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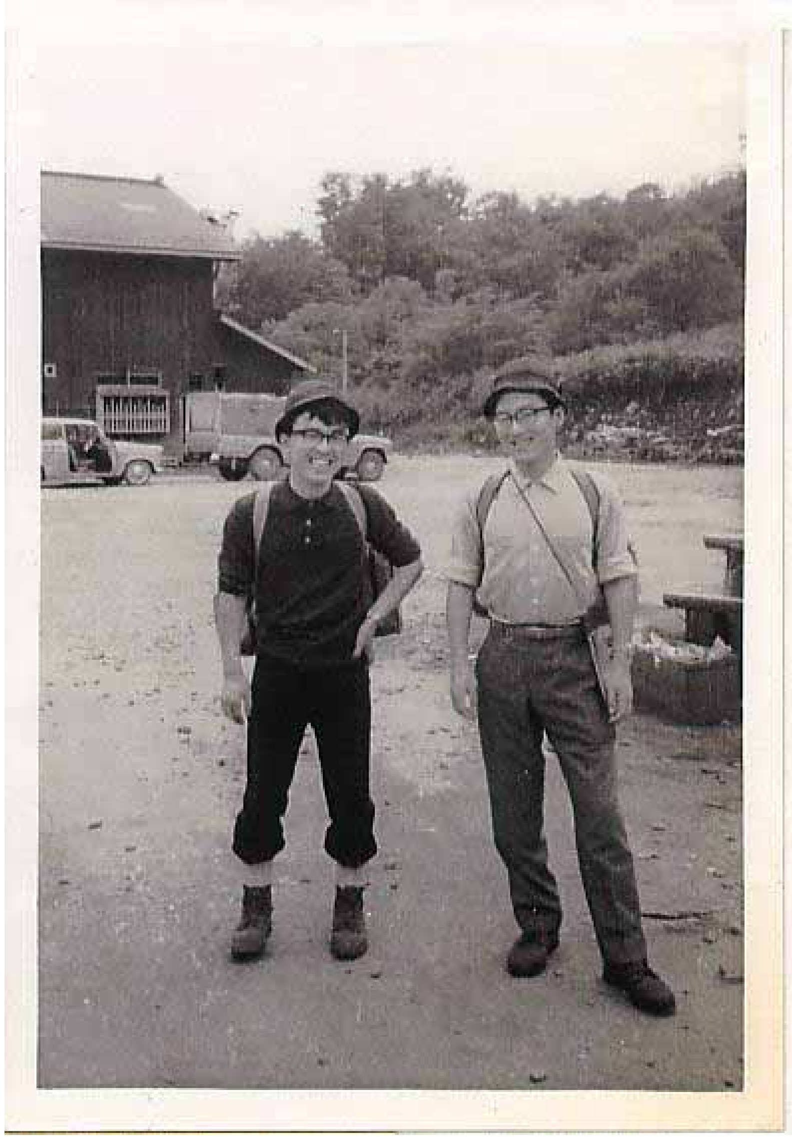

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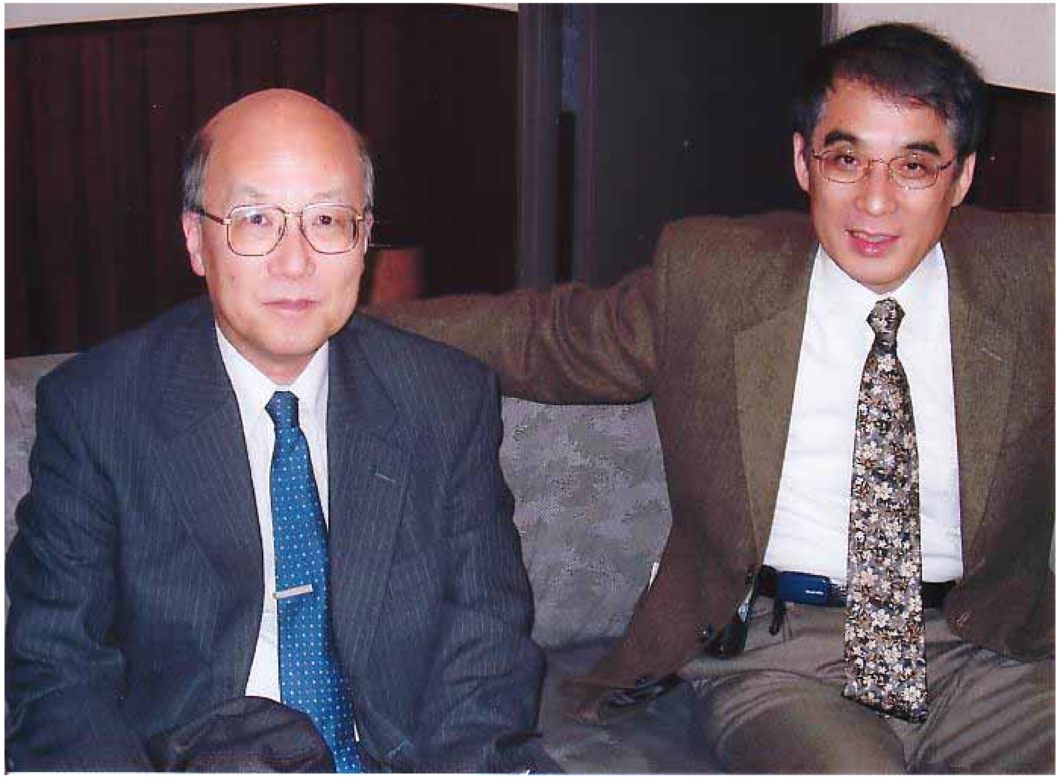

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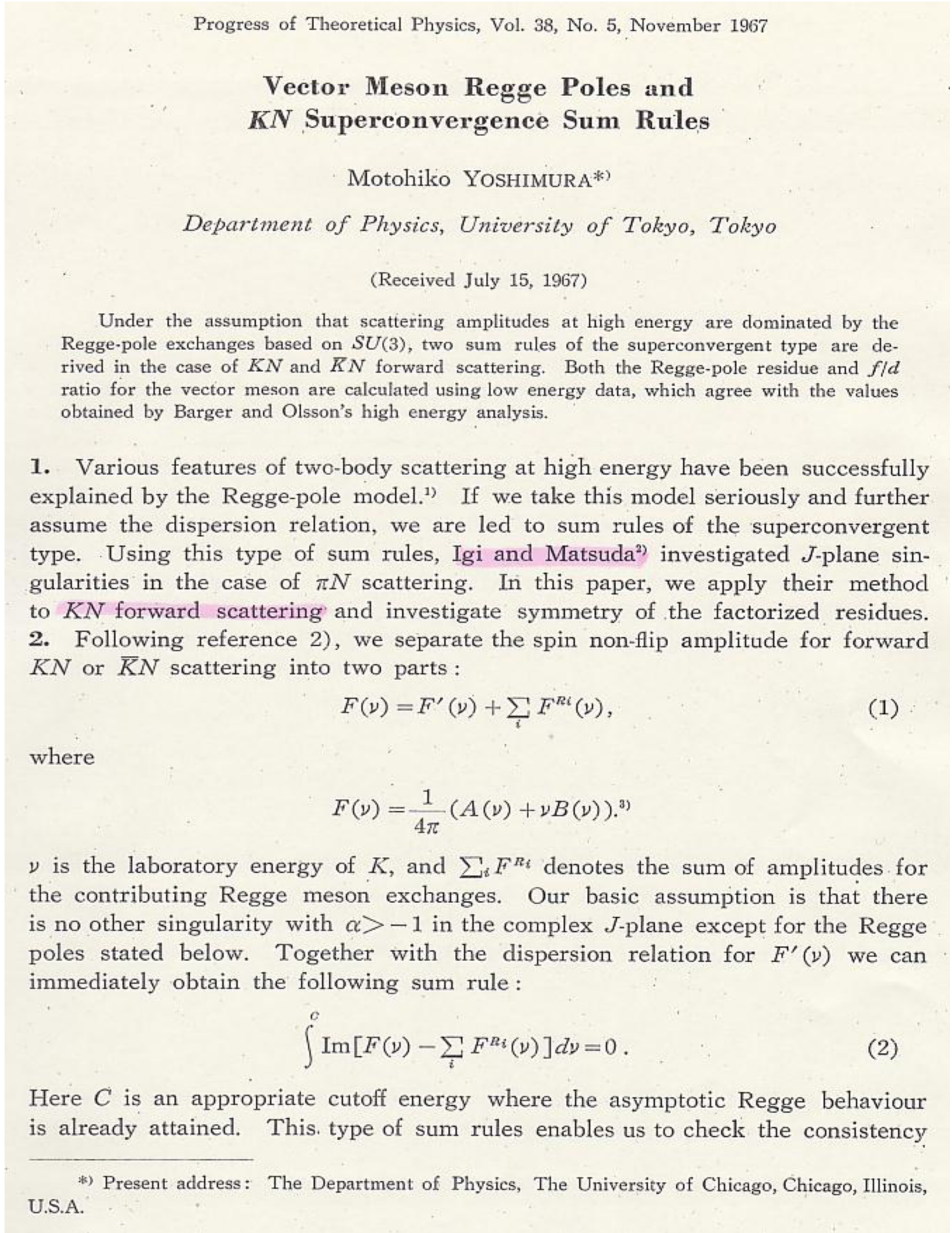

Progress of Theoretical Physics, Vol. 38, No. 5, November 1967

Vector Meson Regge Poles and
***KN* Superconvergence Sum Rules**

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(Received July 15, 1967)

Under the assumption that scattering amplitudes at high energy are dominated by the Regge-pole exchanges based on $SU(3)$, two sum rules of the superconvergent type are derived in the case of KN and $\bar{K}N$ forward scattering. Both the Regge-pole residue and f/d ratio for the vector meson are calculated using low energy data, which agree with the values obtained by Barger and Olsson's high energy analysis.

1. Various features of two-body scattering at high energy have been successfully explained by the Regge-pole model.[1] If we take this model seriously and further assume the dispersion relation, we are led to sum rules of the superconvergent type. Using this type of sum rules, Igi and Matsuda[2] investigated J-plane singularities in the case of πN scattering. In this paper, we apply their method to KN forward scattering and investigate symmetry of the factorized residues.
2. Following reference 2), we separate the spin non-flip amplitude for forward KN or $\bar{K}N$ scattering into two parts:

$$F(\nu)=F'(\nu)+\sum_i F^{R_i}(\nu), \qquad (1)$$

where

$$F(\nu)=\frac{1}{4\pi}(A(\nu)+\nu B(\nu)).^{3)}$$

ν is the laboratory energy of K, and $\sum_i F^{R_i}$ denotes the sum of amplitudes for the contributing Regge meson exchanges. Our basic assumption is that there is no other singularity with $\alpha>-1$ in the complex J-plane except for the Regge poles stated below. Together with the dispersion relation for $F'(\nu)$ we can immediately obtain the following sum rule:

$$\int^{C} \mathrm{Im}[F(\nu)-\sum_i F^{R_i}(\nu)]d\nu=0\,. \qquad (2)$$

Here C is an appropriate cutoff energy where the asymptotic Regge behaviour is already attained. This type of sum rules enables us to check the consistency

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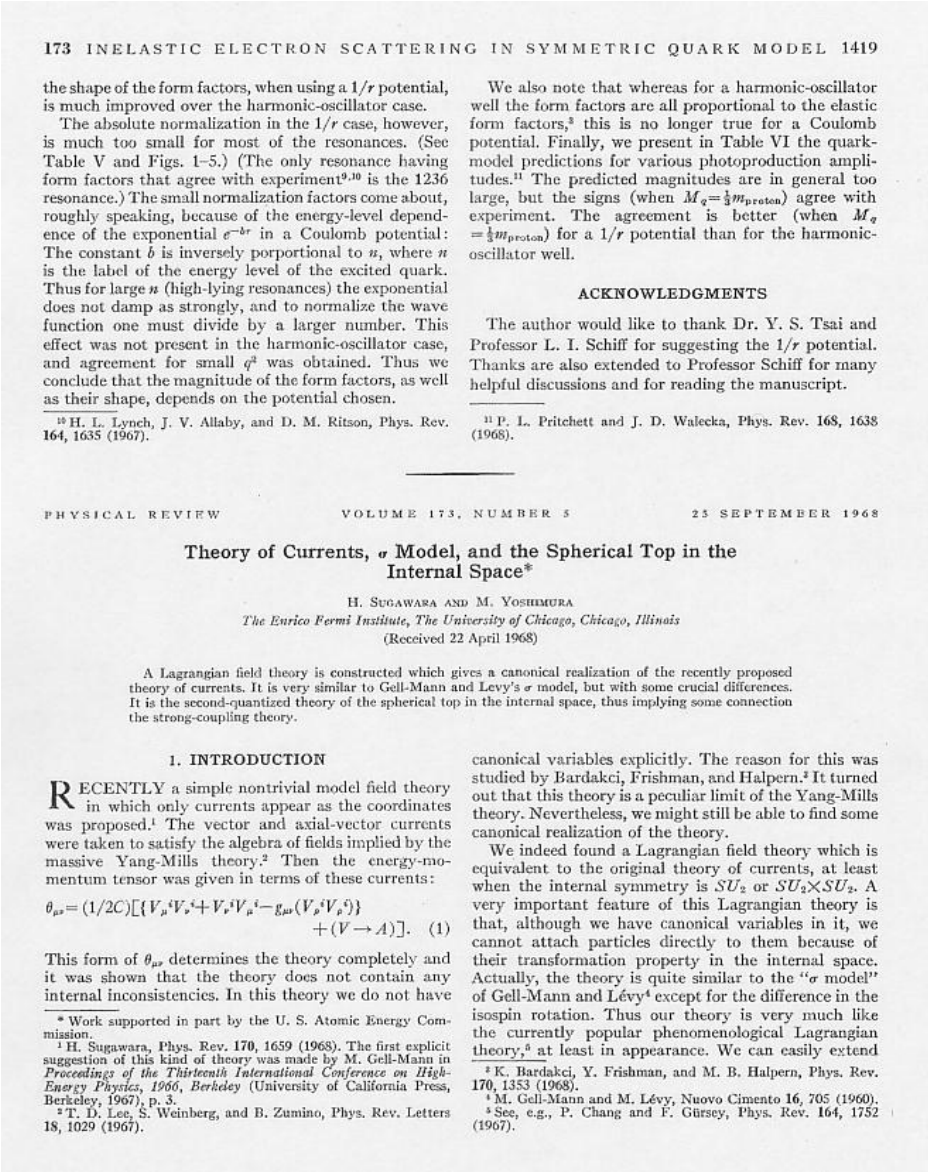
- Analytic S-matrix
- Current algebra and soft pion theorem, with background of quarks
- Duality of resonance and Regge behavior, leading to

Dual resonance model to String model of hadrons

菅原理論



173 INELASTIC ELECTRON SCATTERING IN SYMMETRIC QUARK MODEL 1419

the shape of the form factors, when using a $1/r$ potential, is much improved over the harmonic-oscillator case.

The absolute normalization in the $1/r$ case, however, is much too small for most of the resonances. (See Table V and Figs. 1–5.) (The only resonance having form factors that agree with experiment[9,10] is the 1236 resonance.) The small normalization factors come about, roughly speaking, because of the energy-level dependence of the exponential e^{-br} in a Coulomb potential: The constant b is inversely porportional to n, where n is the label of the energy level of the excited quark. Thus for large n (high-lying resonances) the exponential does not damp as strongly, and to normalize the wave function one must divide by a larger number. This effect was not present in the harmonic-oscillator case, and agreement for small q^2 was obtained. Thus we conclude that the magnitude of the form factors, as well as their shape, depends on the potential chosen.

We also note that whereas for a harmonic-oscillator well the form factors are all proportional to the elastic form factors,[3] this is no longer true for a Coulomb potential. Finally, we present in Table VI the quark-model predictions for various photoproduction amplitudes.[11] The predicted magnitudes are in general too large, but the signs (when $M_q=\frac{1}{3}m_{\text{proton}}$) agree with experiment. The agreement is better (when $M_q=\frac{1}{3}m_{\text{proton}}$) for a $1/r$ potential than for the harmonic-oscillator well.

ACKNOWLEDGMENTS

The author would like to thank Dr. Y. S. Tsai and Professor L. I. Schiff for suggesting the $1/r$ potential. Thanks are also extended to Professor Schiff for many helpful discussions and for reading the manuscript.

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PHYSICAL REVIEW VOLUME 173, NUMBER 5 25 SEPTEMBER 1968

Theory of Currents, σ Model, and the Spherical Top in the Internal Space*

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(Received 22 April 1968)

A Lagrangian field theory is constructed which gives a canonical realization of the recently proposed theory of currents. It is very similar to Gell-Mann and Levy's σ model, but with some crucial differences. It is the second-quantized theory of the spherical top in the internal space, thus implying some connection the strong-coupling theory.

