-nonconserving processes will be adopted here. Instead, it will simply be assumed that these processes are mediated by some unspecified "superheavy" particles, with a characteristic mass M above, say, 10^{14} GeV. Such large masses are indicated by the experimental lower bound[4] on the proton lifetime, and are also required in order that these parti-

1980’s

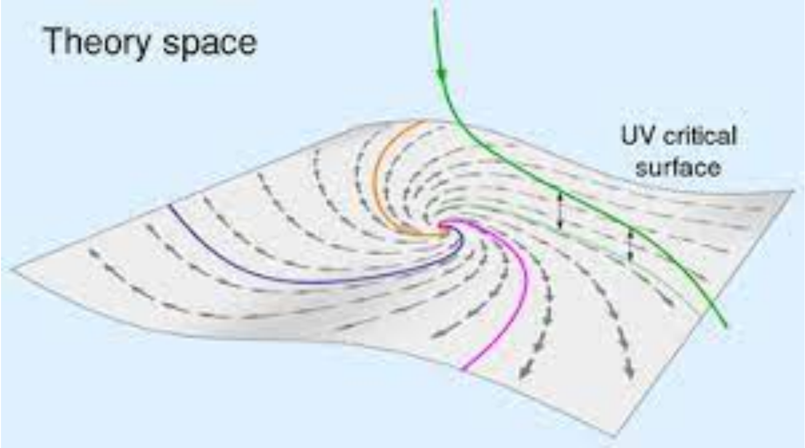
ULTRAVIOLET DIVERGENCES IN QUANTUM THEORIES OF GRAVITATION

•Steven Weinberg(Harvard U.

•)

198042 pages

Part of General Relativity : An Einstein Centenary Survey, 790-831



Asymptotic Safety

Consistent Quantum Gravity if there is a UV fixed point?

Gravitino cosmology

Cosmological Constraints on the Scale of Supersymmetry Breaking

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The gravitino must be either light enough so that ambient gravitinos would not produce too large a cosmic deceleration, or heavy enough so that almost all gravitinos would have decayed before the time of helium synthesis. The second alternative is shown to allow supersymmetry-breaking scales above a model-dependent lower bound of 10^{11} to 10^{16} GeV.

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Supergravity theories[1] necessarily involve a massive spin-$\frac{3}{2}$ particle, the gravitino, whose mass m_g is related to the scale F of spontaneous supersymmetry breaking by the formula[2]

$$m_g = (4\pi/3)^{1/2} F/m_{\rm Pl}, \tag{1}$$

where $m_{\rm Pl}$ is the Planck mass, 1.2×10^{19} GeV. A recent Letter by Pagels and Primack[3] makes the interesting point that the upper bound on the cosmological mass density requires that m_g be less than 1 keV, leading to the upper bound $\sqrt{F} < 2\times10^{6}$ GeV on the scale of spontaneous supersymmetry breaking. This is an important conclusion, because it would mean that supersymmetry, if valid at all, remains unbroken down to energies far below those of order 10^{15} GeV, at which gauge symmetries connecting the strong and electroweak interactions are generally supposed to be broken.[4]

As recognized in Ref. 3, this conclusion applies only if the gravitino is stable enough to survive to the present. It is assumed in Ref. 3 that the gravitino is kept stable by a discrete reflection symmetry,[5] known as "R parity." As conventionally defined, the R parity is even for quarks, leptons, and gauge and Higgs bosons, and odd for their superpartners. The supercurrent is manifestly odd under R parity, so that the gravitino is R odd, and is presumed to be the lightest R-odd particle. In this case, if R parity is conserved, the gravitino may be expected to be absolutely stable.

In this note I wish to examine whether a supersymmetry breakdown at a very high energy such as 10^{15} GeV is really ruled out by the arguments of Ref. 3. First, although R parity is automatically conserved in a wide class of supersymmetric theories, supersymmetric theories do exist in which R parity is not conserved, or at least not with R-parity assignments that would prohibit gravitino decay. Also, even if R parity is an exact symmetry of the Lagrangian, it might be spontaneously broken by whatever mechanism breaks supersymmetry.[6] Further, even if R-parity conservation is exact and not spontaneously broken, how do we know that the gravitino is the lightest R-odd particle? Finally, even if R-parity conservation were exact and not spontaneously broken, and the gravitino were the lightest R-odd particle and consequently absolutely stable, one must still consider whether the annihilation of heavy gravitino pairs might reduce the gravitino mass density to acceptable levels, thus allowing gravitino masses above some lower bound, as is the case for heavy neutrinos.[7]

Let us first dispose of the last issue. Even if R-parity conservation is exact and unbroken, and if the gravitino is the lightest R-odd particle, gravitinos can disappear through annihilation of gravitino pairs, say into $\nu\bar{\nu}$ or $\gamma\gamma$ pairs. Relativistic helicity-($\pm\frac{1}{2}$) gravitinos behave essentially like massless spin-$\frac{1}{2}$ Goldstone fermions,[8] and so these gravitinos annihilate readily at temperatures above m_g, but at these temperatures gravitino pairs are equally readily created in collisions of other particles. Annihilation can only reduce the gravitino population at temperatures below m_g, where the gravitinos are nonrelativistic. At such temperatures the couplings of helicity-($\frac{1}{2}$) as well as -($\pm\frac{3}{2}$) gravitinos are suppressed by powers of $\sqrt{G} = 1/m_{\rm Pl}$. Further, most of the annihilation will take place at temperatures of order m_g, because the thermal average $\langle\sigma v\rangle$ of the product of the gravitino-gravitino annihilation cross section and relative velocity becomes constant for late times, so that the annihilation rate goes like $n_g \propto R^{-3} \propto t^{-2}$, while the cosmic expansion rate $\dot{R}/R$ goes like $(Gm_g n_g)^{1/2} \propto R^{-3/2} \propto t^{-1}$. (This argument can be made more precise by solving the Boltzmann equation for the gravitino number density n_g, as done in Ref. 7.) The dominant contribution to $\langle\sigma v\rangle$ is provided by the conversion of a gravitino pair into a single virtual graviton which then converts into any sufficiently light particle-

Supergravity (hidden sector…)

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Does Gravitation Resolve the Ambiguity among Supersymmetric Vacua?

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Globally supersymmetric theories often have several degenerate supersymmetric vacua. Gravitation splits this degeneracy in such a way that at most one of these vacuum solutions has energy density and cosmological constant equal to zero, while all the rest have negative energy density. Nevertheless, the vacuum with vanishing energy density is stable against decay into the others.

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It is common in supersymmetric theories to find several degenerate vacuum states in which supersymmetry is unbroken and gauge or global symmetries are broken in different ways. For instance, in a supersymmetric SU(N) gauge theory with a single left chiral superfield in the adjoint representation, there are various degenerate vacua in which supersymmetry is unbroken and SU(N) is broken down to $SU(M)\otimes SU(N-M)\otimes U(1)$ or to $SU(N-1)\otimes U(1)$ or not at all. Of course, in the real world supersymmetry *is* broken, but the vacuum ambiguity is nevertheless important for superunified theories in which the scale $\sqrt{F}$ of supersymmetry breaking is much less than the scale M at which the grand gauge group is broken. These ambiguities are not removed by higher-order corrections to the vacuum energy.[1]

One may hope that these ambiguities would be resolved when globally supersymmetric theories are coupled to gravitation. (For theories with scalar field expectation values of order $M \approx 10^{15}$–10^{17} GeV, the gravitational terms in the vacuum energy will be of order GM^6, which is much greater than the energy $F^2/2$ associated with supersymmetry breaking if $\sqrt{F} \ll 10^{13}$ GeV.) It will be shown here that this hope is only partly fulfilled—the different supersymmetric vacua are no longer degenerate, and only one is likely to have vanishing cosmological constant, but most or all of them are stable.