1. INTRODUCTION

RECENTLY a simple nontrivial model field theory in which only currents appear as the coordinates was proposed.[1] The vector and axial-vector currents were taken to satisfy the algebra of fields implied by the massive Yang-Mills theory.[2] Then the energy-momentum tensor was given in terms of these currents:

$$\theta_{\mu\nu}=(1/2C)[\{V_\mu^iV_\nu^i+V_\nu^iV_\mu^i-g_{\mu\nu}(V_\rho^iV_\rho^i)\}+(V\to A)]. \quad (1)$$

This form of $\theta_{\mu\nu}$ determines the theory completely and it was shown that the theory does not contain any internal inconsistencies. In this theory we do not have canonical variables explicitly. The reason for this was studied by Bardakci, Frishman, and Halpern.[3] It turned out that this theory is a peculiar limit of the Yang-Mills theory. Nevertheless, we might still be able to find some canonical realization of the theory.

We indeed found a Lagrangian field theory which is equivalent to the original theory of currents, at least when the internal symmetry is SU_2 or $SU_2\times SU_2$. A very important feature of this Lagrangian theory is that, although we have canonical variables in it, we cannot attach particles directly to them because of their transformation property in the internal space. Actually, the theory is quite similar to the "σ model" of Gell-Mann and Lévy[4] except for the difference in the isospin rotation. Thus our theory is very much like the currently popular phenomenological Lagrangian theory,[5] at least in appearance. We can easily extend

* Work supported in part by the U. S. Atomic Energy Commission.

[1] H. Sugawara, Phys. Rev. **170**, 1659 (1968). The first explicit suggestion of this kind of theory was made by M. Gell-Mann in *Proceedings of the Thirteenth International Conference on High-Energy Physics, 1966, Berkeley* (University of California Press, Berkeley, 1967), p. 3.

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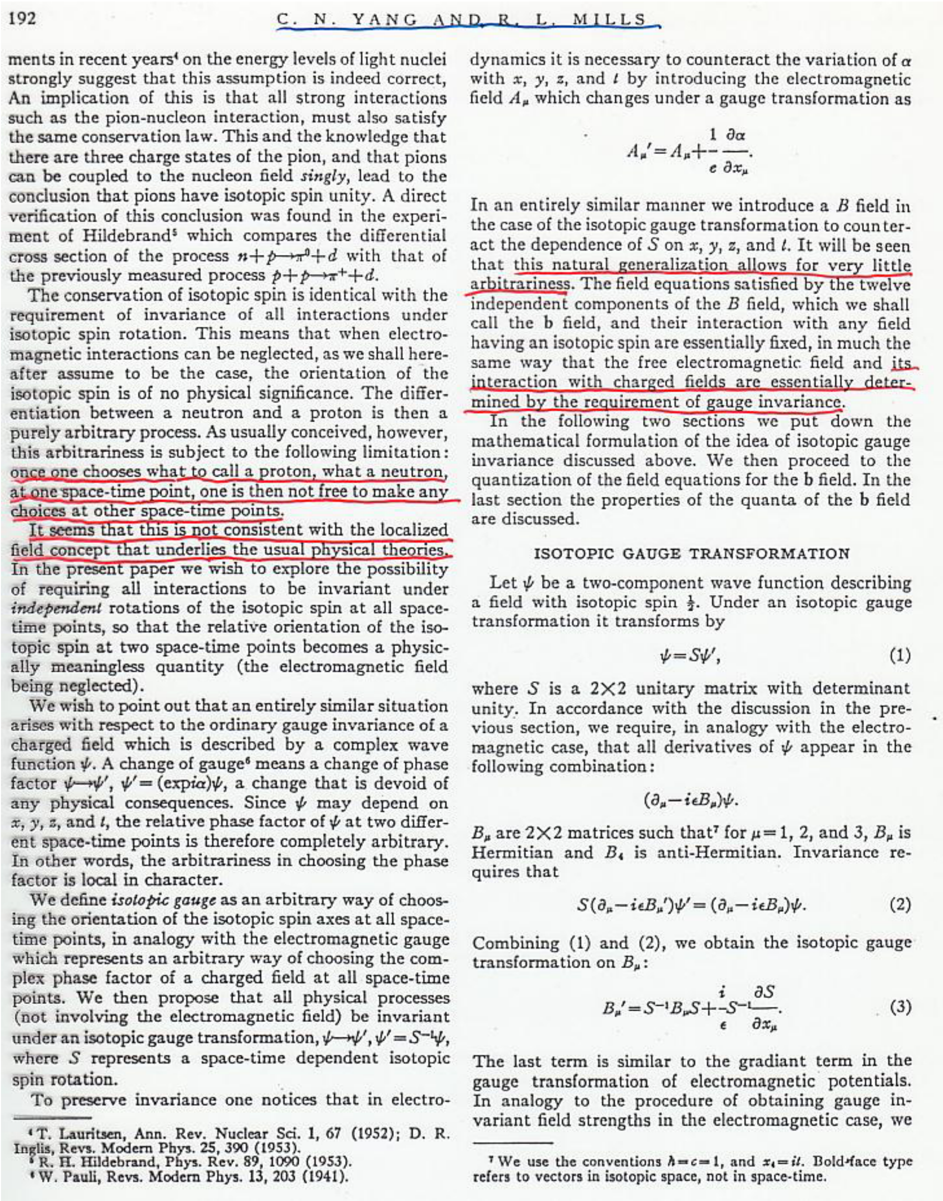
[5] See, e.g., P. Chang and F. Gürsey, Phys. Rev. **164**, 1752 (1967).

- 局所カレントを基本変数として、統一理論の構築をねらう。
- ストリング理論、数理物理へ多大の影響を与えた。

Influential paper 1　Yang-Mills



ments in recent years[4] on the energy levels of light nuclei strongly suggest that this assumption is indeed correct. An implication of this is that all strong interactions such as the pion-nucleon interaction, must also satisfy the same conservation law. This and the knowledge that there are three charge states of the pion, and that pions can be coupled to the nucleon field *singly*, lead to the conclusion that pions have isotopic spin unity. A direct verification of this conclusion was found in the experiment of Hildebrand[5] which compares the differential cross section of the process $n+p\to\pi^0+d$ with that of the previously measured process $p+p\to\pi^++d$.

The conservation of isotopic spin is identical with the requirement of invariance of all interactions under isotopic spin rotation. This means that when electromagnetic interactions can be neglected, as we shall hereafter assume to be the case, the orientation of the isotopic spin is of no physical significance. The differentiation between a neutron and a proton is then a purely arbitrary process. As usually conceived, however, this arbitrariness is subject to the following limitation: once one chooses what to call a proton, what a neutron, at one space-time point, one is then not free to make any choices at other space-time points.

It seems that this is not consistent with the localized field concept that underlies the usual physical theories. In the present paper we wish to explore the possibility of requiring all interactions to be invariant under *independent* rotations of the isotopic spin at all space-time points, so that the relative orientation of the isotopic spin at two space-time points becomes a physically meaningless quantity (the electromagnetic field being neglected).

We wish to point out that an entirely similar situation arises with respect to the ordinary gauge invariance of a charged field which is described by a complex wave function ψ. A change of gauge[6] means a change of phase factor $\psi\to\psi'$, $\psi'=(\exp i\alpha)\psi$, a change that is devoid of any physical consequences. Since ψ may depend on x, y, z, and t, the relative phase factor of ψ at two different space-time points is therefore completely arbitrary. In other words, the arbitrariness in choosing the phase factor is local in character.

We define *isotopic gauge* as an arbitrary way of choosing the orientation of the isotopic spin axes at all space-time points, in analogy with the electromagnetic gauge which represents an arbitrary way of choosing the complex phase factor of a charged field at all space-time points. We then propose that all physical processes (not involving the electromagnetic field) be invariant under an isotopic gauge transformation, $\psi\to\psi'$, $\psi'=S^{-1}\psi$, where S represents a space-time dependent isotopic spin rotation.

To preserve invariance one notices that in electrodynamics it is necessary to counteract the variation of α with x, y, z, and t by introducing the electromagnetic field A_μ which changes under a gauge transformation as

$$A_\mu'=A_\mu+\frac{1}{e}\frac{\partial\alpha}{\partial x_\mu}.$$

In an entirely similar manner we introduce a B field in the case of the isotopic gauge transformation to counteract the dependence of S on x, y, z, and t. It will be seen that this natural generalization allows for very little arbitrariness. The field equations satisfied by the twelve independent components of the B field, which we shall call the **b** field, and their interaction with any field having an isotopic spin are essentially fixed, in much the same way that the free electromagnetic field and its interaction with charged fields are essentially determined by the requirement of gauge invariance.

In the following two sections we put down the mathematical formulation of the idea of isotopic gauge invariance discussed above. We then proceed to the quantization of the field equations for the **b** field. In the last section the properties of the quanta of the **b** field are discussed.

ISOTOPIC GAUGE TRANSFORMATION

Let ψ be a two-component wave function describing a field with isotopic spin $\frac{1}{2}$. Under an isotopic gauge transformation it transforms by

$$\psi=S\psi', \tag{1}$$

where S is a 2×2 unitary matrix with determinant unity. In accordance with the discussion in the previous section, we require, in analogy with the electromagnetic case, that all derivatives of ψ appear in the following combination:

$$(\partial_\mu-i\epsilon B_\mu)\psi.$$

B_μ are 2×2 matrices such that[7] for $\mu=1, 2$, and 3, B_μ is Hermitian and B_4 is anti-Hermitian. Invariance requires that

$$S(\partial_\mu-i\epsilon B_\mu')\psi'=(\partial_\mu-i\epsilon B_\mu)\psi. \tag{2}$$

Combining (1) and (2), we obtain the isotopic gauge transformation on B_μ:

$$B_\mu'=S^{-1}B_\mu S+\frac{i}{\epsilon}S^{-1}\frac{\partial S}{\partial x_\mu}. \tag{3}$$

The last term is similar to the gradiant term in the gauge transformation of electromagnetic potentials. In analogy to the procedure of obtaining gauge invariant field strengths in the electromagnetic case, we

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- 当時はほとんど引用されていなかったが、知る人ぞ知る論文だったよう。関連して、桜井純先生の長大な論文にも影響された。

- 学部4年でセミナー指導を受けた、岡林先生の影響もあったかもしれない。

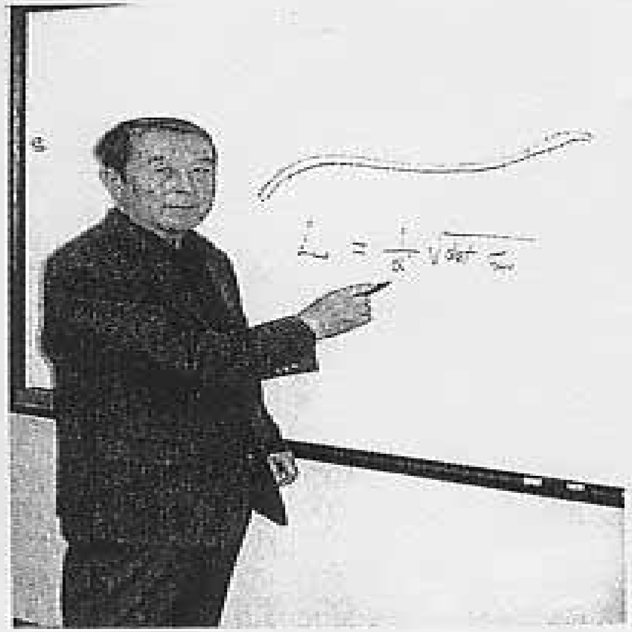
南部先生

南部陽一郎 Yoichiro Nambu

一九二一年福井市生まれ。東大卒業後草創期の大阪市立大に赴任。五二年に渡米、プリンストン高等研究所を経てシカゴ大学教授となる。ベーテ＝サルピーター＝南部方程式の導出、素粒子の超伝導体模型の発案、クォークの「カラー」の前ぶれとなった三重クォーク模型や、閉じ込め理論の先駆となったヒモ模型の提示等、国際物理学界のスプリンターとして、つぎつぎに最先端のテーマに挑戦、自在なアイディアを放出しつづけている。日本の後進の指導にも熱心で帰日もしばしば。七六年オッペンハイマー賞受賞、七八年に文化勲章受章。

なんぶ・よういちろう

シカゴ大学名誉教授。1921年、東京生まれ。東大物理学科卒。大阪市立大学教授などを経て、58年から90年までシカゴ大教授。米国科学賞、ウルフ賞（イスラエル）などを受賞。文化勲章受章。



With Nambu

Here s is the center-of-mass energy squared in $(\mathrm{GeV}/c)^2$; that is, the Regge parameter conventionally called s_0 is taken to be 1.0 $(\mathrm{GeV}/c)^2$. $f(t)$ is assumed to be essentially constant for $|t| < 1.0\ (\mathrm{GeV}/c)^2$, and $\alpha(t)$ is taken to be a linear function of t. This expression serves as a useful parametrization to study shrinkage. For the momentum range 5.4 to 29.4 GeV/c and $|t| < 1.0\ (\mathrm{GeV}/c)^2$ the results are, for np,

$$\alpha(t) = (1.08 \pm 0.06) - (0.86 \pm 0.18)|t|;$$

for pp,

$$\alpha(t) = (1.05 \pm 0.02) - (0.69 \pm 0.05)|t|$$

(Ref. 4); for $\bar{p}p$,

$$\alpha(t) = (0.90 \pm 0.08) + (0.91 \pm 0.38)|t|$$

(Ref. 4). On the basis of this parametrization the np shrinkage appears to be the same as the pp shrinkage within the errors.

It is our pleasure to acknowledge the assistance of the AGS staff, particularly J. Sanford and W. Merkle (deceased), for their assistance in setting up and running the experiment. We would also like to thank G. DeMeester, O. Haas, S. T. Powell, III, R. Seefred, J. Smith, and S. Wilson for their help with the setup.

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AXIAL-VECTOR FORM FACTOR OF NUCLEON DETERMINED FROM THRESHOLD ELECTROPION PRODUCTION*

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(Received 10 November 1969)

Currently available electropion-production data near threshold have been analyzed according to the soft-pion theorem which expresses the threshold cross section in terms of the vector and axial-vector form factors of the nucleon. We determine the axial-vector form factor for q^2 up to ~7 $(\mathrm{GeV}/c)^2$.

As is well known, the condition of partial conservation of axial-vector current (PCAC) and current algebra predict that the amplitudes for electropion production at threshold,[1-5] that is, for

$$e + p \to e + p + \pi^0,$$
$$e + p \to e + n + \pi^+,$$

are expressible in terms of the vector and axial-vector form factors of the nucleon in the ideal soft-pion limit. This enables one to determine the unknown (isovector) axial form factor $G_A(q^2)$ which could otherwise be obtained only from high-energy experiments. Difficulties arise because one has to work with small cross sections along the limits of phase space, taking account of the radiative correction and the deviation from the soft-pion limit due to the finite pion mass. Earlier attempts[5] to determine $G_A(q^2)$ from electroproduction data have been restricted to small

$q^2 \lesssim 0.6$ (GeV/c) with results which are not inconsistent with the crude information available from neutrino experiments.

In this report we use all the currently available data from Stanford Linear Accelerator Center and Massachusetts Institute of Technology (SLAC-MIT),[6,7] DESY,[8] and other groups,[9,10] which cover q^2 ranging up to 7.3 $(\mathrm{GeV}/c)^2$ and the incident-electron energy up to 17.7 GeV. The basic formula can be cast in the form[11]

$$\frac{1}{k^*}\frac{d^2\sigma}{dE_e'd\Omega} = \left(\frac{d\sigma}{d\Omega}\right)_{\mathrm{Mott}} \frac{g^2}{16\pi^2(M_N+m_\pi)^2} \times \left[W_2 + 2\tan^2\frac{\theta}{2}W_1\right] \quad (1)$$

where k^* is the pion momentum in the final πN center-of-mass frame. In the soft-pion limit the two structure functions for the process $n+\pi^+$ are given by

$$W_2 = 2\left[\left(G_A + \frac{q^2}{2M_N^2+q^2}G_M^n\right)^2 + 4M_N^2q^2\left(\frac{G_E^n}{2M_N^2+q^2}\right)^2\right],$$
$$W_1 = 2\left(1+\frac{q^2}{4M_N^2}\right)\left(G_A + \frac{q^2}{2M_N^2+q^2}G_M^n\right)^2. \quad (2)$$

For the process $p+\pi^0$, omit G_A, replace the magnetic and electric vector form factors of the neutron (G_M^n and G_E^n) with those of the proton (G_M^p and G_E^p), and remove the factor 2. The measured threshold cross section is the sum of the two. The isovector axial form factor $G_A(q^2)$ is defined relative to the Kroll-Ruderman value at $q^2 = 0$, with the π-N coupling constant $g^2/4\pi = 14.6$. For the vector form factors the standard dipole parametrization

$$G_E^p = G_M^p/\mu_p = G_M^n/\mu_n = [1+q^2/0.71\ (\mathrm{GeV}/c)^2]^{-2},$$
$$G_E^n = 0,$$

is sufficient for our purpose.

In analyzing the data, the following procedures were taken for handling the necessary corrections.

(1) Radiative correction. Fortunately, the bulk of the radiative correction at the inelastic threshold is caused by the radiative tail from the elastic scattering, which is easy to analyze and isolate. In fact, the results of Ref. 7 indicate the consistency of such a correction procedure, and we simply relied on the corrected data given in the original sources.

(2) Correction for finite pion mass. We adopted a covariant pole-dominance model[12] in which only nucleon- and pion-pole diagrams with vector form factors, and certain axial-vector contact

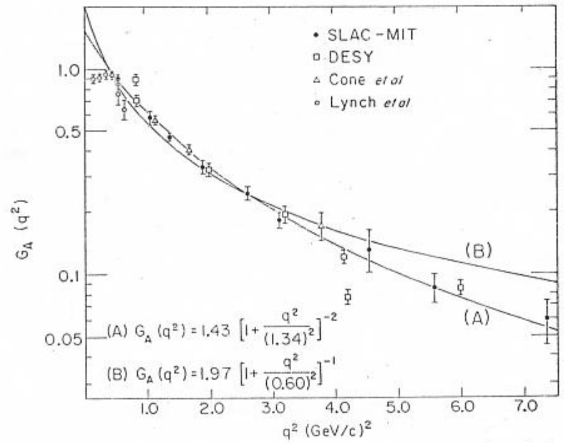

FIG. 1. Axial-vector form factor calculated by using the modified Nambu-Shrauner formula (see text). Best fits were obtained by fitting only those data with higher q^2, i.e., SLAC-MIT, DESY, and Cone et al., disregarding the normalization condition $G_A(0) = 1$.

String model

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Dual-Resonance Model with Quark Spin*†

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(Received 21 September 1970)

We present a dual-resonance model with nontrivial quark spin factors. An amplitude is found which satisfies factorization and eliminates the parity-doubling ghosts. An application to $\pi\pi$ elastic scattering indicates that the positivity condition is not met on the first daughter trajectory if one assumes realistic values for the mass and intercept.

I. INTRODUCTION

RECENT developments of the dual-resonance model[1,2] have revealed a close connection of the duality concept with the quark model.[3–5] It is now possible to embark on the construction of a hadron model out of quarks in the manner represented by the Harari-Rosner quark diagram.[6,7]

A crucial step in this program is a proof of the factorization property of the dual-resonance model. The proof of factorization has been extended to all the daughter trajectories.[8,9] The resonance spectrum in the model has been greatly clarified using the simple device of the harmonic oscillator.[10–12] Roughly speaking, mesons appear to be bound states of the quark and antiquark with a relativistic string between them.

Previous attempts[4,5] to incorporate quark spin into the dual model have suffered from a serious defect. Consistent factorization of the spin factor considered in that approach demands the existence of ghosts associated with negative-parity quarks. On the other hand, a recent work of Carlitz and Kislinger[13] provided a new way to avoid parity doubling of the fermion trajectory in the Van Hove model. Motivated by this work, many people have proposed to dualize the projection operator to eliminate parity-doubling ghosts.[14–16] We will present in this paper a different, but closely related approach to a correct treatment of the quark spin.

Our guiding principle in selecting a spin factor is the simple over-all picture of the dual-resonance model of Refs. 10–12. After constructing a cyclically symmetric amplitude of mesons, we proceed to check factorization of the whole amplitude. A simple quark propagator considered in Sec. II turns out to eliminate the parity-doubling ghosts from the leading trajectory only. A generalization of the propagator is then suggested, and elimination of ghosts from all the trajectories, as well as complete factorization, is proved in Sec. III. The generalized amplitude resembles a recent model of Carlitz, Ellis, Freund, and Matsuda.[16] The main difference lies in our insistence on the original form of the quark projector; therefore, we factorize the meson

* Supported in part by the U. S. Atomic Energy Commission.
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OPERATORIAL FACTORIZATION AND SYMMETRY OF THE SHAPIRO-VIRASORO MODEL*

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Received 28 November 1970

We present an operator formalism which exhibits factorization of the nonplanar dual model proposed first by Virasoro and generalized by Shapiro. The model is shown to contain sufficient gauge operators as well as a selection rule in order to eliminate all the ghosts. A further generalization of the model is also suggested.

Extensive efforts have been devoted to the generalized Veneziano model [1] in order to construct a Feynman-like theory based on duality. It may well be asked, however, whether there exists another class of dual models which can also cope with nonplanar duality in a single representation. A four-point amplitude with the nonplanar duality was proposed some time ago by Virasoro [2], and extended very recently to n-point amplitude by Shapiro [3] in a particular case of the intercept ($\alpha_0 = 2$). The purpose of this paper is to provide an operator formalism which manifests factorization property of the Shapiro's generalization and facilitates more ambitious program such as unitarization. We will also show that the model contains sufficient gauge operators to eliminate all the ghosts through the same mechanism as that considered by Virasoro [4]. No attempt is made to relax the condition of the intercept.

We begin by writing the n-point amplitude of ref. [3] with the intercept $\alpha_0 = 2$:

$$A_n = \left|(z_a-z_b)(z_b-z_c)(z_c-z_a)\right|^2 \times \int \prod_{i \neq a,b,c} d^2 z_i \prod_{i<j} |z_i - z_j|^{2q_i q_j} . \tag{1}$$

The integration region is extended to the whole complex plane without any restriction among the variables, z_i. Poles show up in any two-body channel (ij) since z_i can approach z_j, causing divergence of the integrand. A basic property of the integral representation (1) is that the integrand is invariant under the projective transformation of the variables z_i if the intercept $\alpha_0 = q^2 = 2$. This assures that one can choose any points for a, b and c and any values for z_a, z_b and z_c. In the following discussion we adopt a choice, $z_1 = 0$, $z_n = \infty$, $z_c = 1$ for definiteness. The eq. (1) then reads

$$A_n = \int \prod_{i \neq 1,n,c} d^2 z_i |z_i|^{2q_i q_1} \prod_{2 \leq i < j \leq n-1} |z_i - z_j|^{2q_i q_j} \tag{2}$$

It has been demonstrated in ref. [3] that this leads to the Virasoro amplitude with $\alpha_0 = 2$ for the four point function.

Let us break up the amplitude A_n into $(n-2)!$ terms so that each term corresponds to a definite ordering of z_i in the sense of modulus, for example $\infty > |z_{n-1}| \geq |z_{n-2}| \geq \ldots \geq |z_2|$. We express this as follows:

$$A_n = \sum_{\mathrm{P}} F_n(q_1 \ldots q_n) . \tag{3}$$

The summation is taken over all permutations of the momenta, $q_2 \ldots q_{n-1}$. As will be shown, it turns out that this breaking corresponds to the decomposition of the Virasoro amplitude into the sum of the s- and u- channel pole terms for $n = 4$. For the first term in eq. (3) with $|z_i| \geq |z_{i-1}|$ for any i it is convenient to introduce the polar coordinates

$$z_i = r_i r_{i+1} \cdots r_{c-1} \exp(i\theta_i) \text{ for } 2 \leq i \leq c-1 ;$$

$$z_i^{-1} = r_{c+1} r_{c+2} \cdots r_i \exp(-i\theta_i) \text{ for } c+1 \leq i \leq n-1 .$$

* Research supported by the Air Force Office of Scientific Research, Office of Aerospace Research, US Air Force, under Contract number F44620-70-C-0028.

私の研究遍歴と時代背景　０

- 大学院時代（‘６５－’７０、東大、シカゴ大)

  Dual resonance model (ハドロンの弦模型、のちに万物の超弦模型に発展)

  クオーク模型、Deep inelastic scattering

- ポスドク時代 (‘７０－’７５、バークレイ、ペン、パリ)

  電弱統一理論

隠れ家

- Kabir の弱い相互作用の論文選集
- １９６０年代前半までの重要論文を網羅

- Lee-Yang のパリティ非保存
- (V-A) x (V-A) で完璧な現象論
- 場の理論としてはナンセンス

素晴らしいドラマを予見

Influential paper 2

Reprinted from the Physical Review Letters 19 (1967) 1264-1266

A MODEL OF LEPTONS*

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Massachusetts Institute of Technology, Cambridge, Massachusetts
(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

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

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

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^-

1264

- t' Hooft のプレプリントの出現(1971年秋)直後、知る

- 同時期のFaddeev-Popov論文が技術的に決定的に役にたった。

WS model as a field theory

Reprinted from:
PHYSICAL REVIEW D VOLUME 6, NUMBER 1 1 JULY 1972

Muon Magnetic Moment in a Finite Theory of Weak and Electromagnetic Interactions*

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(Received 22 February 1972)

We calculate the weak-interaction contribution to the muon anomalous magnetic moment of the order $G_W m_\mu^2$ in the Weinberg model for leptons. Using a ξ-limiting procedure we obtain a finite correction to $\frac{1}{2}(g-2)$, which has the value $(1.8-2.2)\times 10^{-9}$.

The unification of weak and electromagnetic interactions has been a fascinating subject of theoretical interest in particle physics. In particular, Weinberg[1] proposed an ingenious Lagrangian model of leptons in which the weak interactions are mediated by massive gauge bosons. In this scheme the leptons (electrons and muons) and the weak intermediate bosons, which are initially massless, acquire masses due to the spontaneous breakdown of a lepton symmetry. The well-known divergence difficulty of theories with massive weak bosons does not occur, because the smooth convergent behavior of the massless gauge theory initially started with is maintained in a nontrivial way even after the introduction of the symmetry breaking.[2,3] In fact, Weinberg[4] calculated a few loop diagrams using propagators for massive vector bosons and confirmed that bad divergences cancel each other among different diagrams.

The purpose of this note is to present a finite additional correction to the muon magnetic moment as predicted by Weinberg's model. We obtain a small but finite correction of the order $G_W m_\mu^2$ which depends on a single unknown parameter. Within the estimated range of this parameter, the finite correction has the values $(1.8-2.2)\times 10^{-9}$ which is to be compared with previous works.[5,6] We have used a ξ-limiting procedure[7] to extract the finite value.

In the model of Ref. 1 we have a neutral boson Z as well as charged bosons $W^\pm$. Their couplings to leptons and their masses are given as[8]

$$\frac{g}{2\sqrt{2}}\,\bar{\mu}\gamma_\alpha(1+\gamma_5)\nu_\mu W^\alpha, \quad \bar{\mu}(g_1\gamma_\alpha+g_2\gamma_\alpha\gamma_5)\mu Z^\alpha,$$

$$M_W=\tfrac{1}{2}\lambda g, \quad M_Z=\tfrac{1}{2}\lambda(g^2+g'^2)^{1/2}, \tag{1}$$

$$g_1=\tfrac{1}{4}(g^2+g'^2)^{1/2}\,\frac{3g'^2-g^2}{g'^2+g^2}, \quad g_2=-\tfrac{1}{4}(g^2+g'^2)^{1/2}.$$

The relation to the electromagnetic charge e and the Fermi constant G_W is

$$e=\frac{gg'}{(g^2+g'^2)^{1/2}}, \quad \frac{G_W}{\sqrt{2}}=\frac{1}{2\lambda^2}. \tag{2}$$

To lowest order in G_W we will evaluate the two Feynman amplitudes depicted in Figs. 1 and 2. All other diagrams of the same order do not contribute to the anomalous magnetic moment. We also ignore the diagram with the scalar field ϕ replacing Z in Fig. 2, because of its small coupling and presumably large mass. The Feynman rules are easily derived from Eq. (14) of Ref. 1, and we use the same manifestly unitary gauge of Ref. 4. The only new rule added here is that we adopt the ξ-limiting procedure for both W and Z. This is needed to avoid ambiguous *finite* contributions that may come from making changes of variables in linearly and quadratically divergent integrals.[9] However, even in the absence of any regularization there are no divergences in the contribution to the anomalous magnetic moment, although some of the integrals are apparently logarithmically divergent.

We write the vertex amplitudes explicitly:

私の研究遍歴　１

- **東北大１（‘７５－’７９）**

  GUT

  Baryogenesis

後に話す

私の研究遍歴　2

- ＫＥＫ(‘７９ー’８８)

  Cosmology and Particle Physics

  Axion (制約、冷たい暗黒物質)

  Neutrino (GUT models, Baryogenesis との関連, oscillation, 超新星爆発),

  (KK) Extra dimensions (cyclic universe)

- 東北大２　と　宇宙線研(’８８ ー‘０４)

  宇宙の熱史の起源: エントロピー生成

  素過程への宇宙環境効果

Constraints from astroparticle physics

PHYSICAL REVIEW D VOLUME 26, NUMBER 8 15 OCTOBER 1982

Astrophysical constraints on a new light axion and other weakly interacting particles

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(Received 24 May 1982)

Constraints on the light axion of Dine, Fischler, and Srednicki are critically reexamined and upper bounds on its mass are derived from stars at various stages of evolution. A conservative upper bound for the axion mass is about 1 eV, while a model-dependent argument gives a better upper bound of mass ~ 0.07 eV. The same argument also applies to coupling of any massless pseudoscalar particle to electrons giving an upper bound of $|g_e| < 1\times 10^{-11}$.

I. INTRODUCTION

The problem of strong *CP* violation has been a long-standing puzzle since the discovery of instanton effects.[1] In 1977 Peccei and Quinn[2] suggested a solution to this problem by observing that the θ parameter responsible for the *CP* violation is made physically meaningless if an extra chiral symmetry is imposed. In their model this new chiral symmetry is spontaneously broken together with the electroweak $SU(2)\times U(1)$, and because of a chiral anomaly a pseudo-Nambu-Goldstone boson emerges.[3] This spinless boson, called the standard axion, has a reasonably well-determined mass and lifetime, which, however, seems excluded by almost all experiments.

More recently, Dine, Fischler, and Srednicki[4] (DFS), based on an earlier idea of Kim,[5] extended the Peccei-Quinn type of symmetry by allowing an $SU(2)\times U(1)$-singlet Higgs boson to break the new chiral symmetry. The axion associated with this chiral symmetry has a mass and a coupling to matter, both suppressed by a large vacuum expectation value (VEV) of the Higgs singlet which is not constrained by the electroweak unification scale of order 100 GeV. It was conjectured[4] that the new scale of chiral breaking should be above 10^9 GeV by using astrophysical constraints[6] on the Kim axion. If the scale is so high, we have no effective laboratory experiment to observe a trace of this new axion.

The purpose of the present paper is to critically reexamine the astrophysical mass bound of the DFS axion and update the old analysis made for Kim's axion by Dicus *et al.*[6] Since the present work overlaps and extends previous works, we shall explain why we embarked on this project. Historically, Sato and Sato[7] were the first who derived a mass bound for a light spinless boson coupling weakly to matter. The most stringent bound that they obtained comes from the cooling of red giants. They only considered as a process of energy loss a Compton-type process, $\gamma+e\rightarrow A+e$, where A couples to electrons via a scalar interaction. However, in stars like red giants the pseudoscalar coupling of the DFS axion is crucial because the energy-loss rate is suppressed by a factor of $(T/m_e)^2$ compared to the scalar coupling. Dicus *et al.* and Vysotsskii *et al.* extended[8,6] the analysis of Sato and Sato to the case of a pseudoscalar axion. Unfortunately, for Kim's axion relevant to our investigation here they[6] mainly discussed another important process, namely the Primakoff process, and derived an upper mass bound based on this process alone. Moreover, most arguments, which have been given so far to derive the bound, depend on stellar models in that the axion energy loss was calculated by using a temperature and density of stars such as those computed without axions. We found that there is no reliable mass bound for the DFS axion immediately available. Indeed, as we shall see in this paper, the most stringent mass bound for the DFS axion is derived from the Compton-type process in red giants, which was not considered in the paper of Dicus *et al.*[6] Even when the Primakoff process dominates as in the case of main-sequence stars, plasma effects turn out to be important, which was ignored by previous authors.

The essential observation in the discussion of astrophysical constraints is that stable stars tend to emit particles with energies less than the thermal energy in the stellar bath. If these particles couple weakly to matter, they easily escape a star and remove its energy too rapidly. The emission of these particles may be suppressed either because it is en-

VOLUME 59, NUMBER 16 PHYSICAL REVIEW LETTERS 19 OCTOBER 1987

Neutrino Burst from SN1987A and the Solar-Neutrino Puzzle

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[(2)]*Research Institute for Fundamental Physics, Kyoto University, Kyoto 606, Japan*
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(Received 8 April 1987; revised manuscript received 14 August 1987)

The prompt ν_e signal from the supernova explosion in the Large Magellanic Cloud presumably detected by Kamiokande II does not necessarily mean that the Mikheyev-Smirnov-Wolfenstein effect on the solar-neutrino flux is not operative. The electron neutrino, once rotated to a different-flavor neutrino in the progenitor star, can come back via the matter oscillation effect in the Earth, or a residual ν_e flux from the progenitor can directly hit the detector, saving the Mikheyev-Smirnov-Wolfenstein explanation of the solar-neutrino problem for a range of mixing parameters.

PACS numbers: 97.60.Bw, 12.15.Ff, 14.60.Gh, 96.60.Kx

The neutrino burst from SN1987A in the Large Magellanic Cloud, first discovered by Hirata *et al.*[1] (the Kamiokande II collaboration) and later confirmed by Bionta *et al.*[2] [the IMB (Irvine-Michigan-Brookhaven) group], gives a unique opportunity to explore the physics of supernova explosions. It is remarkable that gross features of these neutrino events, event rate, average neutrino energy, and time span, agree with theoretical calculations based on conventional models of the stellar collapse.[3] It is the first time in the history of modern science that dynamics of the stellar collapse, on the time scale of less than 10 sec, has been probed, with a positive result.