Before we consider the effects of gravitation, it will be useful to recall the reason why there tend to be several degenerate vacua in globally supersymmetric theories. The potential in such theories has the general form[2]

$$V = J_{ab}^{-1}(z,z^*)F_a(z)F_b(z)^* + \tfrac{1}{2}\textstyle\sum_A |D_A(z,z^*)|^2, \qquad (1)$$

where

$$J_{ab}(z,z^*) = \partial^2 d(z,z^*)/\partial z^a \partial z^{b*}, \qquad (2)$$

$$F_a(z) = \partial f(z)/\partial z^a, \qquad (3)$$

$$D_A(z,z^*) = [\partial\, d(z,z^*)/\partial z^a](t_A)_{ab} z^b. \qquad (4)$$

Here z^a are the complex scalar components of left chiral superfields S^a; t_A is the representation of the Ath gauge generator on these scalars, including a coupling-constant factor; and $f(S)$ and $d(S,S^*)$ are the arbitrary functions whose $\theta_L\theta_L$ and $\theta_L\theta_L\theta_R\theta_R$ terms (F and D terms) appear in the Lagrangian. For renormalizable theories the superpotential $f(z)$ is a cubic polynomial and $d(z,z^*)=|z|^2$; a general d function is included here because renormalizability will not be maintained when we include the effects of gravitation. (With such a general d function, the kinetic term for the scalars is $-J_{ab}\partial_\mu z^a \partial^\mu z^{b*}$.) The gauge group is assumed here to be semisimple, so no Fayet-Iliopoulos terms[3] appear in D_A. A supersymmetric solution is one for which all F_a and all D_A vanish.

Any gauge group G will generally have several "big" subgroups H_n, with the property that if G is spontaneously broken to H_n, then all D_A vanish solely as a consequence of the remaining symmetries in H_n. Inspection of Eq. (4) shows that all D_A vanish if there are no broken generators of G that are neutral under H_n. For example, this is the case if G is SU(N) and H_n is either SU(N) itself, or $SU(N-1)\otimes U(1)$, or $SU(M)\otimes SU(N-M)\otimes U(1)$, or $SU(L)\otimes SU(M)\otimes SU(N-M-L)\otimes U(1)\otimes U(1)$, etc., irrespective of the representations of SU(N) provided by the chiral scalar superfields of the theory. Suppose we constrain the scalar field expectation values to be invariant under any one of such big subgroups H_n. The conditions for a vacuum with unbroken supersymmetry are then just $F_a = 0$. Now, the constraint

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Supergravity as the messenger of supersymmetry breaking

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A systematic study is made of theories in which supergravity is spontaneously broken in a "hidden" sector of superfields that interact with ordinary matter only through supergravity. General rules are given for calculating the low-energy effective potential in such theories. This potential is given as the sum of ordinary supersymmetric terms involving a low-energy effective superpotential whose mass terms arise from integrating out the heavy particles associated with grand unification, plus supersymmetry-breaking terms that depend on the details of the hidden sector and the Kähler potential only through the values of four small complex mass parameters. The result is not the same as would be obtained by ignoring grand unification and inserting small mass parameters into the superpotential from the beginning. The general results are applied to a class of models with a pair of Higgs doublets.

I. INTRODUCTION

It was widely hoped that supersymmetry would turn out to be spontaneously broken at energies no higher than a few hundred GeV, both in order to help in understanding gauge hierarchies and also to allow some chance of confirming supersymmetry experimentally. Unhappily, it has proved difficult to construct satisfactory theories along these lines.[1] We are led to the conclusion that supersymmetry if valid at all is spontaneously broken at energies very much greater than those of $SU(2)\times U(1)$ breaking. But then if any vestige of supersymmetry is to survive at ordinary energies to help establish a gauge hierarchy, the source of supersymmetry breaking must somehow be partly isolated from ordinary particles and interactions.[2]

Recently attention has been drawn to a class of interesting models of this sort.[3–13] In these models, unextended ($N=1$) supersymmetry is broken by very large scalar-field vacuum expectation values (VEV's) of order 10^{19} GeV, but the scalars that have these large VEV's form a "hidden sector," that does not interact directly with the ordinary fields (quarks, leptons, gauge and Higgs bosons, and their superpartners) of the "observable sector." That is, the superpotential of the theory breaks up into a sum of two terms[14,15]

$$f_{\mathrm{TOTAL}}(S,\tilde{S}) = f(S) + \tilde{f}(\tilde{S}),$$

where S^a and $\tilde{S}^h$ are the left-chiral superfields of the observable and hidden sectors, respectively. With a minimal kinetic term and no other interactions, the potential of the scalar (nonauxiliary) components $z^a, \tilde{z}^h$ of $S^a, \tilde{S}^h$ would take the form

$$V(z,\tilde{z}) = \sum_{\text{all } z} \left|\frac{\partial f_{\mathrm{TOTAL}}}{\partial z}\right|^2 = \sum_a \left|\frac{\partial f(z)}{\partial z^a}\right|^2 + \sum_h \left|\frac{\partial \tilde{f}(\tilde{z})}{\partial \tilde{z}^h}\right|^2$$

and the spontaneous breakdown of supersymmetry in the hidden sector could have no effect on the observable sector. In the models of Refs. 3–12 the news that supersymmetry is broken by the $\tilde{z}^h$ VEV's is carried over to the observable superfields by gravity and its superpartners, which interact with both sectors.

In the papers of Ref. 3, a thorough study is presented of a model with a specific linear hidden-sector superpotential $\tilde{f}$, and a specific grand-unified observable sector. Their results exhibit some remarkable features; in particular, the VEV's of the light Higgs scalars are of order of the gravitino mass m_g, and do not depend in any way on the grand-unified mass scale M_{GU}, but do depend on coupling parameters of heavy fields whose masses are of order M_{GU}. However, because the model studied was so specific, and the results were expressed in terms of values for scalar VEV's, it was difficult to see how the decoupling of heavy from light degrees of

Superconductivity explained Using particle physics techniques

Superconductivity for Particular Theorists*)

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No one did more than Nambu to bring the idea of spontaneously broken symmetries to the attention of elementary particle physicists. And, as he acknowledged in his ground-breaking 1960 article "Axial Current Conservation in Weak Interactions", Nambu was guided in this work by an analogy with the theory of superconductivity, to which Nambu himself had made important contributions. It therefore seems appropriate to honor Nambu on his birthday with a little pedagogical essay on superconductivity, whose inspiration comes from experience with broken symmetries in particle theory. I doubt if anything in this article will be new to the experts, least of all to Nambu, but perhaps it may help others, who like myself are more at home at high energy than at low temperature, to appreciate the lessons of superconductivity.

§ 1. Introduction

There is something peculiar about standard textbook treatments of superconductivity. On one hand, the reader learns that superconductors exhibit startling phenomena that can be predicted with extraordinary accuracy. For instance, electrical resistance is so low that currents can circulate for years without perceptible decay; the magnetic flux through a loop of such currents or through a vortex line in a superconductor is quantized to high precision; and a junction between two superconductors at different voltage produces an alternating current with a frequency that is so accurately predicted that it provides the best experimental value of the fundamental constant $e/\hbar$. But in deriving these results, textbooks generally use models that — despite their great historical importance — are surely no better than reasonably good approximations. There are macroscopic models, like that of Ginzburg and Landau, in which cooperative states of electrons are represented with a complex scalar field. And there is the microscopic model of Bardeen, Cooper and Schrieffer, from which the Ginzburg-Landau theory can be derived, and in which electrons appear explicitly, but are assumed to interact only by single-phonon exchange. How can one possibly use such approximations to derive predictions about superconducting phenomena that are of essentially unlimited accuracy?

The answer to this question is well understood by experts in superconductivity. When pressed, they will explain that the high-precision predictions about superconductors actually follow not only from the models themselves, but more generally from the fact that these models exhibit a spontaneous breakdown of electromagnetic gauge invariance in a superconductor. The importance of broken symmetry in superconductivity has been

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Cosmological constant problem

$$R_{\mu\nu} - \frac{1}{2} R g_{\mu\nu} = 8\pi G \langle T_{\mu\nu} \rangle_{\rm vac} \sim \Lambda g_{\mu\nu}$$

$$\Lambda \approx (10^{-3}{\rm eV})^4$$

$$\frac{\Lambda}{M_P^4} \sim 10^{-123} \ll 1 \qquad M_P \sim 10^{19}{\rm GeV}$$

The cosmological constant problem*

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Astronomical observations indicate that the cosmological constant is many orders of magnitude smaller than estimated in modern theories of elementary particles. After a brief review of the history of this problem, five different approaches to its solution are described.

CONTENTS

As I was going up the stair,
I met a man who wasn't there.
He wasn't there again today,
I wish, I wish he'd stay away.

Hughes Mearns

I. INTRODUCTION

Physics thrives on crisis. We all recall the great progress made while finding a way out of various crises of the past: the failure to detect a motion of the Earth through the ether, the discovery of the continuous spectrum of beta decay, the τ-θ problem, the ultraviolet divergences in electromagnetic and then weak interactions, and so on. Unfortunately, we have run short of crises lately. The "standard model" of electroweak and strong interactions currently faces neither internal inconsistencies nor conflicts with experiment. It has plenty of loose ends; we know no reason why the quarks and leptons should have the masses they have, but then we know no reason why they should not.

Perhaps it is for want of other crises to worry about that interest is increasingly centered on one veritable crisis: theoretical expectations for the cosmological constant exceed observational limits by some 120 orders of magnitude.[1] In these lectures I will first review the history of this problem and then survey the various attempts that have been made at a solution.

*Morris Loeb Lectures in Physics, Harvard University, May 2, 3, 5, and 10, 1988.

[1]For a good nonmathematical description of the cosmological constant problem, see Abbott (1988).

II. EARLY HISTORY

After completing his formulation of general relativity in 1915–1916, Einstein (1917) attempted to apply his new theory to the whole universe. His guiding principle was that the universe is static: "The most important fact that we draw from experience is that the relative velocities of the stars are very small as compared with the velocity of light." No such static solution of his original equations could be found (any more than for Newtonian gravitation), so he modified them by adding a new term involving a free parameter λ, the cosmological constant:[2]

$$R_{\mu\nu} - \tfrac{1}{2} g_{\mu\nu} R - \lambda g_{\mu\nu} = -8\pi G T_{\mu\nu} \; . \tag{2.1}$$

Now, for $\lambda > 0$, there was a static solution for a universe filled with dust of zero pressure and mass density

$$\rho = \frac{\lambda}{8\pi G} \; . \tag{2.2}$$

Its geometry was that of a sphere S_3, with proper circumference $2\pi r$, where

$$r = 1/\sqrt{8\pi\rho G} \; , \tag{2.3}$$

so the mass of the universe was

$$M = 2\pi^2 r^3 \rho = \frac{\pi}{4} \lambda^{-1/2} G^{-1} \; . \tag{2.4}$$

In some popular history accounts, it was Hubble's discovery of the expansion of the universe that led Einstein to retract his proposal of a cosmological constant. The real story is more complicated, and more interesting.

One disappointment came almost immediately. Einstein had been pleased at the connection in his model between the mass density of the universe and its geometry, because, following Mach's lead, he expected that the mass distribution of the universe should set inertial frames. It was therefore unpleasant when his friend de Sitter, with whom Einstein remained in touch during the war, in 1917 proposed another apparently static cosmological model with no matter at all. (See de Sitter, 1917.) Its line element (using the same coordinate system as de Sitter, but in a different notation) was

$$d\tau^2 = \frac{1}{\cosh^2 Hr} [dt^2 - dr^2 - H^{-2} \tanh^2 Hr (d\theta^2 + \sin^2\theta \, d\varphi^2)] \; , \tag{2.5}$$

[2]The notation used here for metrics, curvatures, etc., is the same as in Weinberg (1972).

Anthropic prediction of the value of the cosmological constant (dark energy) !!

Anthropic Bound on the Cosmological Constant

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In recent cosmological models, there is an "anthropic" upper bound on the cosmological constant Λ. It is argued here that in universes that do not recollapse, the only such bound on Λ is that it should not be so large as to prevent the formation of gravitationally bound states. It turns out that the bound is quite large. A cosmological constant that is within 1 or 2 orders of magnitude of its upper bound would help with the missing-mass and age problems, but may be ruled out by galaxy number counts. If so, we may conclude that anthropic considerations do not explain the smallness of the cosmological constant.

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Our knowledge of the present expansion rate of the Universe indicates that the effective value Λ of the cosmological constant is vastly less than what would be produced by quantum fluctuations[1] in any known realistic theory of elementary particles. In view of the continued failure to find a microscopic explanation of the smallness of the cosmological constant, it seems worthwhile to look for a solution in other, "anthropic," directions.[2] Perhaps Λ must be small enough to allow the Universe to evolve to its present nearly empty and flat state, because otherwise there would be no scientists to worry about it. Without having a definite framework for such reasoning, one can at least point to four lines of current cosmological speculation, in which anthropic considerations could set bounds on the value we observe for the effective cosmological constant:

(a) The effective cosmological constant may evolve very slowly, perhaps because of slow changes in the value of some scalar field, as in the model of Banks.[3] In this case, it would be natural to expect that for some very long epoch the cosmological constant would remain near zero. The question then is, why do we find ourselves in such an epoch? As remarked by Banks, the answer may be anthropic: Perhaps only in such epochs is life possible.

(b) The Universe may evolve through a very large number of first-order phase transitions, in which bubbles form within bubbles within bubbles . . . , each bubble having within it a smaller value of the vacuum energy, and hence of the effective cosmological constant. If the steps in vacuum energy are very small, then it would be natural to expect that there would be some phase in which the effective cosmological constant is correspondingly small. Abbott[4] has suggested a scalar-field theory with a potential that has an infinite number of closely spaced local minima; bubbles form within bubbles as the scalar-field value jumps from one minimum to the next. Recently Brown and Teitelboim[5] have proposed a model in which a similar sequence of phase transitions occurs, but in which the bubble walls are elementary membranes coupled to a three-form gauge field, with the difference in cosmological constants between the inside and outside of each membrane caused by the differences in the values of the four-form field strength.[6] In models of the type discussed in Refs. 4 and 5 it may not be strictly necessary to invoke the anthropic principle because gravitational effects can stop the process of bubble formation when the vacuum energy is about to become negative.[7] However, it takes an enormously long time to reach this final stage, and anthropic arguments may be needed to explain why we are not still in an earlier stage, with large effective cosmological constant.

(c) Fluctuations in scalar fields can trigger cosmic inflation in regions of the Universe where the fields happen to be large. Except near the edges, the inflationary region would appear to its inhabitants as a separate subuniverse. In this region further fluctuations can produce new inflations, and so on. This has been studied by Linde,[8] who remarks that the physical constants of the subuniverse in which we live may be in part constrained by the requirement that life could arise in such a subuniverse.

(d) Quantum fluctuations in the very early Universe may cause incoherence between different terms in the state vector of the Universe; each term would then in effect represent a separate universe. Such a picture has been considered by Hawking.[9] Our own Universe could correspond to any one of the terms in the state vector, subject only to the anthropic condition, that it be a universe in which life could develop.

Without committing ourselves to any one of these cosmological models, it seems appropriate at least to ask, just what limit does the anthropic principle place on the effective cosmological constant Λ?

Fortunately, at least for $\Lambda > 0$, the anthropic principle provides a rather sharp upper bound on Λ. This is because in a continually expanding universe, the cosmological constant (unlike charges, masses, etc.) can affect the evolution of life in only one way. Without undue anthropocentrism, it seems safe to assume that in order for any sort of life to arise in an initially homogeneous and isotropic universe, it is necessary for sufficiently large gravi-

Experimental tests of Quantum Mechanics!