A closer examination of these events, however, reveals some unusual features that seem difficult to reconcile with the standard calculation. These may have interesting astrophysical and particle-physics implications.[4-6] In this paper we shall pay particular attention to the first forward events of Kamiokande II suggestive of the prompt neutronization burst and examine what they mean in the context of basic properties of the neutrino and how they are related to the solar-neutrino problem.

Recall that the basic process of detection in the water Cherenkov facility is $\nu_e+e\rightarrow\nu_e+e$ for the electron-type neutrinos and $\bar\nu_e+p\rightarrow e^++n$ for electron-type antineutrinos. The similar processes induced by ν_H and $\bar\nu_H$ ($H=\mu$ or τ) are unlikely to occur, since they have smaller cross sections. The former reaction (ν_e+e) is characterized by the directionality of the recoil electron in a forward cone of about 15°, while the latter ($\bar\nu_e+p$) yields an isotropic distribution of e^+ for neutrino energy of $\simeq 10$ MeV. It is thus natural to associate the first one or two forward (within $18°\pm18°$ and $15°\pm27°$ cone) events of the Kamiokande II observation with the prompt ν_e burst. The probability of finding two forward events within 42° out of randomly distributed $\bar\nu_e$ events is small, $\simeq 0.6\%$. The standard calculation[3] also supports this interpretation: Other types of neutrinos are not much emitted at the first instant. A potentially serious problem[5] that may be raised with this interpretation is that in the calculation of Wilson and co-workers, the yield of prompt ν_e events is much less ($\simeq 0.3$ event in Kamiokande II) and the observed duration of $\simeq 100$ ms between the first two events is too large. These two features are, however, nicely explained in the advective overturn model of Arnett.[7] This uncertainty in astrophysical models casts a doubt on interpreting the second event as the ν_e signal. The ambiguity is hoped to be resolved by future observations, but for the following analysis we shall assume that the first one or two events were caused by $\nu_e e$ scattering, mentioning parameter ranges in two cases. As pointed out in Ref. 4 and also by Walker and Schramm[8] prior to the supernova event, the prompt ν_e signal appears then to rule out the Mikheyev-Smirnov-Wolfenstein mechanism[9] of neutrino oscillation as a possible explanation of the solar-neutrino deficit,[10] because the ν_e burst generated at the core is converted to another type of neutrino (ν_μ or ν_τ) in passing through the outer region of the progenitor star, which is not dissimilar to the sun in its density.

This conclusion rests on the assumption that the conversion is very efficient in the progenitor star and that nothing drastic happens until the converted neutrino arrives at the detector. We have examined carefully whether this is true and, surprisingly, found that there are two possibilities to save the Mikheyev-Smirnov-Wolfenstein explanation: a possibility of the prompt ν_H being converted back to ν_e within the Earth, and the possibility of a sizable ν_e residual in the progenitor. These two cases can occur in different parameter regions of δm^2 and $\sin 2\theta$ that can then be tested in forthcoming experiments. These parameter regions differ somewhat, depending on whether one accepts the second event as due to $\nu_e e$ scattering.

The effects of neutrino oscillation in the Earth have been discussed in the literature.[11] As an idealization,

ミクロ物理学の進展 (最近40年間)

- 記述言語としての、場の量子論と
  ダイナミックスとしての、ゲージ場の確立
- 統一理論候補として、スーパーストリング
- フレーバー物理の進歩：標準模型の確立とその彼方
  ６　クオーク、　タウレプトン
  ＣＰの破れ現象の拡大：　Ｋ　－＞　Ｂ
  ニュートリノの質量と混合

ミクロ物理学の進展　２

- 宇宙物理の包含

  B-genesis, L-genesis

  インフレーション

  ダークマターとダークエネルギー

ミクロ法則が

熱い宇宙の始まりと終わりを支配

ビッグバン宇宙

- ハッブルの膨張則：　$v = \frac{r}{150\text{億年}}$
- 3 °K 黒体輻射
- 軽元素合成

最近の話題

- 物質・反物質不均衡 (バリオン非対称)
- インフレーション
- 暗黒物質
- 大規模構造
- 3 °K 非等方性

素粒子の大統一理論との関連が注目

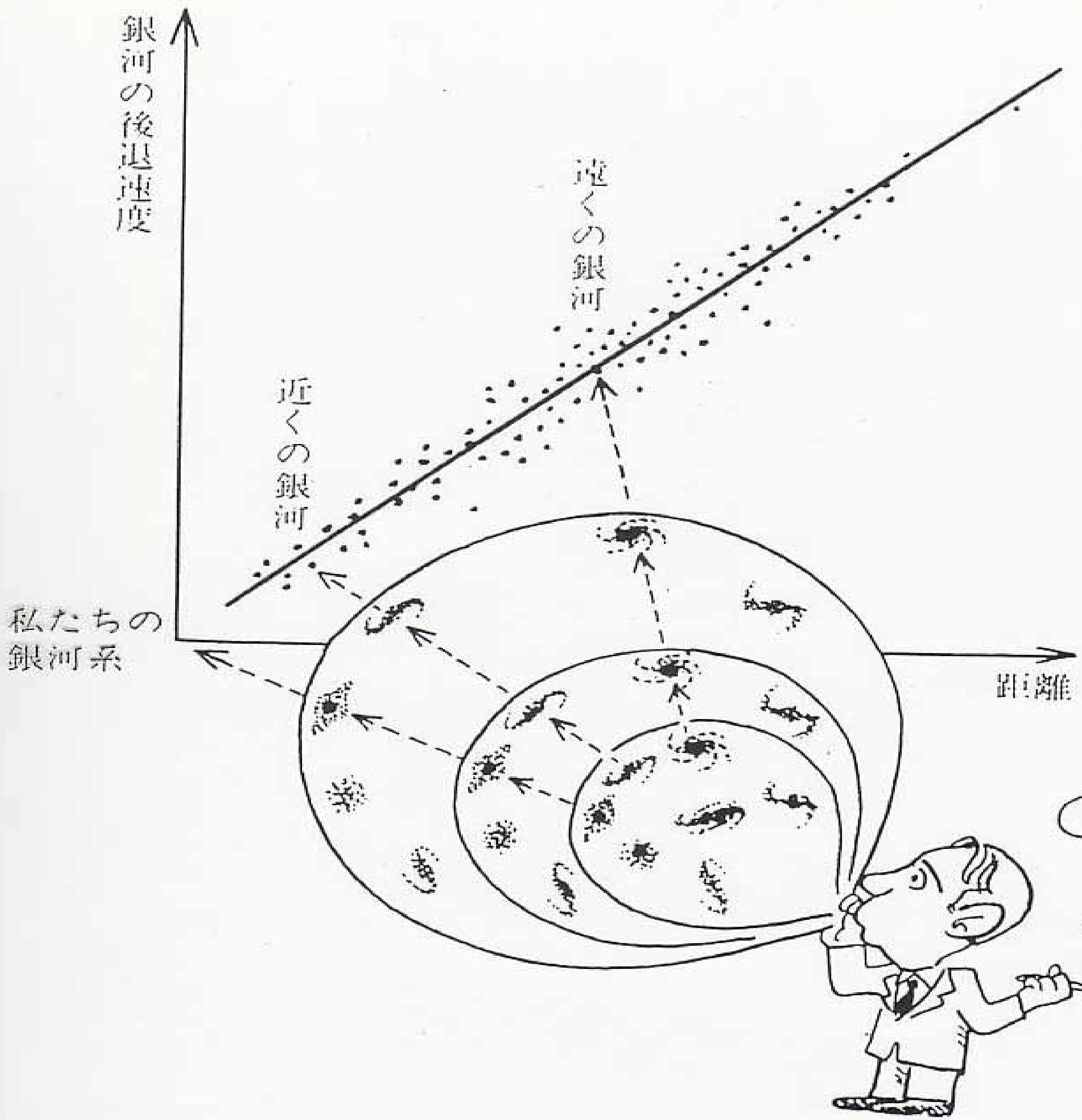
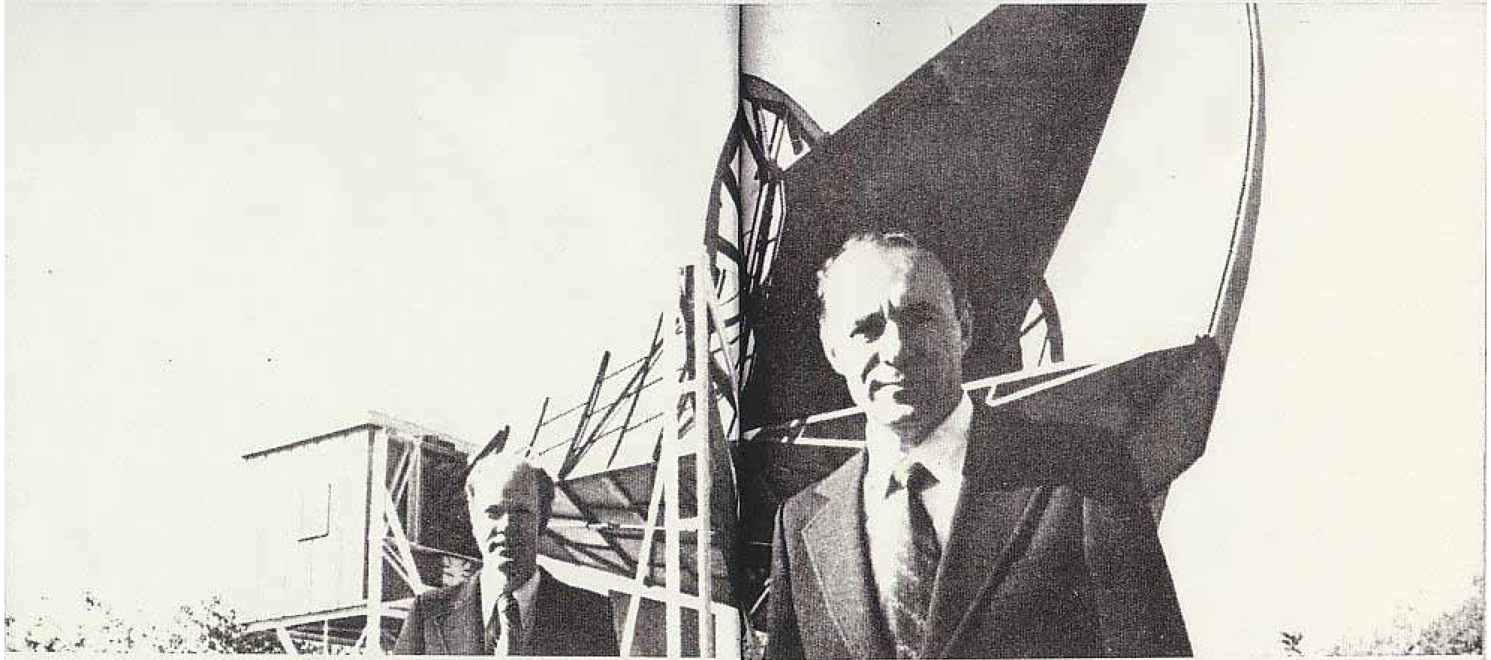
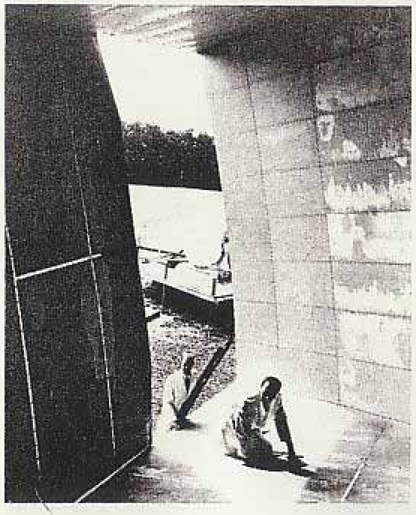


図 4.6　ハッブルの法則．遠くの銀河ほど速い速度で私たちから遠ざかっている．

Penzias and Wilson





Above: THE HOLMDEL RADIO TELESCOPE: Arno Penzias (right) and Robert W. Wilson (left) are shown here with the 20-foot horn antenna used by them in 1964-65 in their discovery of the 3° K cosmic microwave radiation background. This telescope is at the Holmdel, New Jersey, site of the Bell Telephone Laboratories. (Bell Telephone Laboratories Photograph)

Left: INSIDE THE HOLMDEL RADIO TELESCOPE: Penzias is shown here taping the joints of the 20-foot horn antenna at Holmdel, with Wilson looking on. This was part of an effort to eliminate any possible source of electrical noise from the antenna structure that might account for the 3° K microwave static observed in 1964-65. All such efforts only succeeded in reducing the observed microwave noise intensity very slightly, and the conclusion became inescapable that this microwave radiation is really of astronomical origin. (Bell Telephone Laboratories Photographs)

The Sept 29th 1963
Gamow Dacha
785 - 6th Street
Boulder, Colorado

Dear Dr. Penzias,

~~Send~~ Thank you for sending me your paper on 3 °K radiation It is very nicely written except that "early history" is not "quite complete". The theory of, what is now knows as "primeval fireball" was first developed by me in 1946 (Phys. Rev. 70, 572, 1946; 74, 505, 1948; Nature 162, 680, 1948). The prediction of the numerical value of the present (residual) temperature could ~~can~~ be found in Alpher & Hermann's paper (Phys. Rev. 75, 1093, 1949) who estimate it as 5 °K, and in my paper (Kong Dansk. Ved. Sels. 27 no 10, 1953) with the estimate of 7 °K. Even in my popular book "Creation of Universe" (Viking 1952) you can find (on p. 42) the formula $T = 1.5 \cdot 10^{10}/t^{1/2}$ °K, and the upper limit of 50 °K. Thus, you see the world did not start with almighty Dicke.

Sincerely G. Gamow

図 3　G. ガモフの手紙．奇妙なことに日付が間違って 1963 年となっている*.

部分が第一回転準位にあることが観測された．この準位と基底状態のエネルギー差は波長 2.64 mm に対応している．原子スペクトルについてのヘルツベルグの標準的な本 (Herzberg 1950) の中に次のような記述がある："$K=0$ と $K=1$ の線の強度比から回転温度 2.3°K が結論される．この温度はもちろん限られた意味をもつにすぎない"．