ANNALS OF PHYSICS 194, 336–386 (1989)

Testing Quantum Mechanics

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Received March 6, 1989

This paper presents a general framework for introducing nonlinear corrections into ordinary quantum mechanics, that can serve as a guide to experiments that would be sensitive to such corrections. In the class of generalized theories described here, the equations that determine the time-dependence of the wave function are no longer linear, but are of Hamiltonian type. Also, wave functions that differ by a constant factor represent the same physical state and satisfy the same time-dependence equations. As a result, there is no difficulty in combining separated subsystems. Prescriptions are given for determining the states in which observables have definite values and for calculating the expectation values of observables for general states, but the calculation of probabilities requires detailed analysis of the method of measurement. A study is presented of various experimental possibilities, including the precession of spinning particles in external fields, experiments of Stern–Gerlach type, and the broadening and de-tuning of absorption lines.

CONTENTS

1. INTRODUCTION AND SUMMARY

Considering the pervasive importance of quantum mechanics in modern physics, it is odd how rarely one hears of efforts to test quantum mechanics experimentally with high precision. It is true that over the last decade there have been a number of experimental tests [1] of predictions that distinguish quantum mechanics from

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1990’s

Effective field theories in nuclear physics !!

Nuclear forces from chiral lagrangians

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The method of phenomenological lagrangians is used to derive the consequences of spontaneously broken chiral symmetry for the forces among two or more nucleons.

The forces among nucleons have been studied as much as anything in physics. Much of this work has necessarily been phenomenological: scattering data and deuteron properties are used to determine a two-nucleon interaction, which can then be used as an input to multi-nucleon calculations. As more and more has been learned about the meson spectrum, efforts have been increasingly aimed at calculating the nuclear potential as an expansion in terms of decreasing range arising from the exchange of one or more mesons of various types, but the number of free parameters rises rapidly as more and more meson types are included, especially if one attempts to extend these calculations to forces involving more than two nucleons. This paper applies methods [1] based on the chiral symmetry of quantum chromodynamics to derive an expansion of the potential among any number of low energy nucleons in powers of the nucleon momenta, which is related to but not identical with the expansion in terms of increasing range. It is not clear which expansion will be more useful in dealing with the two-nucleon problem, but the expansion in powers of momenta gives far more specific information about multi-nucleon potentials.

The lagrangian that we shall use in this work will be taken as the most general possible lagrangian involving pions and low-energy nucleons consistent with spontaneously broken chiral symmetry and other known symmetries. It is given by an infinite series of terms with increasing numbers of derivatives and/or nucleon fields, with the dependence of each term on the pion field prescribed by the rules of broken chiral symmetry. Other degrees of freedom, such as heavy vector mesons, Δ's, and antinucleons, are "integrated out": their contribution is buried in the coefficients of the series of terms in the pion–nucleon lagrangian. We shall also integrate out nucleons with momenta greater than some scale Q, which requires that these coefficients in the lagrangian be Q-dependent. Later we will consider how to make a judicious choice of Q; for the moment, it will be enough to specify that Q is substantially less than m_ρ. Any detailed model such as that of Skyrme [2] (also see ref. [3]) that embodies broken chiral symmetry will give results that are consistent with ours, but less general; in particular, we do not specify any particular higher-derivative terms in the lagrangian such as those that are introduced to stabilize skyrmions, but instead we consider all possible terms, with any numbers of derivatives, that are allowed by the symmetries of strong interactions.

Now consider the S-matrix for a scattering process with N incoming and N outgoing nucleons, all with momenta no larger than Q. The non-relativistic nature of the problem makes it appropriate to apply "old-fashioned" time-ordered perturbation theory: there is an energy denominator for every intermediate state, instead of a propagator for every internal particle line. The energy denominators associated with intermediate states involving just N nucleons are small, of order $Q^2/2m_N$, as compared with Q for the

[1] Research supported in part by the Robert A. Welch Foundation and NSF Grant PHY 9009850.

The Cosmological Constant and probabilities!!

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LIKELY VALUES OF THE COSMOLOGICAL CONSTANT

Hugo Martel[1,2], Paul R. Shapiro[1,3], and Steven Weinberg[4,5]

ABSTRACT

In theories in which the cosmological constant takes a variety of values in different "subuniverses," the probability distribution of its observed values is conditioned by the requirement that there must be someone to measure it. This probability is proportional to the fraction of matter which is destined to condense out of the background into mass concentrations large enough to form observers. We calculate this "collapsed fraction" by a simple, pressure-free, spherically symmetric, nonlinear model for the growth of density fluctuations in a flat universe with arbitrary value of the cosmological constant, applied in a statistical way to the observed spectrum of density fluctuations at recombination. From this, the probability distribution for the vacuum energy density ρ_V for Gaussian random density fluctuations is derived analytically. (The conventional quantity λ_0 is the vacuum energy density in units of the critical density at present, $\lambda_0 = \rho_V/\rho_{\mathrm{crit},0}$, where $\rho_{\mathrm{crit},0} = 3H_0^2/8\pi G$.) It is shown that the results depend on only one quantity, $\sigma^3\bar{\rho}$, where σ^2 and $\bar{\rho}$ are the variance and mean value of the fluctuating matter density field at recombination, respectively. To calculate σ, we adopt the flat CDM model with nonzero cosmological constant and fix the amplitude and shape of the primordial power spectrum in accordance with data on cosmic microwave background anisotropy from the COBE satellite DMR experiment. A comparison of the results of this calculation of the likely values of ρ_V with present observational bounds on the cosmological constant indicates that the small, positive value of ρ_V (up to 3 times greater than the present cosmic mass density)

2000's

Efective theories in Cosmology!

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Quantum contributions to cosmological correlations

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The "in-in" formalism is reviewed and extended, and applied to the calculation of higher-order Gaussian and non-Gaussian correlations in cosmology. Previous calculations of these correlations amounted to the evaluation of tree graphs in the in-in formalism; here we also consider loop graphs. It turns out that for some though not all theories, the contributions of loop graphs as well as tree graphs depend only on the behavior of the inflaton potential near the time of horizon exit. A sample one-loop calculation is presented.

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I. INTRODUCTION

The departures from cosmological homogeneity and isotropy observed in the cosmic microwave background and large scale structure are small, so it is natural that they should be dominated by a Gaussian probability distribution, with bilinear averages given by the terms in the Lagrangian that are quadratic in perturbations. Nevertheless, there is growing interest in the possibility of observing non-Gaussian terms in various correlation functions [1], such as an expectation value of a product of three temperature fluctuations. It is also important to understand the higher-order corrections to bilinear correlation functions, which appear in Gaussian correlations.

Until now, higher-order cosmological correlations have been calculated by solving the classical field equations beyond the linear approximation. As will be shown in the Appendix, this is equivalent to calculating sums of tree graphs, though in a formalism different from the familiar Feynman graph formalism. For instance, Maldacena [2] has calculated the non-Gaussian average of a product of three scalar and/or gravitational fields to first order in their interactions, which amounts to calculating a tree graph consisting of a single vertex with 3 attached gravitational and/or scalar field lines.

This paper will discuss how calculations of cosmological correlations can be carried to arbitrary orders of perturbation theory, including the quantum effects represented by loop graphs. So far, loop corrections to correlation functions appear to be much too small ever to be observed. The present work is motivated by the opinion that we ought to understand what our theories entail, even where in practice its predictions cannot be verified experimentally, just as field theorists in the 1940s and 1950s took pains to understand quantum electrodynamics to all orders of perturbation theory, even though it was only possible to verify results in the first few orders.

There is a particular question that will concern us. In the familiar calculations of lowest-order Gaussian correlations, and also in Maldacena's tree-graph calculation of non-Gaussian correlations, the results depended only on the behavior of the unperturbed inflaton field near the time of horizon exit. Is the same true for loop graphs? If so, it will be possible to calculate the loop contributions with some confidence, but we can learn little new from such calculations. On the other hand, if the contribution of loop graphs depends on the whole history of the unperturbed inflaton field, then calculations become much more difficult, but potentially more revealing. In this case, it might even be that the loop contributions are much larger than otherwise expected.