CN

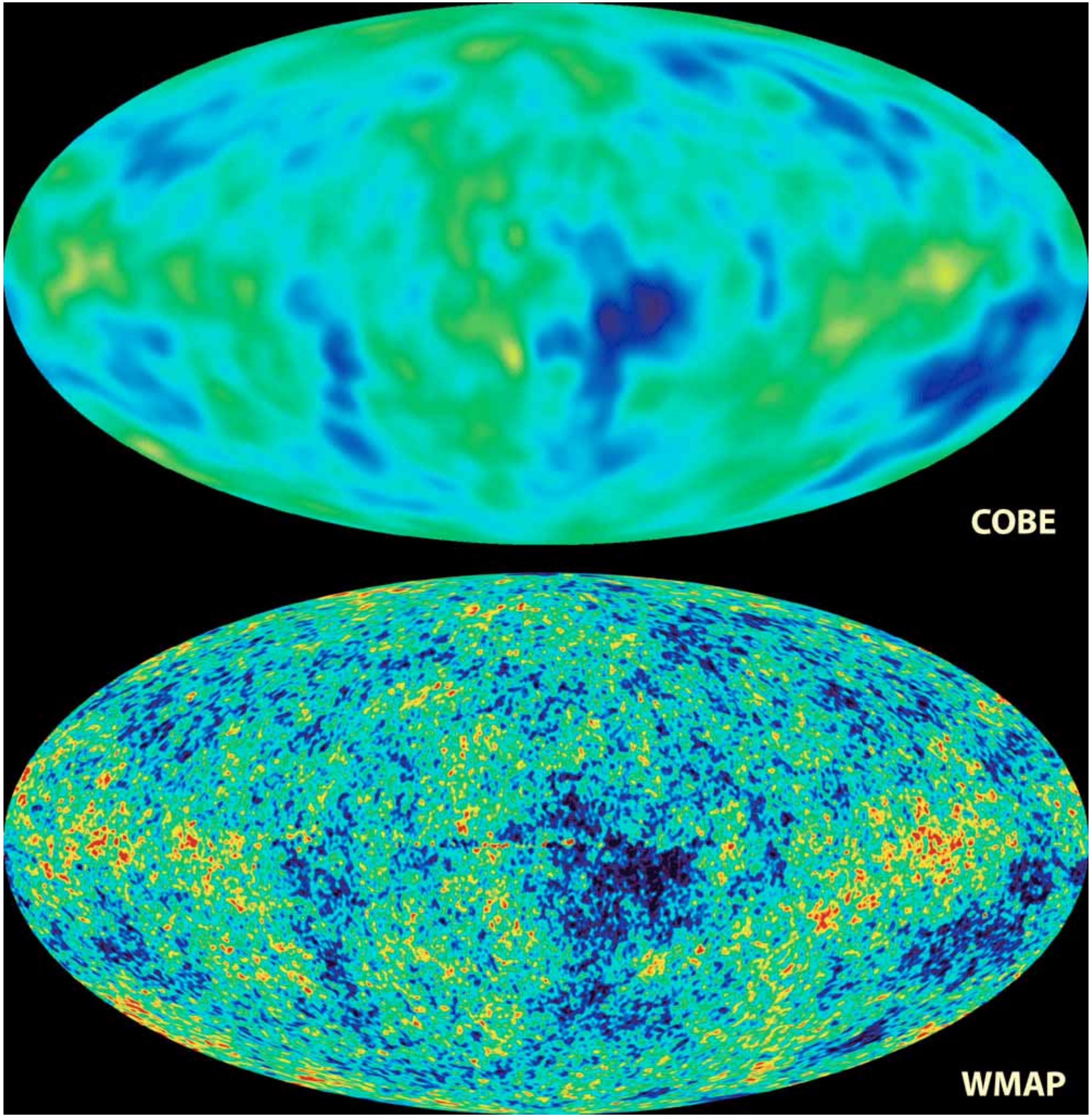
McKellar, 1941

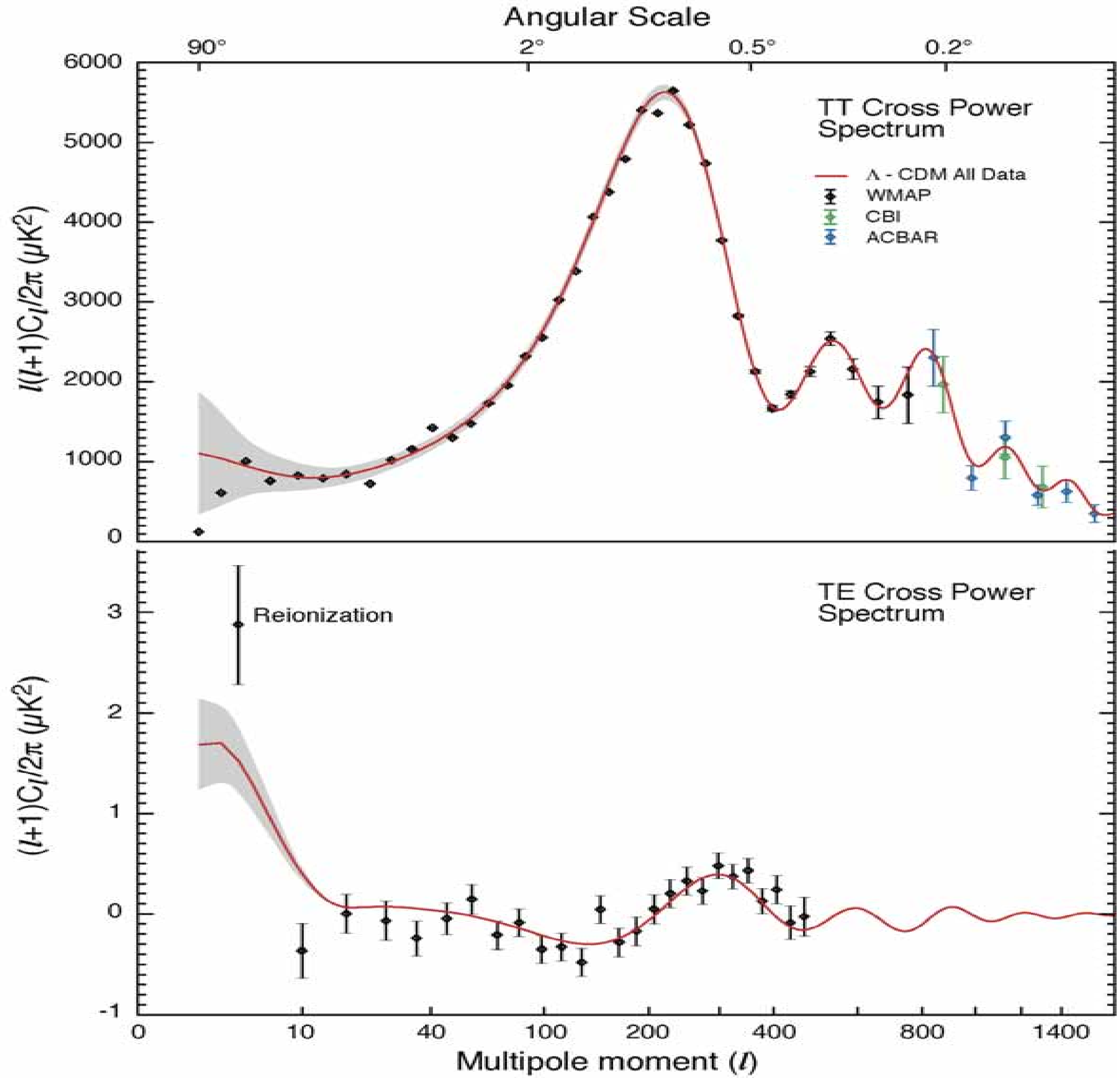
ところが，衝突による励起や，より高い準位からの遷移ではこのような占有比を維持し

$E(J=1) - E(J=0)$
$\sim (2{,}64\,\mathrm{mm})^{-1}$

* 訳は章末にある．

COBEとの比較





温度揺らぎ

WMAP
の観測

温度揺らぎ
と偏光の
相関

Cosmological Parameters

$$h = 0.72 \pm 0.05 (\mathrm{WMAP})$$

$$= 0.71^{+0.04}_{-0.03} (\mathrm{all})$$

$$\Omega_M h^2 = 0.14 \pm 0.02 (\mathrm{WMAP})$$

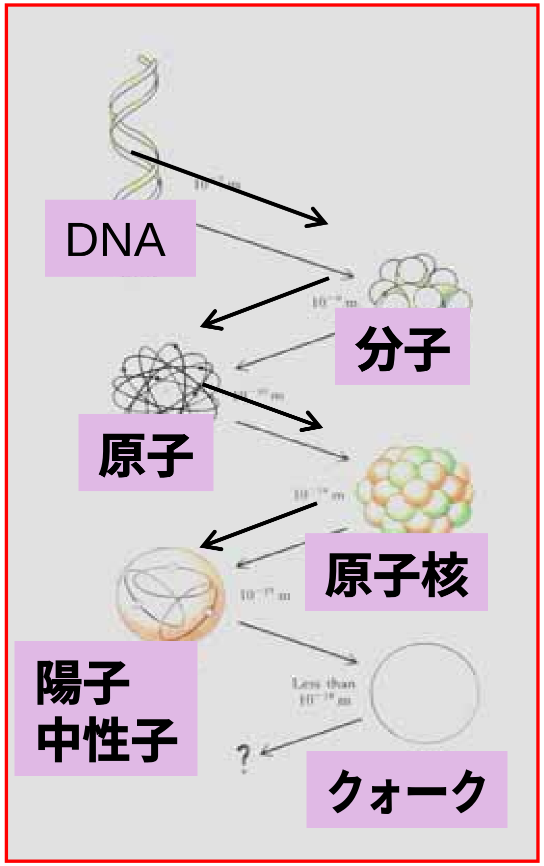
$$= 0.135^{+0.008}_{-0.009} (\mathrm{all})$$

$$\Omega_B h^2 = 0.024 \pm 0.001 (\mathrm{WMAP})$$

$$= 0.0224 \pm 0.0009 (\mathrm{all})$$

$$\Omega_{tot} = 1.02 \pm 0.02 (\mathrm{WMAP} + \mathrm{SN}, \mathrm{or}, \mathrm{HST}, \mathrm{2DF})$$

物質の究極要素


DNA
分子
原子
原子核
陽子
中性子
クォーク
?


クォーク
レプトン
アップ・
クォーク
ダウン・
クォーク
電子
電子ニュートリノ
第1世代
チャーム・
クォーク
ストレンジ・
クォーク
ミューオン
ミューオン・
ニュートリノ
第2世代
トップ・
クォーク
ボトム・
クォーク
タウ
タウ・
ニュートリノ
第3世代

クオーク と レプトン

3世代、それぞれ2種類、計6種類のクオークとレプトンの世界

第1世代

$$\begin{pmatrix} \text{電子} \\ \text{電子ニュートリノ} \end{pmatrix} \Updownarrow \quad \text{弱い相互作用}$$

$$\begin{pmatrix} u & u & u \\ d & d & d \end{pmatrix}$$

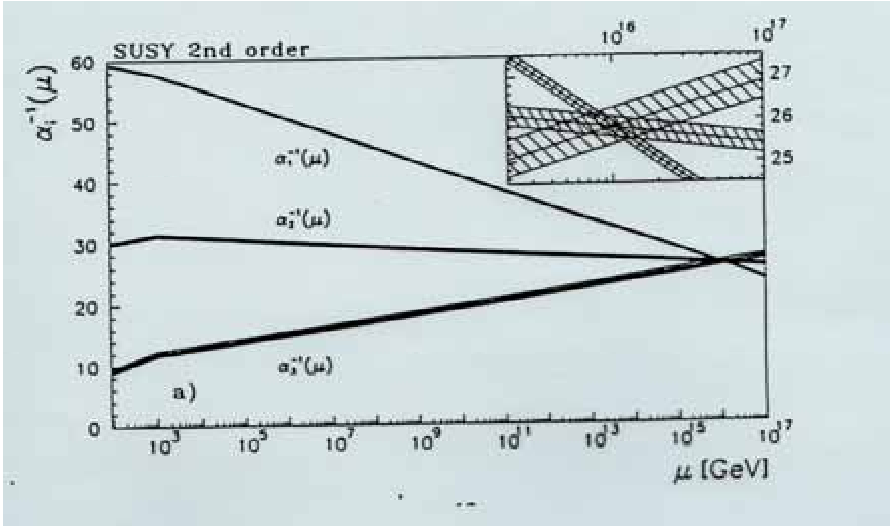
強い相互作用

$$\Leftrightarrow$$

Recent developments in particle physics

2 hints towards unification $@ \;\square\; 10^{15}\ GeV$

SUSY coupling unification



Neutrino mass via seesaw

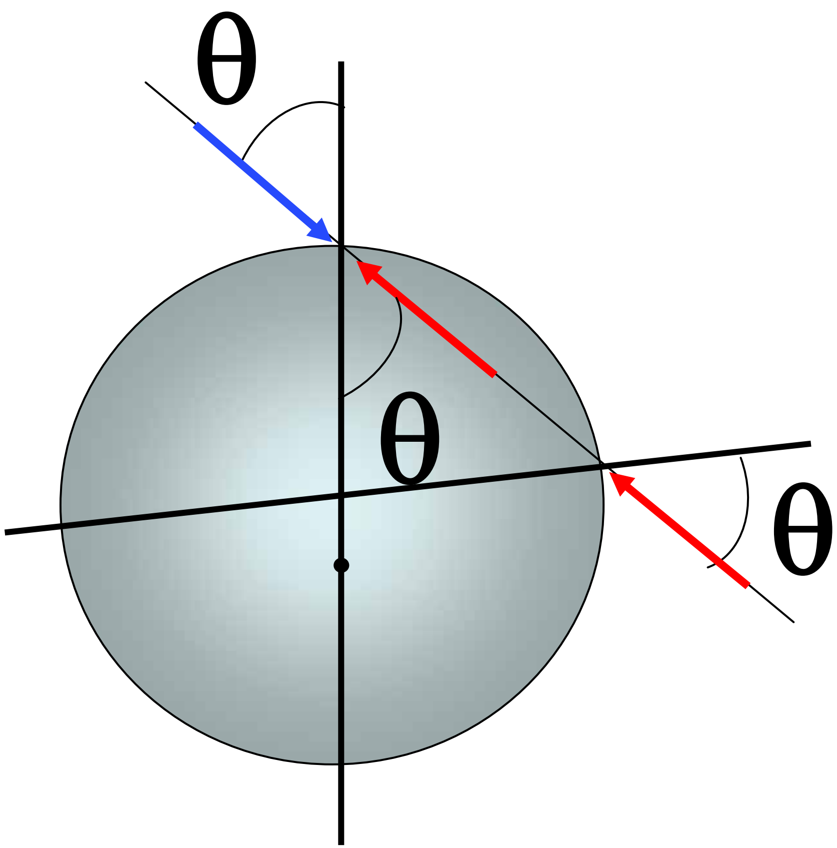
$$m_\nu = \frac{m_{q,l}^{\;2}}{M_{new\ physics}}$$

$$\Rightarrow \quad M_{new\ physics} = \frac{(100\,GeV)^2}{10^{-2}\ eV} = 10^{15}\ GeV$$

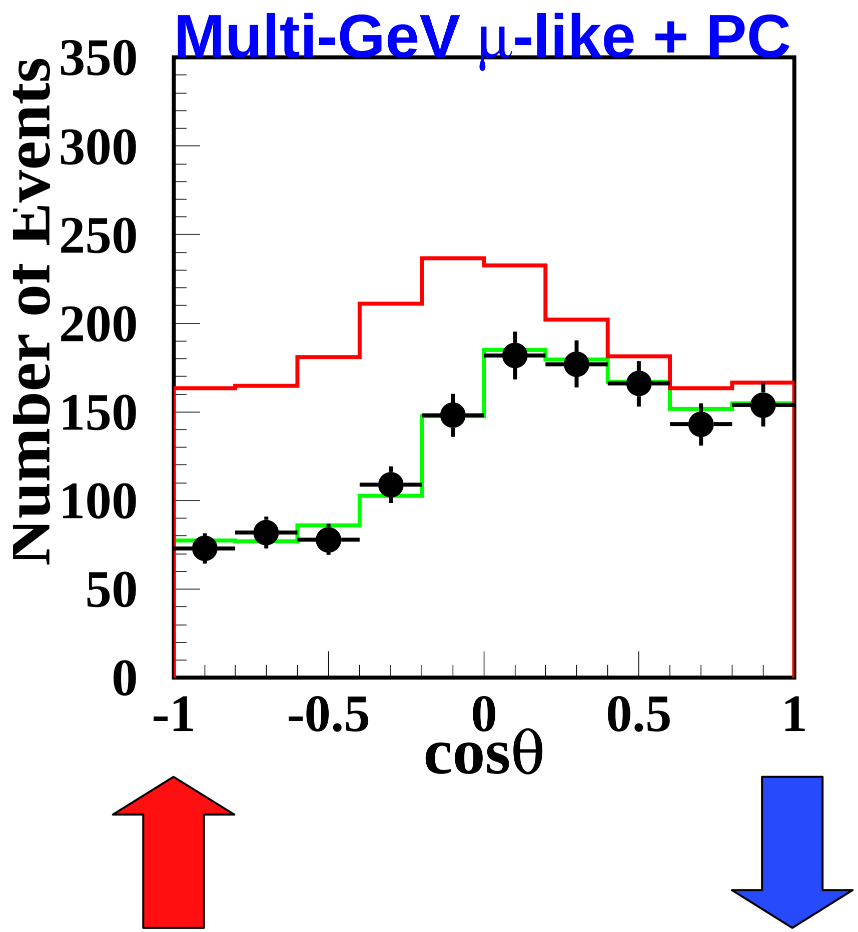
上下非対称性の発見

強度計算に依存しない

ニュートリノ振動の距離依存性を観測　(〜波長の観測)