The appropriate formalism for dealing with this sort of problem is the "in-in" formalism originally due to Schwinger [3]. Schwinger's presentation is somewhat opaque, so this formalism is outlined (and extended) in an Appendix. In Sec. II we summarize those aspects of this formalism that are needed for our present purposes. Section III introduces a class of theories to serve as a basis of discussion, with a single inflaton field, plus any number of additional massless scalar fields with only gravitational interactions and vanishing expectation values. In Sec. IV we prove a general theorem about the late-time behavior of cosmological correlations at fixed internal as well as external wave numbers. Section V introduces a class of unrealistic theories to illustrate the problems raised by the integration over internal wave numbers, and how these problems may be circumvented. In Sec. VI we return to the theories introduced in Sec. III, and we show that the conditions of the theorem proved in Sec. IV are satisfied for these theories. This means that, to all orders of perturbation theory, if ultraviolet divergences cancel in the integrals over internal wave numbers, then cosmological correlations do indeed depend only on the behavior of the unperturbed inflation field near the time of horizon exit in the cases studied. We can also find other theories in which this result does not apply, as for instance by giving the additional scalar fields a self-interaction. Section VII presents a sample one-loop calculation of a cosmological correlation.

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Effective field theory of inflation

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Effective field theory for inflation

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The methods of effective field theory are used to study generic theories of inflation with a single inflaton field. For scalar modes, the leading corrections to the $\mathcal{R}$ correlation function are found to be purely of the k-inflation type. For tensor modes the leading corrections to the correlation function arise from terms in the action that are quadratic in the curvature, including a parity-violating term that makes the propagation of these modes depend on their helicity. These methods are also briefly applied to nongeneric theories of inflation with an extra shift symmetry, as in so-called ghost inflation.

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I. GENERIC THEORIES OF INFLATION

Observations of the cosmic microwave background and large scale structure are consistent with a simple theory of inflation [1] with a single canonically normalized inflaton field $\varphi_c(x)$, described by a Lagrangian

$$\mathcal{L}_0 = \sqrt{g}\left[-\frac{M_P^2}{2}R - \frac{1}{2}g^{\mu\nu}\partial_\mu\varphi_c\partial_\nu\varphi_c - V(\varphi_c)\right], \quad (1)$$

where $g \equiv -\mathrm{Det}g_{\mu\nu}$, $M_P \equiv 1/\sqrt{8\pi G}$ is the reduced Planck mass, and $V(\varphi_c)$ is a potential down which the scalar field rolls more-or-less slowly. With this theory, the strength of observed fluctuations in the microwave background matter density indicates that the cosmic expansion rate $H \equiv \dot{a}/a$ and the physical wave number k/a at horizon exit, when these are equal, have the value [2] $H = k/a \approx \sqrt{\epsilon} \times 2 \times 10^{14}$ GeV where ϵ is the value of $-\dot{H}/H^2$ at this time, and a is the Robertson–Walker scale factor. Hence H and k/a at horizon exit are likely to be much less than $M_P \simeq 2.4 \times 10^{18}$ GeV, and even considerably less than a plausible grand unification scale $\approx 10^{16}$ GeV. This provides a justification after the fact for using a Lagrangian (1) with a minimum number of spacetime derivatives. [As is well known, (1) is the most general Lagrangian density for gravitation and a single scalar field with no more than two spacetime derivatives. An arbitrary function of φ multiplying the first term could be eliminated by a redefinition of the metric, and an arbitrary function of φ multiplying the second term could be eliminated by a redefinition of φ.]

But H and k/a at horizon exit are not entirely negligible compared with whatever fundamental scale characterizes the theory underlying inflation, and at earlier times k/a is exponentially larger than at horizon exit, so it is worth considering the next corrections to (1). We assume that (1) is just the first term in a generic effective field theory, in which terms with higher derivatives are suppressed by negative powers of some large mass M, characterizing whatever fundamental theory underlies this effective field theory. Rather than committing ourselves to any particular underlying theory, we will simply assume that all constants in the higher derivative terms of the effective Lagrangian take values that are powers of M indicated by dimensional analysis, with coefficients roughly of order unity. Because H and k/a are so large during inflation, observations of fluctuations produced during inflation provide a unique opportunity for detecting effects of higher derivative terms in the gravitational action.

To get some idea of the value of M, we note that the unperturbed canonically normalized scalar field $\bar{\varphi}_c$ described by the Lagrangian (1) has a time derivative $\dot{\bar{\varphi}}_c = \sqrt{2\epsilon}M_P H$, so the change in $\bar{\varphi}_c$ during a Hubble time $1/H$ at around the time of horizon exit is of order $\dot{\bar{\varphi}}_c/H = \sqrt{2\epsilon}M_P$. If we are to use effective field theory to study fluctuations at about the time of horizon exit in generic theories in which the dependence of the action on φ_c is unconstrained by symmetry principles or by other consequences of an underlying theory, and if (1) is at least a fair first approximation to the full theory, then the mass M that is characteristic of the effective field theory of inflation cannot be much smaller than $\sqrt{2\epsilon}M_P$, for if it were then there would be no limit on the size of higher-derivative terms containing many powers of φ_c/M. It follows that the expansion parameter H/M in this class of theories is no greater than $H/\sqrt{2\epsilon}M_P \simeq 6 \times 10^{-5}$, whatever the value of ϵ.

We will tentatively assume here that M is of order $\sqrt{2\epsilon}M_P$, in which case the coefficients of the higher-derivative terms in the effective Lagrangian have to be taken as arbitrary functions of φ_c/M. This is likely to be the case if ϵ is not too small, say of order 0.02, since then there is not much difference between $\sqrt{2\epsilon}M_P$ and M_P, and M is unlikely to be much larger than M_P. (The considerations presented below would still be valid if M were instead much larger than $\sqrt{2\epsilon}M_P$, as for instance if $M \approx M_P$ and ϵ is very small, but then we would have to count powers of φ_c/M as well as numbers of derivatives in judging how much the various higher-derivative terms are

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2010’s

Goldstone Bosons and ΔN_{eff}

Goldstone Bosons as Fractional Cosmic Neutrinos

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It is suggested that Goldstone bosons may be masquerading as fractional cosmic neutrinos, contributing about 0.39 to what is reported as the effective number of neutrino types in the era before recombination. The broken symmetry associated with these Goldstone bosons is further speculated to be the conservation of the particles of dark matter.

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The correlations of temperature fluctuations in the cosmic microwave background depend on the effective number $N_{\rm eff}$ of neutrino species present in the era before recombination. Although observations are certainly consistent with the expected value $N_{\rm eff}=3$, there have been persistent hints in the data that the effective number may be somewhat greater. WMAP9 together with ground-based observations (WMAP9 + eCMB) [1] gave $N_{\rm eff}=3.89\pm 0.67$, while Planck together with the WMAP9 polarization data and ground-based observations (Planck + WP + highL) [2] gives $N_{\rm eff}=3.36\pm 0.34$, both at the 68% confidence level. Is it possible that some nearly massless weakly interacting particle is masquerading as a fractional cosmic neutrino?

As a candidate for an imposter fractional neutrino, one naturally thinks of Goldstone bosons, associated with the spontaneous breakdown of some exact or nearly exact global continuous symmetry. They would, of course, be massless or nearly massless, and the characteristic derivative coupling of Goldstone bosons would make them weakly interacting at sufficiently low temperatures.

Since Fermi statistics reduces the energy density of neutrinos relative to massless bosons by a factor $7/8$, and $N_{\rm eff}$ lumps antineutrinos with neutrinos, a neutral Goldstone boson might look like $(1/2)/(7/8)=4/7$ of a neutrino. But for this to be true, there is an important qualification: the Goldstone bosons must remain in thermal equilibrium with ordinary particles until after the era of muon annihilation, so that the temperature of the Goldstone bosons matches the neutrino temperature. If Goldstone bosons went out of equilibrium much earlier, then neutrinos but not Goldstone bosons would have been heated by the annihilation of the various species of particles of the standard model (SM), and the contribution of Goldstone bosons to $N_{\rm eff}$ would be much less than $4/7$. As we shall see, there is a plausible intermediate possibility that the contribution of Goldstone bosons to $N_{\rm eff}$ would be $(4/7)(43/57)^{4/3}=0.39$. To judge when the Goldstone bosons went out of thermal equilibrium, we need a specific theory [3].