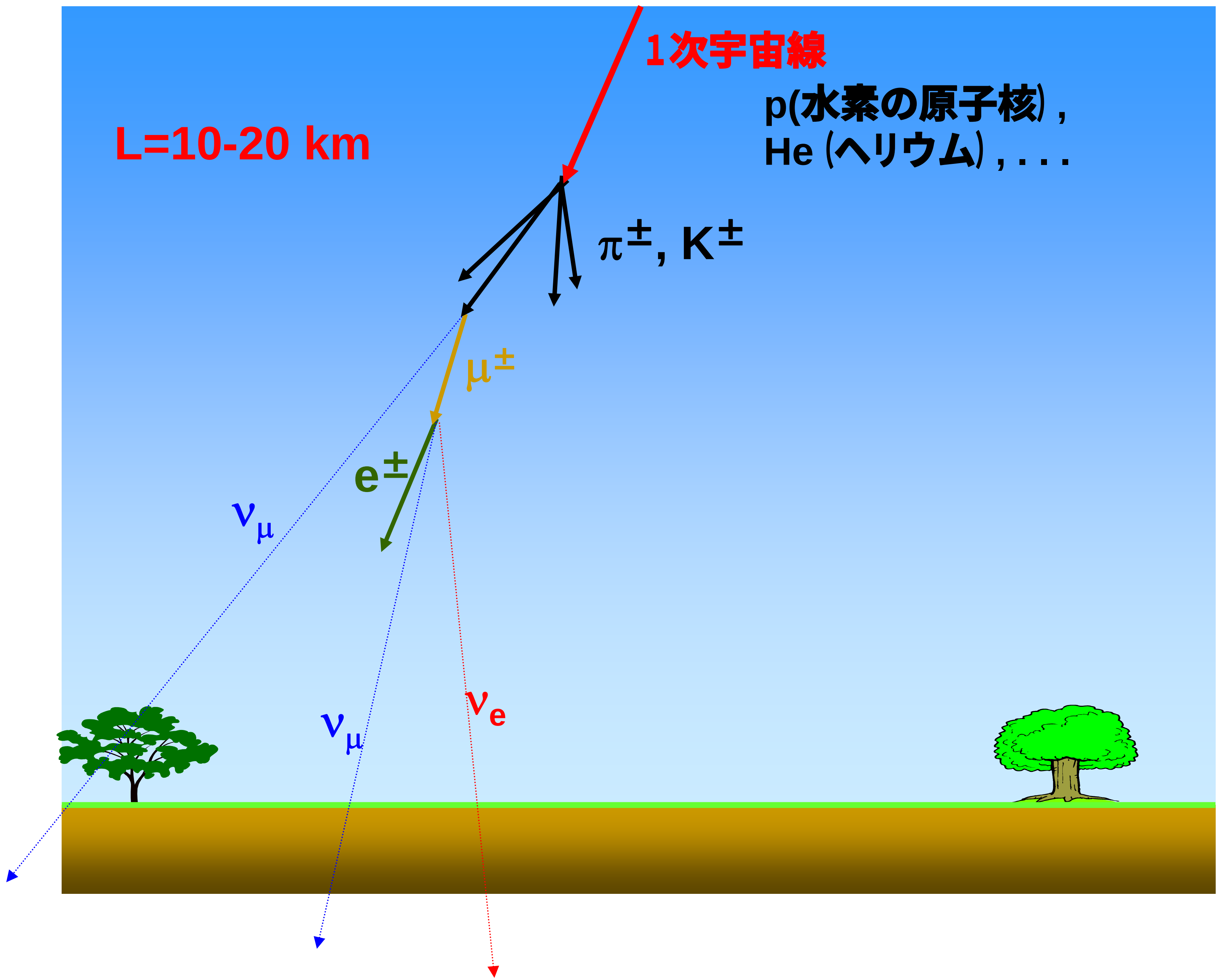


大気ニュートリノの角度分布
振動していなければ、上下対称である。



上向き　下向き

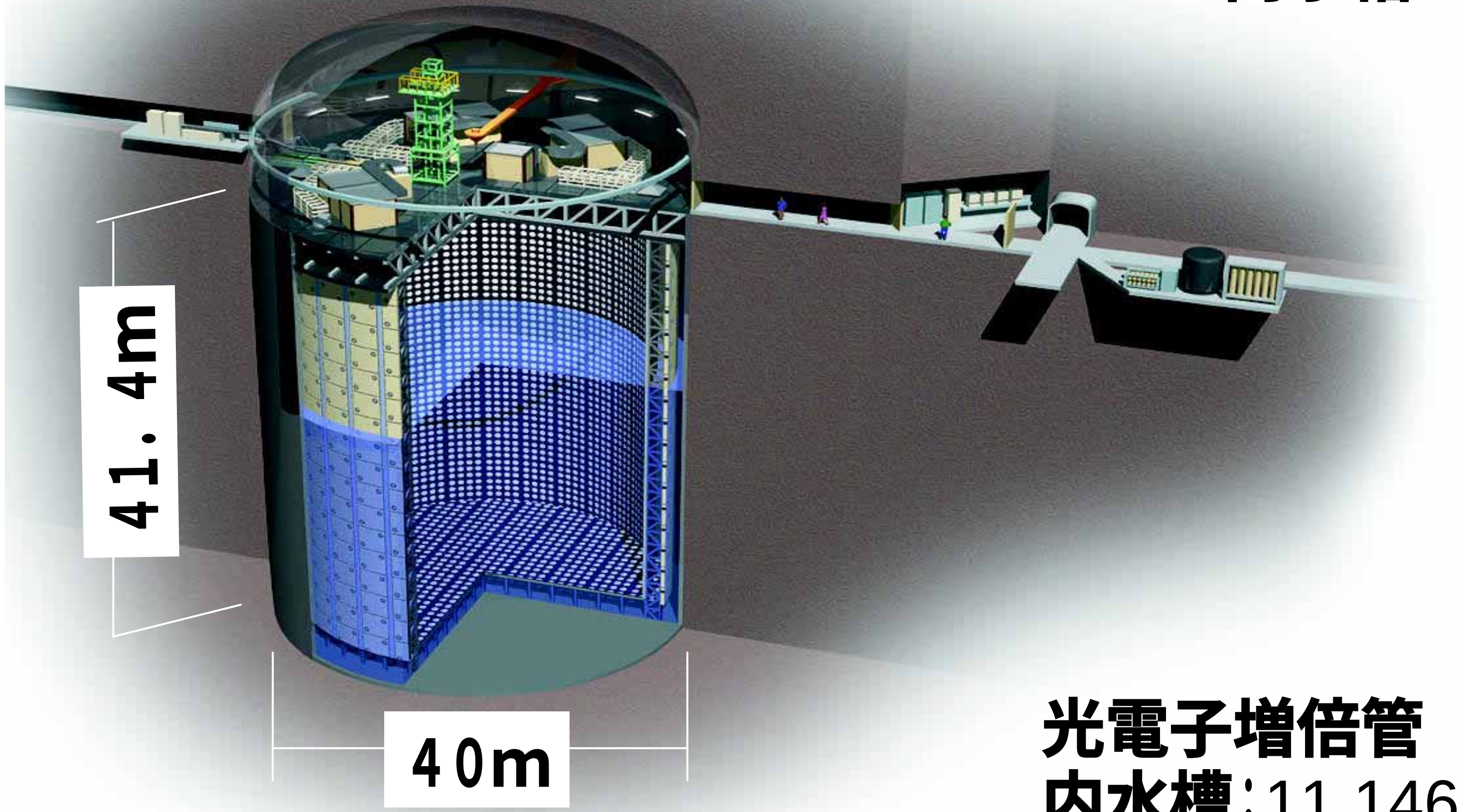
大気ニュートリノ



スーパーカミオカンデ

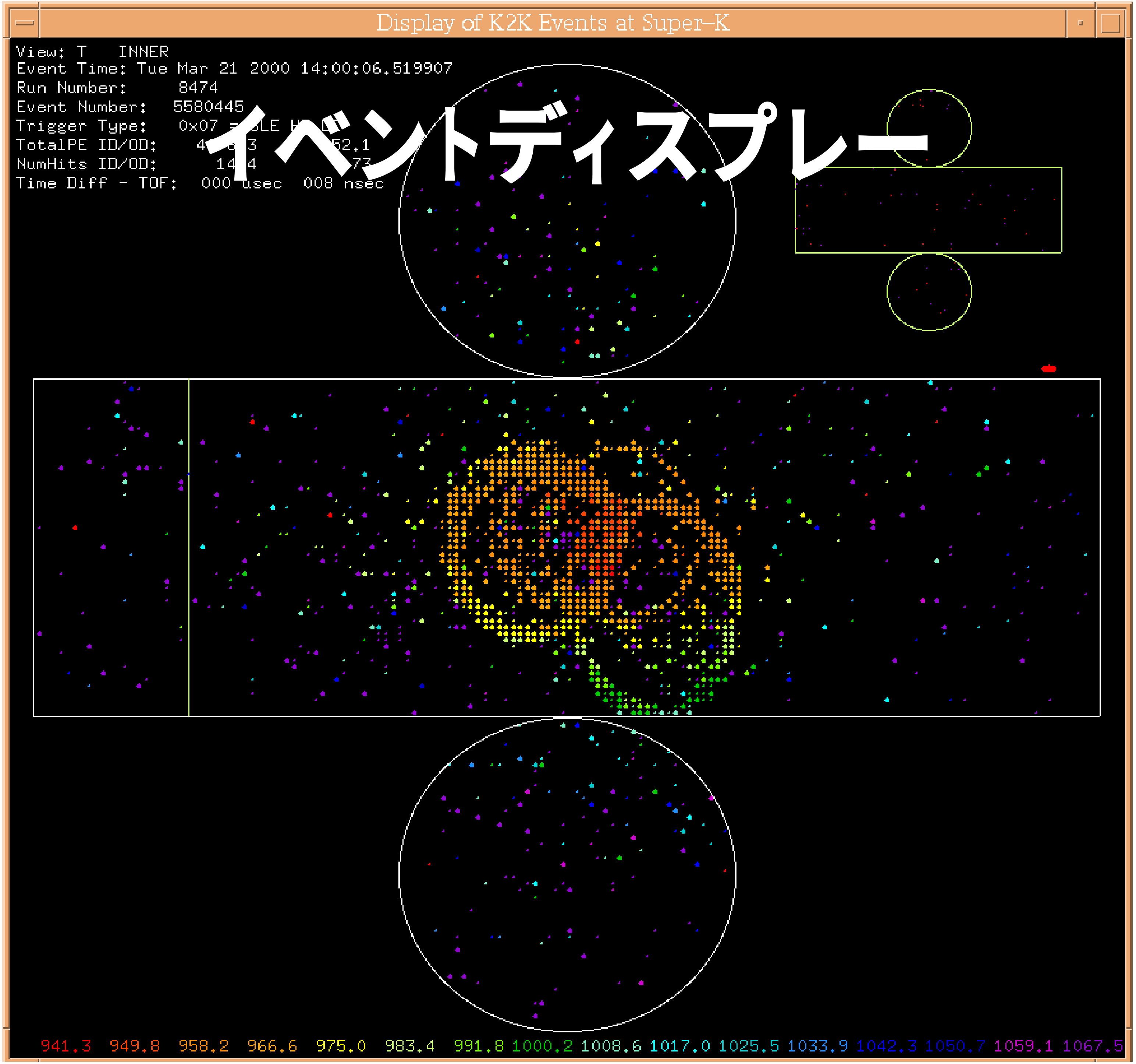
（１９９６年4月完成）

全質量　：　50,000 トン
内水槽　：　32,000 トン



光電子増倍管
内水槽：11,146本(直径50cm)
外水槽：　1,850本(直径20cm)

イベントディスプレー


Display of K2K Events at Super-K
View: T INNER
Event Time: Tue Mar 21 2000 14:00:06.519907
Run Number: 8474
Event Number: 5580445
Trigger Type: 0x07
TotalPE ID/OD:
NumHits ID/OD:
Time Diff - TOF: 000 usec 008 nsec
941.3 949.8 958.2 966.6 975.0 983.4 991.8 1000.2 1008.6 1017.0 1025.5 1033.9 1042.3 1050.7 1059.1 1067.5

大統一のもう一つの要素

- バリオン数の非保存

クオークとレプトンの同一視

結果：

陽子崩壊
宇宙の物質・反物質不均衡

Proton instability inevitable ?

No, but • • •

Progress of Theoretical Physics, Vol. 58, No. 3, September 1977

Muon Number Nonconservation in a Unified Scheme of All Interactions

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(Received April 30, 1977)

We present a unified gauge model based on the group $SU(6)$ that contains as commuting subgroups the $SU(3)\times U(1)$ of weak interactions and the color $SU(3)$ of strong interactions. In this scheme nonconservation of separate electron and muon number naturally arises through mixing of flavors in the quark sector, while the proton remains stable.

It was pointed out some time ago[1] that a unified gauge model of all interactions based on a simple group possesses an attractive feature of explaining both the smallness of the fine structure constant and the asymptotic freedom presumably needed to reproduce simple results of the naive quark parton model. A remarkable property of such a scheme is that it necessarily involves weak gauge bosons mediating transitions between a lepton and a quark. Exchange of this class of weak bosons generally causes nonconservation of separate electron and muon number due to different orientations of flavors in the lepton and quark sector. In the original $SU(5)$ model of Georgi and Glashow[1] the same exchange is also responsible for proton decay which, then, must be enormously suppressed, making the muon number nonconservation negligible for all practical purposes. This situation is drastically changed in a scheme in which the proton is forbidden to decay by some natural conservation law. The process such as $\mu\to e\gamma$ may then occur with a rate somewhat smaller than that of the present experimental limit. In this note we shall give an example of such models by unifying the vectorlike $[SU(3)\times U(1)]_W$ weak interaction model[2] recently proposed by the present author and the color $SU(3)_C$ gauge theory[3] into a $SU(6)$ gauge scheme. The model thus unified predicts decays such as $\mu\to e\gamma$ and $K\to\mu e$ with similar rates. It is entirely possible[4] that these decays are made to occur in much the same way as $\lambda\to n\gamma$ and $\bar{\lambda}n\to l\bar{l}$ by introducing an arbitrary Cabibbo-like mixing within the lepton sector, but we would find it more attractive if the existing mixing in the quark sector implies the muon number nonconservation in the other sector.

In a couple of previous papers[2,5] we showed that the $SU(3)\times U(1)$ may be a relevant weak interaction group beyond the minimal standard scheme.[6] The group $SU(6)$ then emerges as a natural choice of the unified group. In classifying fundamental fermions we demand that the leptons l and the quarks q respectively belong to $(\underline{3},\underline{1})$ and $(\underline{3},\underline{3})$ representations of the subgroup $SU(3)_W\times SU(3)_C\subset SU(6)$. The simplest representation of $SU(6)$ that contains both is $\underline{15}$, which

and (RR) since the former imply the same helicities to the lepton and antilepton, hence are kinematically favored. The amplitude for K_L-decay is recast using h_i's defined in Table I,

$$\tfrac{1}{2}\tilde{G}_F c_K[h_1(\bar{\mu}e+\bar{e}\mu)+h_2(\bar{\mu}\gamma_5 e-\bar{e}\gamma_5\mu)+h_3\bar{e}e+h_4\bar{\mu}\mu], \tag{7}$$

where $c_K=\langle 0|\bar{n}\gamma_5\lambda|K^0\rangle$. We have included contribution of this mechanism to the ordinary lepton pair decay, $e\bar{e}$ and $\mu\bar{\mu}$. A similar result for K_S-decay is also easily derived. Equation (7) yields decay rates normalized to $K^+\to\mu^+\nu_\mu$ with decay constant f_K, $(c_K\tilde{G}_F/2f_K m_\mu G_F\sin\theta)^2$ times $(h_1^2+h_2^2)$, $1.10h_3^2$, $0.86h_4^2$ for $K_L\to\mu\bar{e}$ (or $\bar{\mu}e$), $e\bar{e}$, $\mu\bar{\mu}$, respectively. Numerical factors are due to available phase spaces. Unfortunately, this result contains an unknown factor c_K whose estimate would require detailed knowledge of strong interactions. As a crude estimate of order of magnitude we might attempt to use a result of the naive quark model combined with PCAC, $c_K=f_K m_K^2/(m_\lambda+m_n)$.

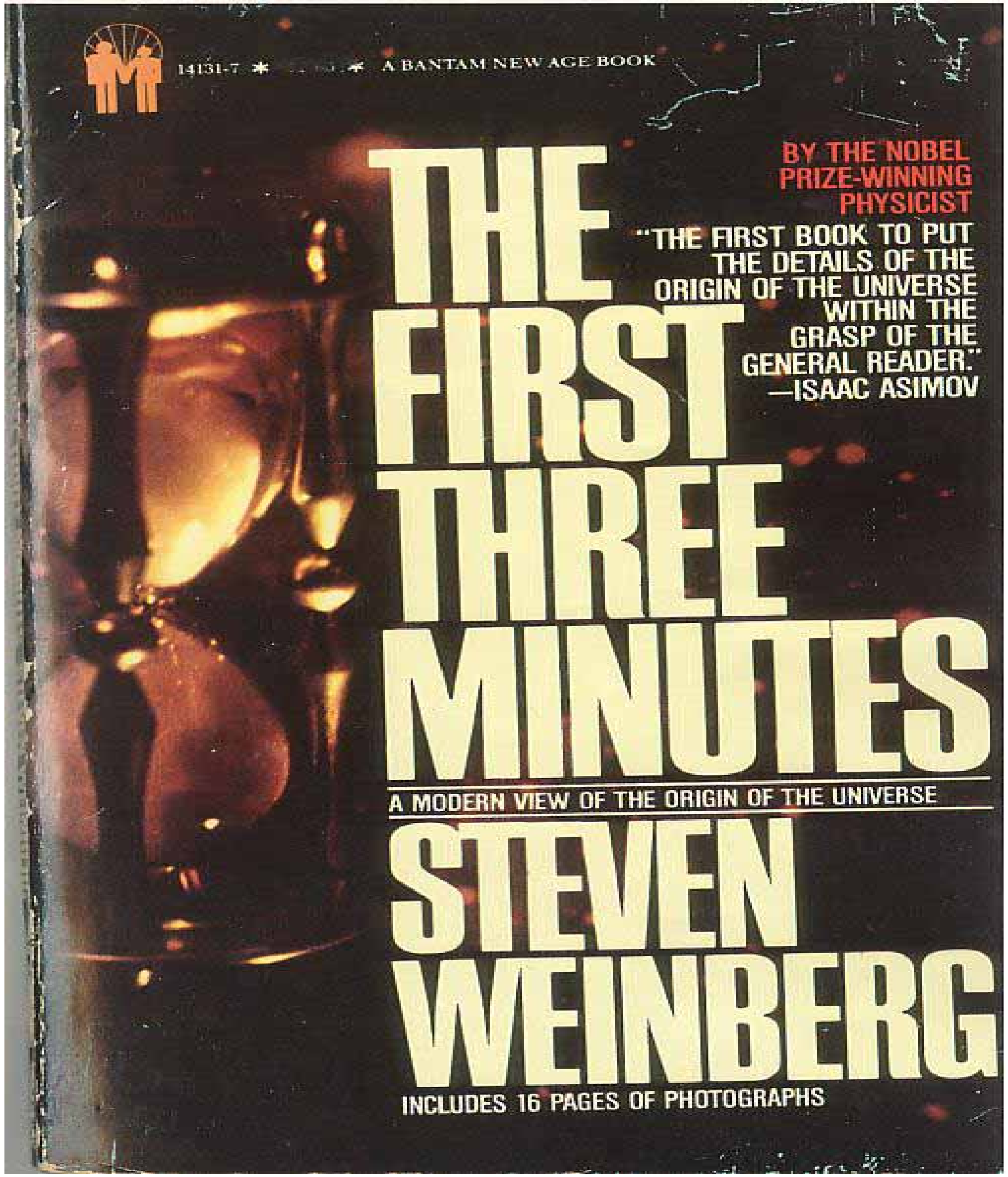
We now use experimental limits[8] to set bounds on the fundamental parameters of this model. The decay mode $K_L\to\mu\bar{\mu}$ is observed at the rate of branching ratio 10^{-8}, but the conventional contribution of 2γ intermediate states is expected to be of this order for $\mu\bar{\mu}$. We therefore ignore this mode in the following analysis. Numbers of experimental upper limits[8] used here are 2.2×10^{-8}, 2.0×10^{-9}, 2.0×10^{-9} for the branching ratios of $\mu\to e\gamma$, $K_L\to\mu\bar{e}\,(e\bar{\mu})$, $K_L\to e\bar{e}$, respectively. These give, together with the naive expectation of quark model, the following bounds,

$$\tilde{G}_F(f_1^2+f_2^2)^{1/2}/G_F m_\mu<2.5\times10^{-4},$$
$$\tilde{G}_F(h_1^2+h_2^2)^{1/2}m_K^2/G_F m_\mu(m_\lambda+m_n)<1.2\times10^{-5},$$
$$\tilde{G}_F h_3 m_K^2/G_F m_\mu(m_\lambda+m_n)<1.1\times10^{-5}. \tag{8}$$

Despite crude estimation of c_K it might not be too absurd to say that all the upper bounds to $\tilde{G}_F/G_F$ obtained from (8) are equal within a factor of 10 except the Cabibbo angle involved (see Table I). The Cabibbo factors are $\cos^2\theta:\sin^2\theta:\cos^2\theta$ for $(\mu\to e\gamma):(K_L\to\mu e):(K_L\to ee)$ in the models (A) and (D), and $\sin^2\theta:\cos^2\theta:\sin^2\theta$ in the models (B) and (C). If the process $\mu\to e\gamma$ is observed at a rate not significantly smaller than that of the present limit, we would then conclude that the models (B) and (C) are inconsistent with the present limit of $K_L\to\mu e$.