We will consider the simplest possible broken continuous symmetry, a $U(1)$ symmetry associated with the conservation of some quantum number W. All fields of the standard model are supposed to have $W=0$. To allow in the simplest way for the breaking of this symmetry, we introduce a single complex scalar field $\chi(x)$, neutral under $SU(3)\otimes SU(2)\otimes U(1)$, which carries a nonvanishing value of W. With this field added to the standard model, the most general renormalizable Lagrangian is

$$\mathcal{L}=-\frac{1}{2}\partial_\mu\chi^\dagger\partial^\mu\chi+\frac{1}{2}\mu^2\chi^\dagger\chi-\frac{1}{4}\lambda(\chi^\dagger\chi)^2-\frac{g}{4}(\chi^\dagger\chi)(\varphi^\dagger\varphi)+\mathcal{L}_{\rm SM},\tag{1}$$

where μ^2, g, and λ are real constants; $\mathcal{L}_{\rm SM}$ is the usual Lagrangian of the standard model; and $\varphi=(\varphi^0,\varphi^-)$ is the standard model's scalar doublet. Experience with the linear σ model shows that with a Lagrangian like (1), there are several diagrams in each order of perturbation theory that must be added up in order to give matrix elements that agree with theorems governing soft Goldstone bosons. To avoid this, it is better to separate a massless Goldstone boson field $\alpha(x)$ and a massive "radial" field $r(x)$ by defining

$$\chi(x)=r(x)e^{2i\alpha(x)},\tag{2}$$

where $r(x)$ and $\alpha(x)$ are real, with the phase of $\chi(x)$ adjusted to make $\langle\alpha(x)\rangle=0$. (The 2 in the exponent is for future convenience.) The Lagrangian (1) then takes the form

$$\mathcal{L}=-\frac{1}{2}\partial_\mu r\partial^\mu r+\frac{1}{2}\mu^2r^2-\frac{1}{4}\lambda r^4-2r^2\partial_\mu\alpha\partial^\mu\alpha-\frac{g}{4}r^2(\varphi^\dagger\varphi)+\mathcal{L}_{\rm SM}.\tag{3}$$

The $SU(2)\otimes U(1)$ symmetry of the standard model is of course broken by a nonvanishing vacuum expectation value of the field φ^0, with a real zeroth-order value $\langle\varphi\rangle\simeq 247$ GeV. The $U(1)$ symmetry of W conservation is also broken if $(\mu^2-g\langle\varphi\rangle^2)/\lambda$ is positive, in which case r gets a real vacuum expectation value, given in zeroth order by

$$\langle r\rangle=\sqrt{m_r^2/2\lambda},\qquad m_r^2\equiv\mu^2-g\langle\varphi\rangle^2/2.\tag{4}$$

In this formalism, the interaction of Goldstone bosons with the particles of the standard model arises entirely from

2020’s

Last years:

- Models of quark and lepton masses
- Massless particles in extra dimensions
- Foundations of Quantum Mechanics
- Gravitational waves…

Models of lepton and quark masses

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A class of models is considered in which the masses only of the third generation of quarks and leptons arise in the tree approximation, while masses for the second and first generations are produced respectively by one-loop and two-loop radiative corrections. So far, for various reasons, these models are not realistic.

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I. INTRODUCTION

In the Standard Model the masses of quarks and leptons take values proportional to the coupling constants in the interaction of these fermions with scalar fields, constants that in the context of this model are entirely arbitrary. But the peculiar hierarchical pattern of lepton and quark masses seems to call for a larger theory, in which in some leading approximation the only quarks and leptons with nonzero mass are those of the third generation, the tau, top, and bottom, with the other lepton and quark masses arising from some sort of radiative correction. Such theories were actively considered [1] soon after the completion of the Standard Model, but interest in this program seems to have lapsed subsequently [2].

This paper will explore in detail a class of models of this sort, based on a different symmetry group. These models are not realistic, for reasons that will be spelled out later, but it is hoped that they may help to revive interest in this program, and to lay out some of the methods and problems that it confronts.

II. GAUGE AND SCALAR FIELDS

If the spontaneous breakdown of the electroweak symmetry gave masses only to the quarks and leptons of the third generation in the tree approximation, then nothing in the Standard Model would generate masses for the first and second generations in higher orders of perturbation theory. To get masses for the second and first generations by emission and absorption of some sort of gauge bosons, we would need to expand the gauge symmetry group. In order for these masses to be much less than the zeroth order masses of the third generation, we would need the gauge coupling constants to be relatively small, more or less like the electroweak couplings. If these new gauge couplings together with those of the Standard Model all descended from some theory such as a string theory or a unified gauge theory in which they were all equal at some very high energy, then in order to have small couplings at accessible energies the new gauge group would have to be a direct product of simple subgroups with smaller beta functions than for the $SU(3)$ of QCD—that is, most likely only $SO(3)$ and/or $SO(2)$. After some attempts, what seems to work best is $SO_L(3)\otimes SO_R(3)$, with the three generations of left-handed quark and lepton $SU(2)\otimes U(1)$ doublets forming separate representations $(3,1)$ of $SO_L(3)\otimes SO_R(3)$, and the three generations of right-handed quarks and charged leptons furnishing separate representations $(1,3)$. [We label representations of $SO(3)$ by their dimensionality.] Though we shall concentrate on this gauge group, our analysis will deal with problems that would have to be encountered in any attempt to interpret the hierarchy of quark and lepton masses as radiative corrections.

In order for scalar fields to have renormalizable couplings to these quarks and leptons, they would have to form 9 electroweak doublets

$$\begin{pmatrix} \Phi^+_{ia} \\ \Phi^0_{ia} \end{pmatrix}, \tag{1}$$

transforming as $(3,3)$ representations of $SO_L(3)\otimes SO_R(3)$. [Here superscripts indicate charges; subscripts i, j, etc. are $SO_L(3)$ vector indices running over the values 1, 2, 3; subscripts a, b, etc. are $SO_R(3)$ vector indices. also running over the values 1, 2, 3.] Emission and absorption of the corresponding spinless particles also produces radiative corrections to the quark and lepton masses. As we shall see in the next section, while keeping the mass of the Standard Model Higgs boson and the weak coupling constant at their known values, we can take all the other scalar particles and the new vector bosons to be heavy enough to have escaped detection. But the calculation in Sec. IV shows that the radiative corrections to masses do not disappear when the new scalar and vector bosons become very heavy.