Another important consequence of the unified $SU(6)$ scheme is existence of a new class of hadrons, which we tentatively call superhadrons. These are $SU(3)_C$ singlet states that contain t's; baryons, $B^*(tqq)$, $B^{**}(ttq)$, $B^{***}(ttt)$ and mesons $M^*(t\bar{q})$, $\bar{M}^*(q\bar{t})$. These superhadrons are integrally charged and may be produced in association by strong or electromagnetic interactions. In their decay the $U(1)$ color charge $Q_{\rm col}$ defined by (3) is absolutely conserved. This conservation also assures stability of the proton unless a superhadron is less massive than that. These charges are $(0,1,0,0,-1,-2,-1)$ for $(l,B,M,B^*,B^{**},B^{***},M^*)$, respectively. It appears reasonable to assume that the least massive super-

Popular, but influential:
"The first three minutes"



From the first three minutes

0.06 years (or 22 days) for the temperature to drop to 10 million degrees, then it took another six years for the temperature to drop to one million degrees, another 600 years for the temperature to drop to 100,000 degrees, and so on. The whole time that it took the universe to cool from 100 million degrees to 3,000° K (i.e., to the point where the contents of the universe were just about to become transparent to radiation) was 700,000 years. (See figure 8.) Of course, when I write here of "years" I mean a certain number of absolute time units, as, for instance, a certain number of periods in which an electron makes an orbit around the nucleus in a hydrogen atom. We are dealing with an era long before the earth would begin its tours around the sun.

If the universe in the first few minutes was really composed of precisely equal numbers of particles and antiparticles, they would all have annihilated as the temperature dropped below 1,000 million degrees, and nothing would be left but radiation. There is very good

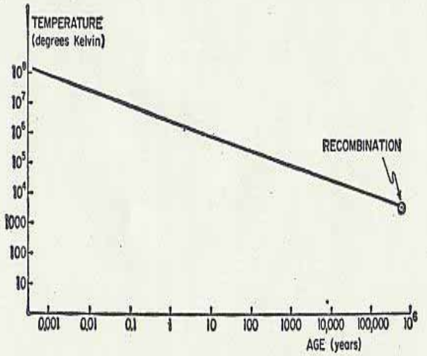


Figure 8. *The Radiation-Dominated Era.* The temperature of the universe is shown as a function of time, for the period from just after the end of nucleosynthesis to the recombination of nuclei and electrons into atoms.

evidence against this possibility—we are here! There must have been some excess of electrons over positrons, of protons over antiprotons, and of neutrons over antineutrons, in order that there would be something left over after the annihilation of particles and antiparticles to furnish the matter of the present universe. Up to this point in this chapter I have purposely ignored the comparatively small amount of this leftover matter. This is a good approximation if all we want is to calculate the energy density or the expansion rate of the early universe; we saw in the preceding chapter that the energy density of nuclear particles did not become comparable to the energy density of radiation until the universe had cooled to about 4,000° K. However, the small seasoning of leftover electrons and nuclear particles has a special claim to our attention, because they dominate the contents of the present universe, and in particular, because they are the main constituents of the author and the reader.

As soon as we admit the possibility of an excess of matter over antimatter in the first few minutes, we open up the problem of determining a detailed list of ingredients for the early universe. There are literally hundreds of so-called elementary particles on the list published every six months by the Lawrence Berkeley Laboratory. Are we going to have to specify the amounts of each one of these types of particle? And why stop at elementary particles—do we also have to specify the numbers of different types of atoms, of molecules, of salt and pepper? In this case, we might well decide that the universe is too complicated and too arbitrary to be worth understanding.

Fortunately, the universe is not that complicated. In order to see how it is possible to write a recipe for its contents, it is necessary to think a little more about what is meant by the condition of thermal equilibrium. I have already emphasized how important it is that the universe has passed through a state of thermal equilib-

to an infinite electric field. But whether the universe is open or closed, it is safe to say that the cosmic electric charge per photon is negligible.

The baryon number per photon is also easy to estimate. The only *stable* baryons are the nuclear particles, the proton and neutron, and their antiparticles, the antiproton and antineutron. (The free neutron is actually unstable, with an average life of 15.3 minutes, but nuclear forces make the neutron absolutely stable in the atomic nuclei of ordinary matter.) Also, as far as we know, there is no appreciable amount of antimatter in the universe. (More about this later.) Hence, the baryon number of any part of the present universe is essentially equal to the number of nuclear particles. We observed in the preceding chapter that there is now one nuclear particle for every 1,000 million photons in the microwave radiation background (the exact figure is uncertain), so the baryon number per photon is about one thousand-millionth (10^{-9}).

This is really a remarkable conclusion. To see its implications, consider a time in the past when the temperature was above ten million million degrees (10^{13} ° K), the threshold temperature for neutrons and protons. At that time the universe would have contained plenty of nuclear particles and antiparticles, about as many as photons. But the baryon number is the *difference* between the numbers of nuclear particles and antiparticles. If this difference were 1,000 million times smaller than the number of photons, and hence also about 1,000 million times smaller than the *total* number of nuclear particles, then the number of nuclear particles would have exceeded the number of antiparticles by only one part in 1,000 million. In this view, when the universe cooled below the threshold temperature for nuclear particles, the antiparticles all annihilated with corresponding particles, leaving the tiny excess of particles over antiparticles as a residue which would eventually turn into the world we know.

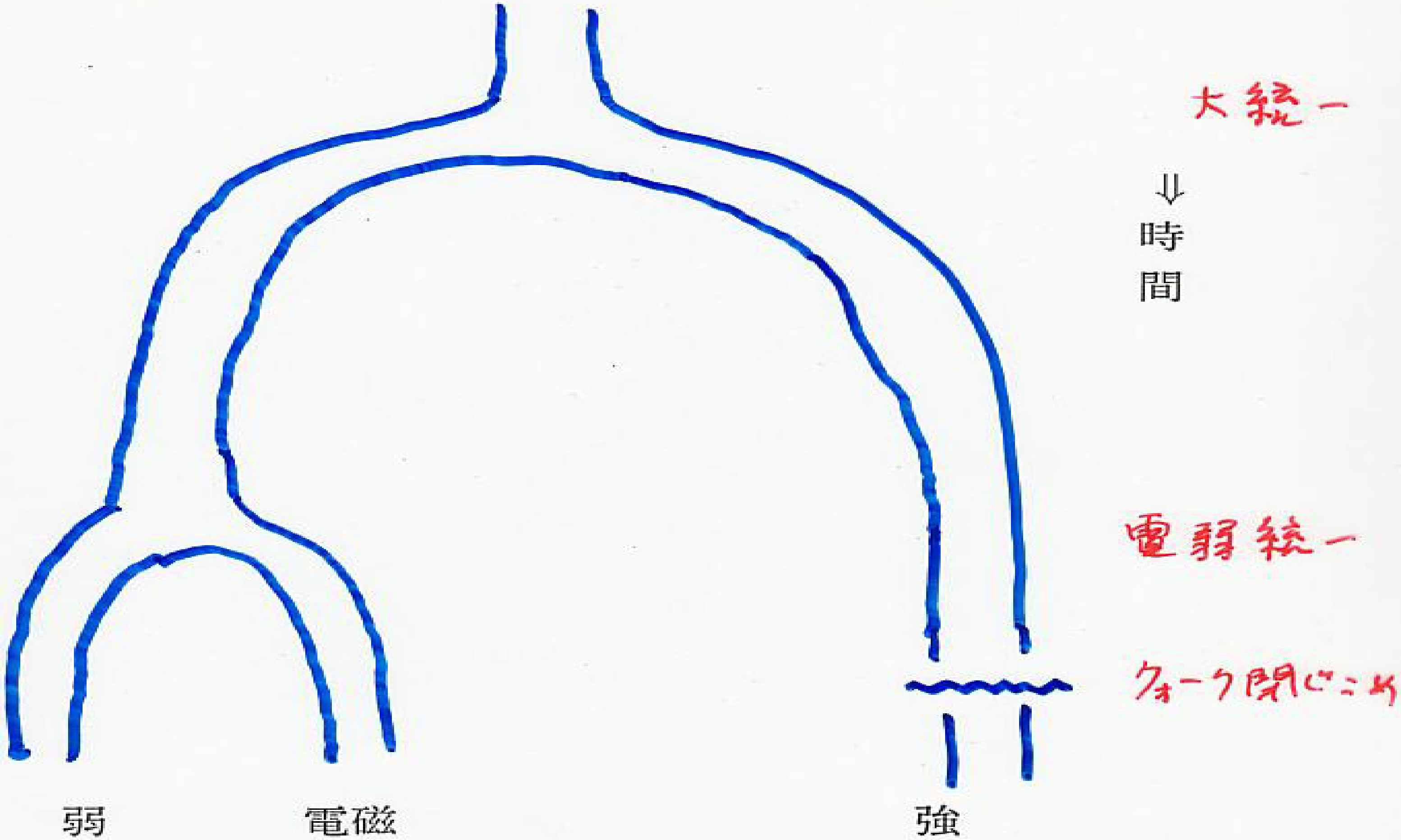
The occurrence in cosmology of a pure number as small as one part per 1,000 million has led some theorists to suppose that the number really is zero—that is, that the universe really contains an equal amount of matter and antimatter. Then the fact that the baryon number per photon appears to be one part in 1,000 million would have to be explained by supposing that, at some time before the cosmic temperature dropped below the threshold temperature for nuclear particles, there was a segregation of the universe into different domains, some with a slight excess (a few parts per 1,000 million) of matter over antimatter, and others with a slight excess of antimatter over matter. After the temperature dropped and as many particle-antiparticle pairs as possible annihilated, we would be left with a universe consisting of domains of pure matter and domains of pure antimatter. The trouble with this idea is that no one has seen signs of appreciable amounts of antimatter anywhere in the universe. The cosmic rays that enter our earth's upper atmosphere are believed to come in part from great distances in our galaxy, and perhaps in part from outside our galaxy as well. The cosmic rays are overwhelmingly matter rather than antimatter—in fact, no one has yet observed an antiproton or an antinucleus in the cosmic rays. In addition, we do not observe the photons that would be produced from annihilation of matter and antimatter on a cosmic scale.

Another possibility is that the density of photons (or, more properly, of entropy) has not remained proportional to the inverse cube of the size of the universe. This could happen if there were some sort of departure from thermal equilibrium, some sort of friction or viscosity which could have heated the universe and produced extra photons. In this case, the baryon number per photon might have started at some reasonable value, perhaps around one, and then dropped to its present low value as more photons were produced. The

宇宙の相転移

- 宇宙初期には，低温では自発的に破れる対称性が，高温のために回復.
- 力の分化とクォーク閉じ込めの開放.
- 時間の経過とともに，

  大統一 ⇒ 電弱統一 ⇒ クォーク閉じ込め



Baryo-genesis papers

VOLUME 41, NUMBER 5 PHYSICAL REVIEW LETTERS 31 JULY 1978

Unified Gauge Theories and the Baryon Number of the Universe

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(Received 27 April 1978)

I suggest that the dominance of matter over antimatter in the present universe is a consequence of baryon-number–nonconserving reactions in the very early fireball. Unified guage theories of weak, electromagnetic, and strong interactions provide a basis for such a conjecture and a computation in specific SU(5) models gives a small ratio of baryon- to photon-number density in rough agreement with observation.

It is known that the present universe is predominantly made of matter, at least in the local region around our galaxy, and there has been no indication observed[1] that antimatter may exist even in the entire universe. I assume here that in our universe matter indeed dominates over antimatter, and I ask within the framework of the standard big-bang cosmology[2] how this evolved from an initially symmetric configuration, namely an equal mixture of baryons and antibaryons. Since the baryon number is not associated with any fundamental principle of physics,[3] such an initial value seems highly desirable. I find in this paper that generation of the required baryon number is provided by grand unified gauge theories[4] of weak, electromagnetic, and strong interactions, which predict simultaneous violation of baryon-number conservation and *CP* invariance. More interestingly, my mechanism can explain why the ratio of the baryon- to the photon-number density in the present universe is so small, roughly of the order[2] of 10^{-8}–10^{-10}.

The essential point of my observation is that in the very early, hot universe the reaction rate of baryon-number–nonconserving processes, if they exist, may be enhanced by extremely high temperature and high density. In gauge models discussed below, the relevant scale of temperature is given by the grand unification mass around 10^{16} GeV where fundamental constituents, leptons and quarks, begin to become indistinguishable. This mass is high enough to make futile virtually all attempts to observe proton decay in the present universe: proton lifetime $\gg 10^{30}$ year.[5] Instead, if my mechanism works, we may say that a fossil of early grand unification has remained in the form of the present composition of the universe.

The laws obeyed by the hot universe at temperatures much above a typical hadron mass (~ 1 GeV) might, at first sight, appear hopelessly complicated because of many unknown aspects of hadron dynamics. Recent developments of high-energy physics, however, tell that perhaps the opposite is the case. At such high temperatures and densities hadrons largely overlap and an appropriate description of the system is given in terms of pointlike objects—quarks, gluons, leptons, and any other fundamentals. The asymptotic freedom[6] of the strong interaction and weakness of the other interactions further assure[7] that this hot universe is essentially in a thermal equilibrium state made of almost freely moving objects. I shall assume that this simple picture of the universe is correct up to a temperature close to the Planck mass, $G_N^{-1/2} \sim 10^{19}$ GeV, except possibly around the two transitional regions where spontaneously broken weak-electromagnetic and grand-unified gauge symmetries become re-

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ORIGIN OF COSMOLOGICAL BARYON ASYMMETRY

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A new version of cosmological baryon generation is examined. In this scheme the baryon asymmetry is caused by the nonequilibrium decay of an X (leptoquark) and $X^\dagger$ boson of $\sim 10^{16}$ GeV, which takes place after an equilibrium period of baryon nonconserving two-body reactions. This mechanism imposes a severe constraint on grand unified theories; both upper and lower limits to the unification mass are derived.