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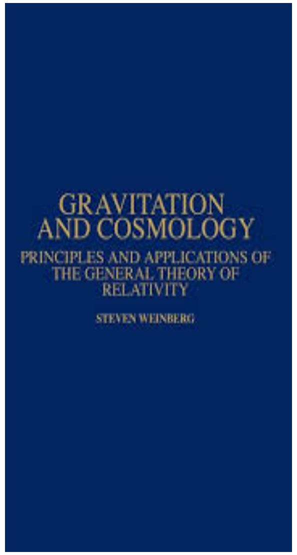
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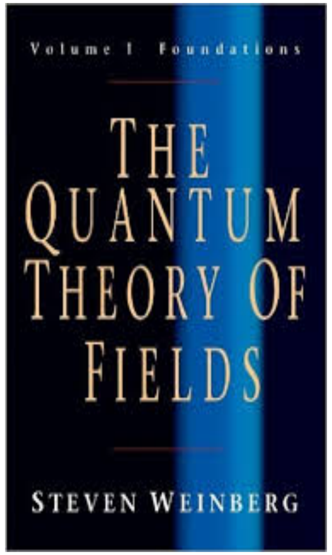
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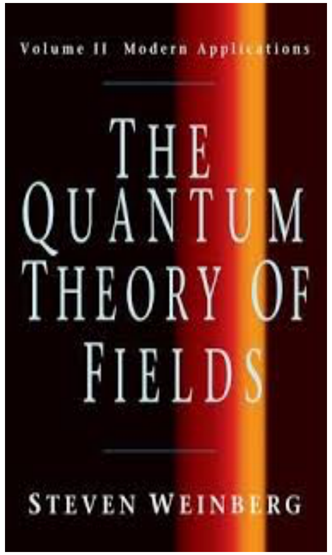
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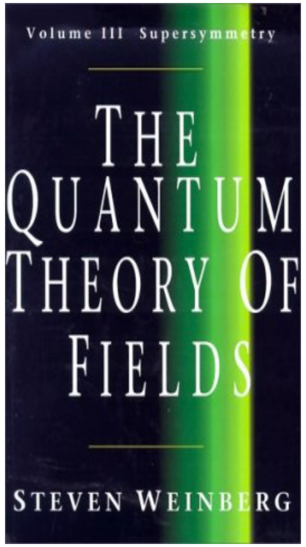
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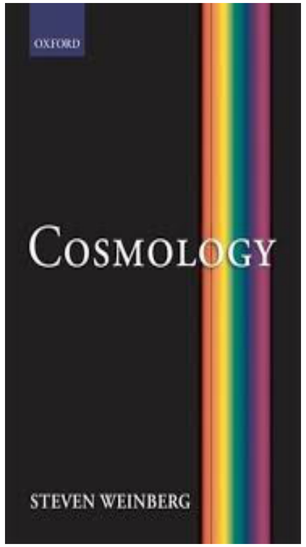
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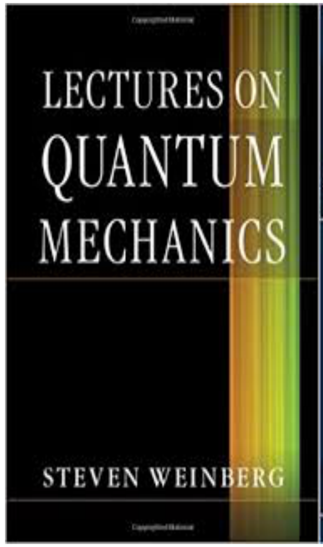
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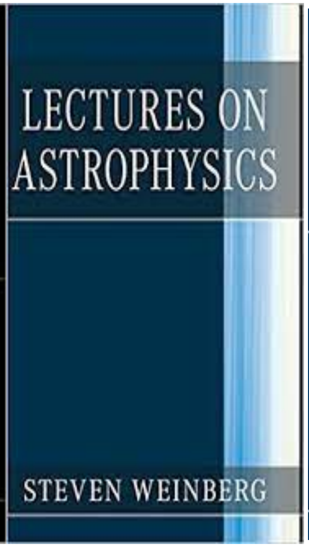
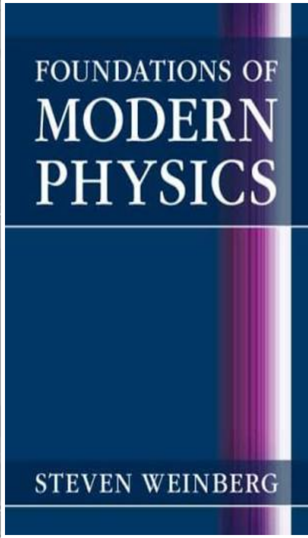
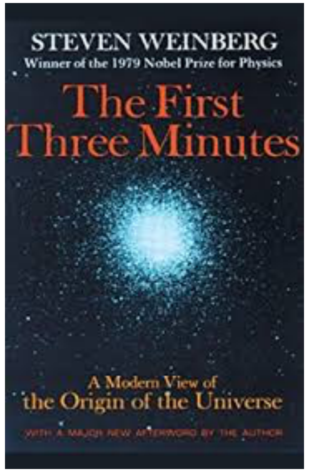
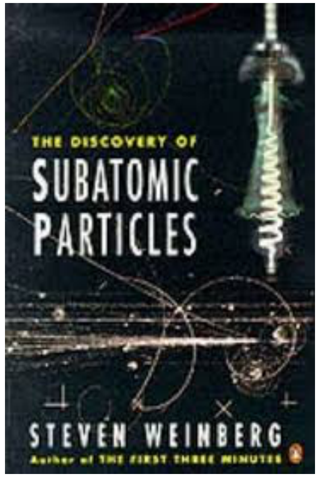
STEVEN WEINBERG

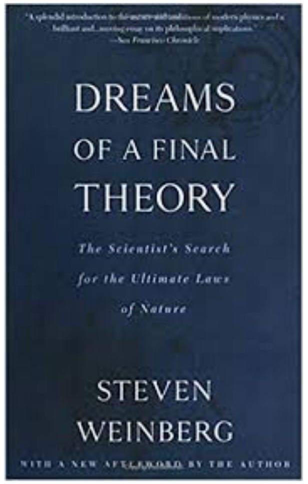

LECTURES ON
ASTROPHYSICS
STEVEN WEINBERG

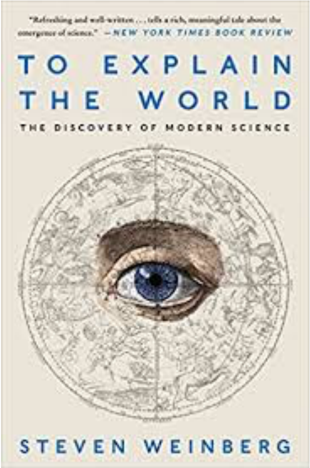

FOUNDATIONS OF
MODERN
PHYSICS
STEVEN WEINBERG

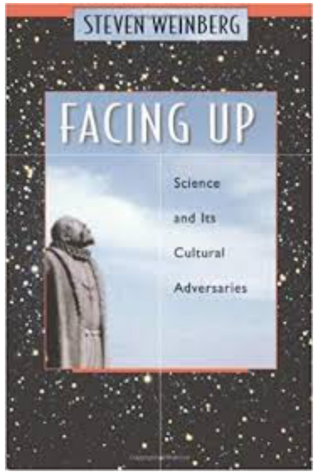
Popular Science Books

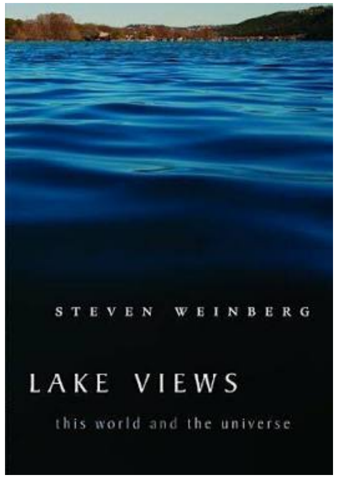


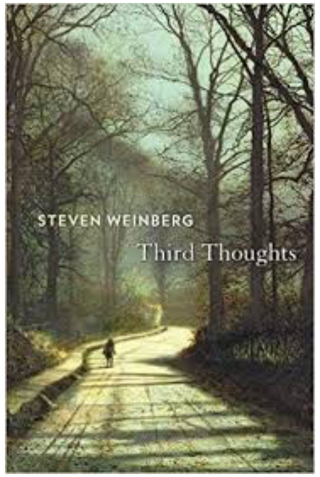


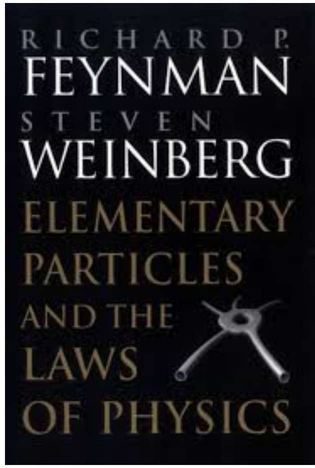












Four Golden Lessons
(advice for graduate students)

- Nobody knows everything and you don’t have to

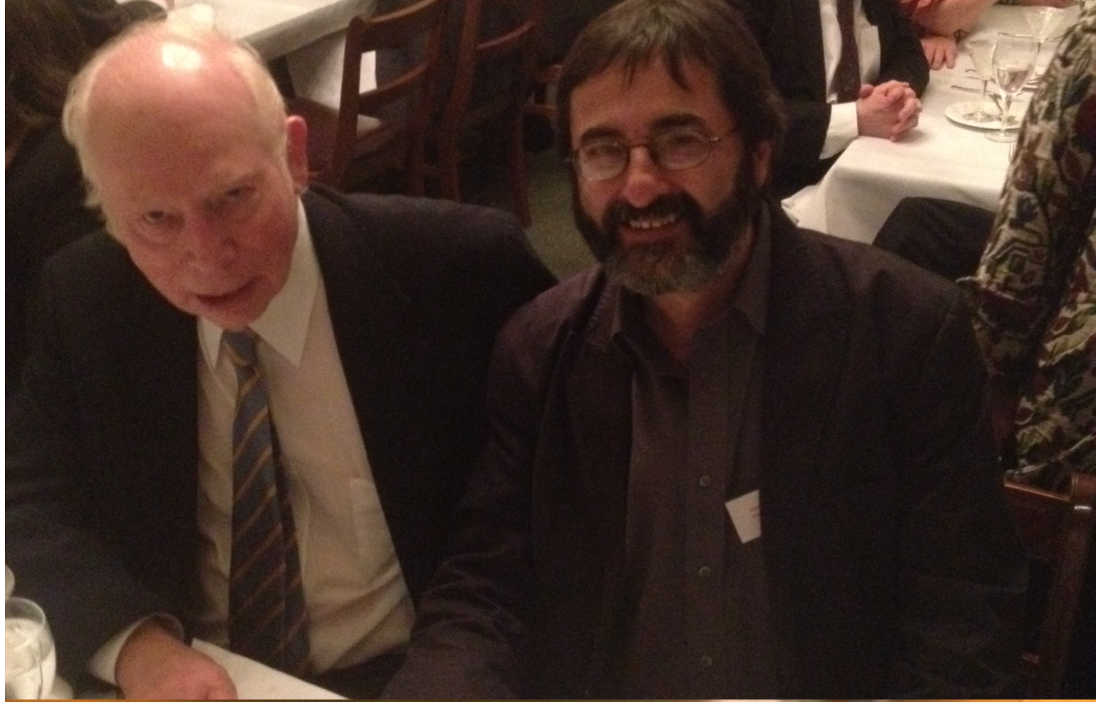
- My advice is to go for the messes — that's where the action is

- To forgive yourself for wasting time

- To learn something about the history of science

Personal Reminiscences

- His lectures
- My thesis advisor
- Support for my career
- Personal anecdotes
- Support to ICTP
- Unexpected collaboration



Joseph Polchinski: A Biographical Memoir

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Joseph Polchinski (1954-2018), one of the the leading theoretical physicists of the past 50 years, was an exceptionally broad and deep thinker. He made fundamental contributions to quantum field theory, advancing the role of the renormalization group, and to cosmology, addressing the cosmological constant problem. Polchinski's work on D-branes revolutionized string theory and led to the discovery of a nonperturbative quantum theory of gravity. His recent, incisive reformulation of the black hole information paradox presents us with a profound challenge. Joe was deeply devoted to his family, a beloved colleague and advisor, an excellent writer, and an accomplished athlete.

I. INTRODUCTION

Joseph Polchinski ranks among the greatest theoretical physicists of his generation. His interests were centered on particle physics and quantum gravity, particularly string theory, and he made epochal contributions in these areas. (Unusually for a theorist, his most revolutionary works came between the ages of 40 and 60.) But he was uncommonly broad. His work substantially impacted a range of fields from cosmology to condensed matter physics.

Polchinski was the very opposite of the caricature of the narrowly focussed theorist, who explores abstruse mathematical structures in an out-of-touch quest for elegance and beauty. He was a full-blooded physicist who cared about understanding Nature, by whatever means he could muster. He was heard saying "I'm a theoretical physicist first, string theorist second." (At the time, shortly after the second superstring revolution that his work had largely triggered, this was not a universal sentiment in the elated community.) As a great pianist is more than a dazzling virtuoso, Polchinski's technical prowess allowed him to perform the most challenging calculations, but always in service of a deeper vision, and in pursuit of a more profound understanding of how the universe works.

When Polchinski came of age as a physicist, the Standard Model of particle physics had recently been worked out. It explained all observed forces but gravity, within the framework of quantum mechanics. The great remaining task since then has been to unify gravity and quantum mechanics. It takes some courage to decide to work on quantum gravity: one expects that the final theory will be simple only at energies and distances that are far out of reach technologically. Many reasonable hypotheses can be probed at best indirectly, through subtle effects.

Against these odds, Joe made profound and lasting contributions to science. From our early (and personal) vantage point,[1] Joe's most significant works include the "string theory landscape" as a solution of the cosmological constant problem, and his contributions to the renormalization group and to the black hole information paradox. And through the breadth and depth of its impact, Polchinski's 1995 discovery of D-branes truly stands out.

D-branes mark a watershed in theoretical physics. Before D-branes, string theory was formulated and understood largely at a perturbative level, as a sum over diagrams. The theory offered a promising approach to quantum gravity—one could compute how gravitons scatter—and to unification, with the hope of deriving the observed symmetries, forces, and particles, from the rigid structure provided by the theory.

But string theory could not be applied to strongly gravitating systems such as black holes, or to cosmology. Thus, many pressing questions could not be addressed: what is the origin of black hole entropy? What happened at the big bang? Polchinski was the first to recognize that the symmetries of string theory require the existence of D-branes, extended objects that are intrinsically non-perturbative.

This insight led to dramatic progress on several deep and long-standing problems. Polchinski took the lead in

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[1] We urge the reader to study Polchinski's own recollections [1], which are truly fascinating. They are refreshingly and poignantly honest about the difficulties he faced in his life and career. As such they can be a valuable resource for all of us, and especially to young people entering the field. Other excellent obituaries and tributes written by physicists include Refs. [2–5].

Some memorable quotes

With or without religion, you would have good people doing good things and evil people doing evil things. But for good people to do evil things, that takes religion..

Science doesn't make it impossible to believe in God, it just makes it possible not to believe in God

If history is any guide at all, it seems to me to suggest that there is a final theory. In this century we have seen a convergence of the arrows of explanation, like the convergence of meridians toward the North Pole.

All logical arguments can be defeated by the simple refusal to reason logically

The effort to understand the universe is one of the very few things which lifts human life a little above the level of farce and gives it some of the grace of tragedy.

The more the universe seems comprehensible, the more it seems pointless.

If there is no point in the universe that we discover by the methods of science, there is a point that we can give the universe by the way we live, by loving each other, by discovering things about nature, by creating works of art. And that — in a way, although we are not the stars in a cosmic drama, if the only drama we're starring in is one that we are making up as we go along, it is not entirely ignoble that faced with this unloving, impersonal universe we make a little island of warmth and love and science and art for ourselves. That's not an entirely despicable role for us to play.

My Favourite Quotes

The Guideline

Our purpose in theoretical physics is not to describe the world as we find it, but to explain-in terms of a few fundamental principles- why the world is the way it is.

Steven Weinberg

…The other possibility, which I have to admit is *a priori* more likely, is that at very high energy we will run into really new physics, not describable in terms of a quantum field theory. I think that by far the most likely possibility is that this will be something like a string theory.

Our mistake is not that we take our theories too seriously, but that we do not take them seriously enough. It is always hard to realize that these numbers and equations we play with at our desks have something to do with the real world.

Organic
WE HAVE 144 ORGANIC VARIETIES
32
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Thank You!!!

Steven Weinberg