It has recently been suggested [1–3] that the apparent excess [4] of cosmological baryons over antibaryons may be explained by baryon nonconserving processes near the Planck time. As an interesting consequence of such a mechanism the observed ratio [4] (10^{-8}–10^{-10}) of baryon to photon number may be related to microscopic quantities that characterize these processes, hence the ratio may not be an arbitrary parameter in cosmology. Besides the baryon, *C* and *T* noninvariance necessary for this explanation, recent investigations, as particularly emphasized in refs. [5–8], have made it clear that departure from thermal equilibrium is essential to generate a net baryon number. Reactions among light (masses $\ll$ leptoquark bosons) fermions considered in ref. [1] seem hardly able to produce an appreciable amount of baryon excess because the likely effect of nonequilibrium reactions of almost massless particles will simply be to red-shift [6] the temperature in an expanding universe. This can easily be seen [9] by using Boltzmann-like equations linearized in the deviation from the thermal distribution and taking into account the unitarity constraint [6–8]. As an alternative to light fermion reactions, Weinberg [8] and also Toussaint et al. [6] have more recently suggested that the decay of heavy leptoquark bosons (X's) may produce an appreciable baryon asymmetry due to the presence of a threshold. An important point to note is that below the threshold the decay dominates over its inverse decay and the linearized rate equation cannot be applied unlike the previous case. In Weinberg's scheme gravitational thermal equilibrium at the Planck time is also assumed to explain the initial condition of vanishing baryon number. The purpose of this note is to elaborate on this decay mechanism and propose as a new mechanism of the initial condition light fermion reactions in equilibrium that precede the X-decay. This makes it unnecessary to assume the presence of gravitational thermal equilibrium, and gives an interesting constraint on the parameters of grand unified theories (GUT) [10,11], independently of the details of specific models. Our mechanism strongly suggest that the relation m_X (mass of leptoquark gauge boson) $\approx \alpha^2 m_p$ (the Planck mass $\approx 1.2 \times 10^{19}$ GeV) is not accidental, but a necessary consequence of baryon generation.

Consider below its threshold m_X the decay of leptoquark bosons X into diquarks $\overline{q}\overline{q}$ and leptoquarks $q\ell$, and the inverse process. In this paper we assume exact conservation of color, hence a color triplet X may dominantly decay into the modes just listed. In principle, X can be a Higgs boson, but we ignore this possibility by assuming that a leptoquark gauge boson gives a more important contribution, which is true if the masses of fermions (quarks or leptons) $\ll$ mass of the ordinary weak boson (≈ 100 GeV). Also for simplicity, we do not consider the case that the mass of the colored Higgs boson $H \ll m_X$, hence $m(H) \approx m_X$ in the following. The effective annihilation rate of X

YKIS lecture　　著書

KEK-TH 30

August 1981

COSMOLOGICAL BARYON PRODUCTION AND RELATED TOPICS

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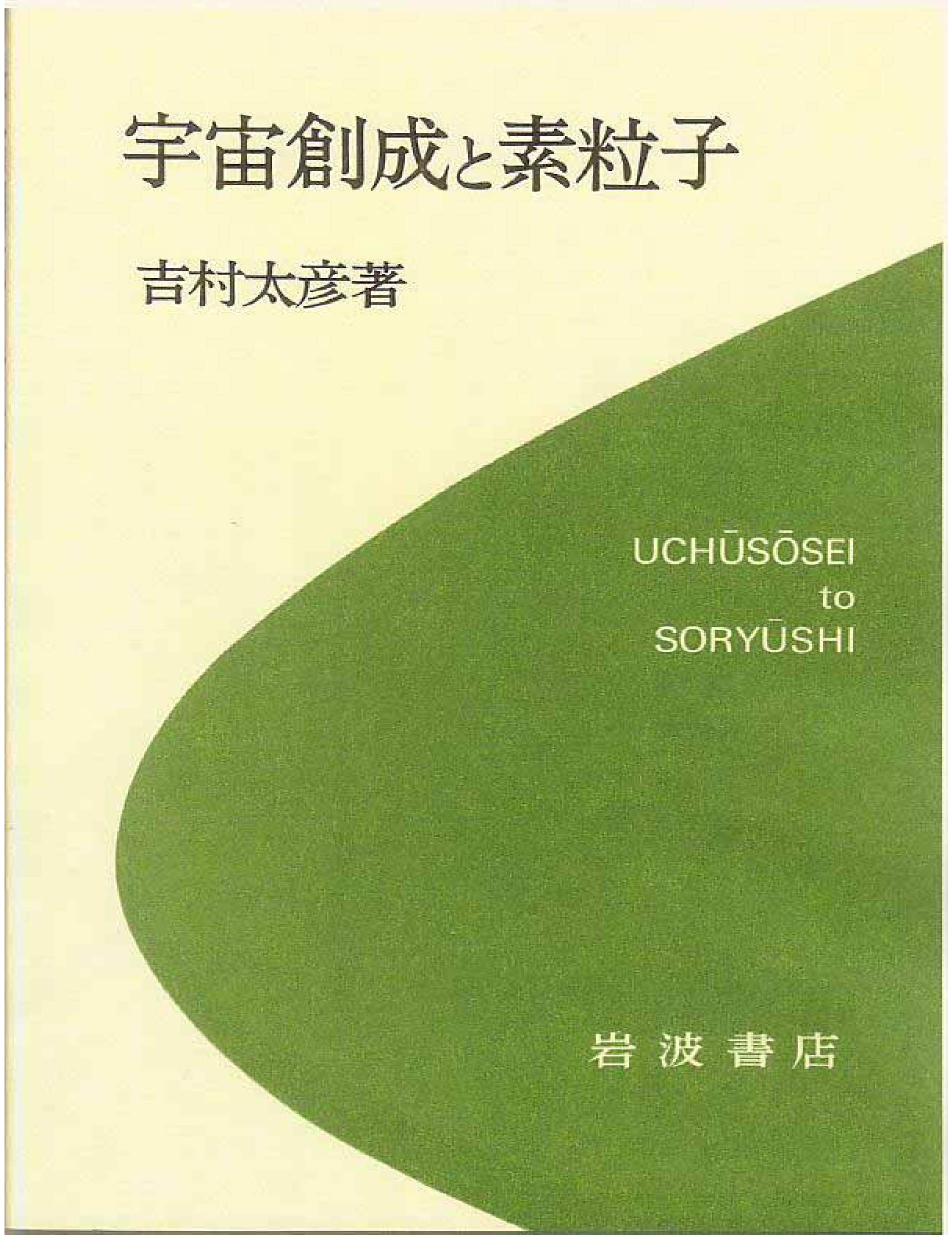
Lectures delivered at 4th Kyoto Summer Institute on Grand Unified Theories and Related Topics, June 29 - July 3, 1981 (Kyoto), and to be published in Proceedings.



Sakharov’s response to my paper

To dear M. Yoshimura.
Department of Physics
Tohoku University, Sendai 980
Japan

3/VI A.C.

Baryon asymmetry of the Universe

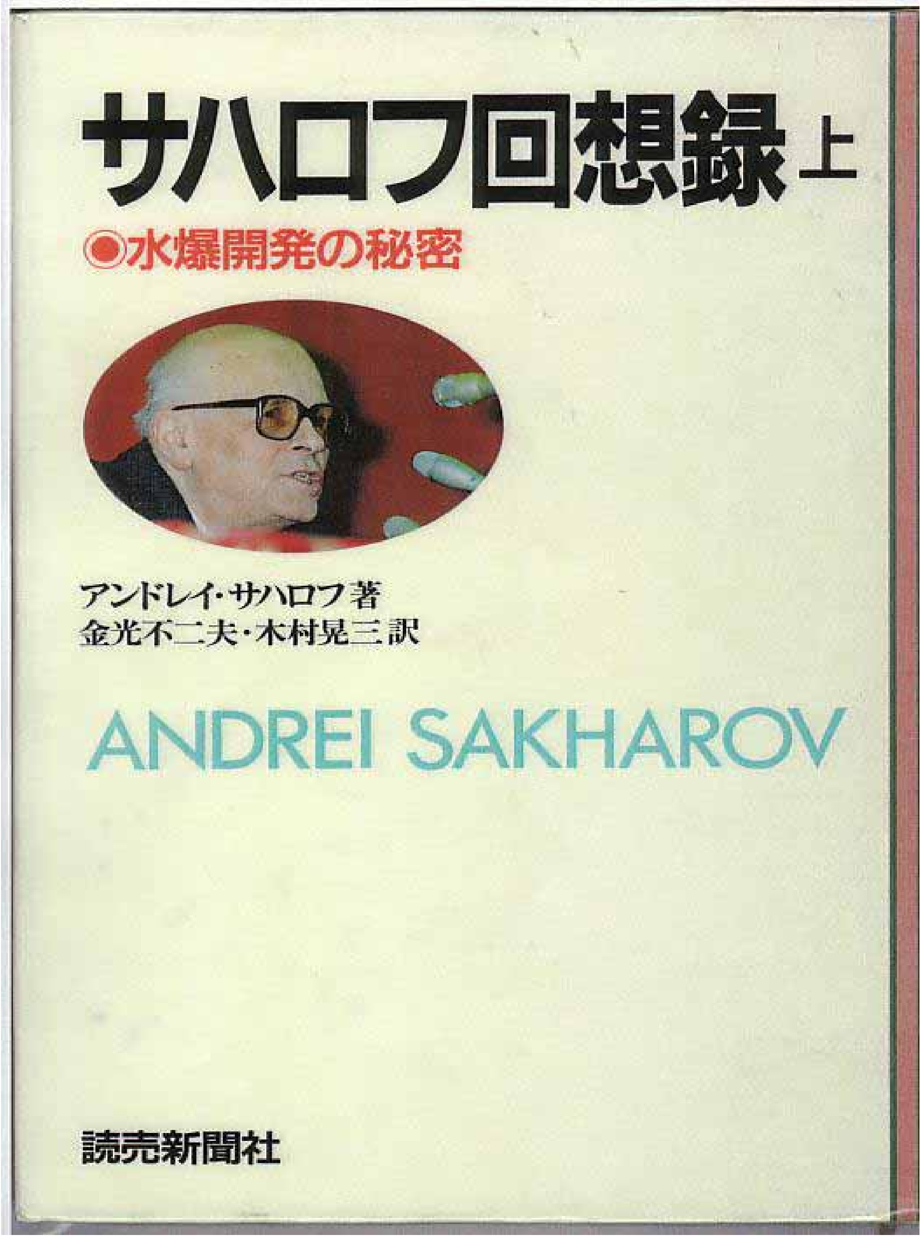
A.D.Sakharov
P.N.Lebedev Physical Institute, Academy of Sciences of the USSR.

A possible process of the appearance of baryon and anti-lepton excess at the early stage of the charge-neutral hot Universe expansion in the unified gauge theory of strong, weak and electromagnetic interactions is discussed. By the estimate presented here the baryon asymmetry $A = \frac{N_B}{N_\gamma}$ (the ratio of the mean baryon density to the relic radiation quantum density, to an accuracy of the numerical factor equal to the ratio of the number of baryons to the initial entropy of the hot Universe, in the same co-moving volume) is equal, in the order of magnitude to $A \sim \alpha^3 \vartheta^3 \delta_c$. The value $\alpha = g^2$ is the gauge field interaction constant, ϑ is the quantity of the order of the Cabibbo angle, δ_c is the phase of complex quark mixing. The numerical coefficient in this formula may contain an additional small parameter. Some considerations are expressed concerning the many-sheet model of the Universe suggested before by the author.

I. Introduction. Estimation of the effect.

In 1966 the author expressed the supposition that the observed baryon (and supposed lepton) asymmetry of the Universe appears at the early stage of cosmological expansion from the initial neutral charge state. Such a process is possible due to the effects of CP-invariance violation in non-stationary conditions of expansion if the baryon and lepton charge conservation is supposed to be broken [1].

Historical account



を否定したり、クォークが整数の電荷を持つという理論を支持すると言う必要がなかったことに気が付いた。

量子色力学

実際には、クォークが分数の電荷を持つというはるかに美しい理論が正しいことは、その当時でさえすでにほぼ確定的だった。その後の発展は、強い相互作用に関する力学理論である、いわゆる量子色力学を含むこの理論を立証した。この理論ではクォークは色と呼ばれる自由度を与えられ、量子色力学という名前もここからきている。量子力色学（ＱＣＤ）は、質量その他のハドロン（バリオンと中間子）の性質を記述する上で大きな成功をおさめた。ＱＣＤの重要な特徴は、クォークの閉じこめである。すなわちクォークは、原子から電子を分離するようには、バリオンあるいは中間子から分離することができない。それは、距離とともに減少しない力によってクォークを保持する「弦」ができるからである。

私は、自分の失言を訂正しようと、前日の外国人科学者を会議場で探したが、彼の名前を知らなかったのでついに探すことができなかった。

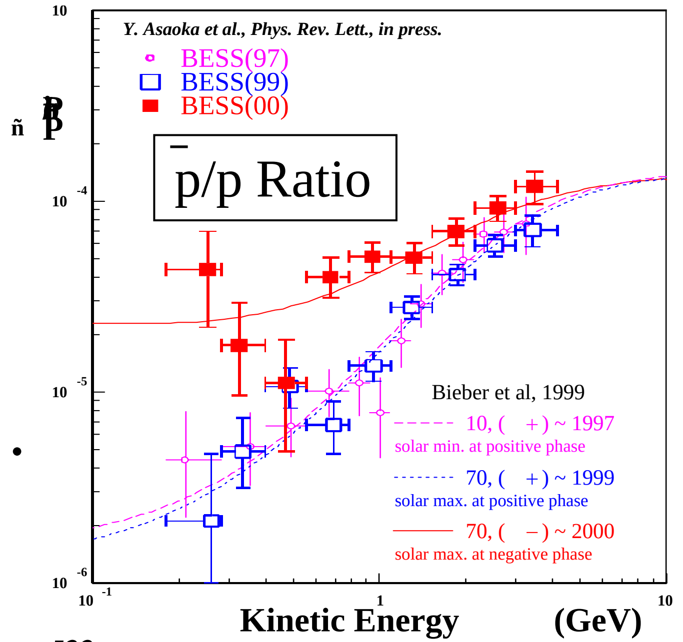
七八年に吉村太彦は、大統一理論を宇宙のバリオン非対称性に結びつけた重要な論文を発表した。これは私もやればできたはずなのに、やらなかったものだ。大統一理論とバリオン非対称性との関係の問題については、吉村以前にすでにソ連のイグナチエフ、クラスニコフ、クジミン、タフヘリージェが国際学会で発表していたが、しかしいちばん広く認められたのは吉村の論文である。これらの論文は強い印象を与え、これに刺激されてとりわけＸボゾン、Ｙボゾンやいわゆ

Absence of antimatter and problem with symmetric cosmology

- Observational evidence against symmetric cosmology

$$\frac{\overline{He}}{He} < 10^{-6} \quad \text{near earth}$$

low energy $\bar{p}$ spectrum



No evidence of γ from $\overline{N} + N \to \pi^0 + \cdots$

- Theoretical problem with B-symmetric cosmology

$$\frac{n_B}{n_\gamma} = \frac{n_{\bar{B}}}{n_\gamma} = \frac{O[100]}{m_N m_{pl} \langle \sigma v \rangle_{N\bar{N}}} \approx 10^{-18} \qquad @T \approx \frac{m_N}{50}$$

much smaller than observed 10^{-10}

No working model of domain separation

Generation of B-asymmetry

- Key quantity

$$\left(\frac{n_B}{n_\gamma}\right)_{\text{after annihilation}} = O[1] \times \left(\frac{B-\overline{B}}{B+\overline{B}}\right)_{\text{before annihilation}}$$

Observation $\quad \dfrac{n_B}{n_\gamma} = O\left[10^{-10}\right]$

imply 1 excess of B out of 10^{10} pairs

How to produce the asymmetry：
3 conditions

in the early universe

Necessary ingredients

$\cancel{B}$ $\quad\quad$ $\cancel{CP}$ $\quad\quad$ *out of equilibrium*

Need of arrow of time

without suppression of inverse process,

$$\Delta B = \left(\Delta B\right)_{\rightarrow} + \left(\Delta B\right)_{\leftarrow} = 0$$

Out of equlibrium condition: case of heavy particle decay

$$X \to \overline{q}\overline{q}, ql$$

- One way decay, no inverse decay

$$H > \Gamma(= \alpha m_X) \qquad @T = m_X$$

$$H = \frac{1.6\sqrt{N}T^2}{m_{pl}}$$

Otherwise, Boltzmann suppression by $n_X \propto \exp(-m_X/T)$

Typically leading to

$$m_X > O[0.01]\alpha m_{pl} \approx 10^{15}GeV$$

Need for high unification scale

Reheating after inflation $\qquad T_{RH} > m_X$

Delicacy of CP: Quantum interference

Baryon excess from a pair of particle and antiparticle process, e.g. $X \quad \overline{X}$

$$\left| g_1 f_1 + g_2 f_2 + \cdots \right|^2 - \left| g_1^* f_1 + g_2^* f_2 + \cdots \right|^2$$

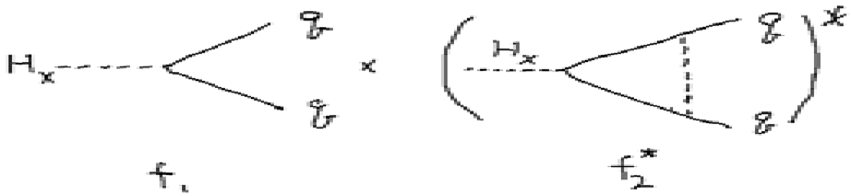
$$= -4\, \mathrm{Im}(g_1 g_2^*)\, \mathrm{Im}(f_1 f_2^*) + \cdots$$

$\mathrm{Im}(g_1 g_2^*) \neq 0$ CP violation

$\mathrm{Im}(f_1 f_2^*) \neq 0$ Rescattering phase

Interference computed by Landau-Cutkovsky rule

Σ_q



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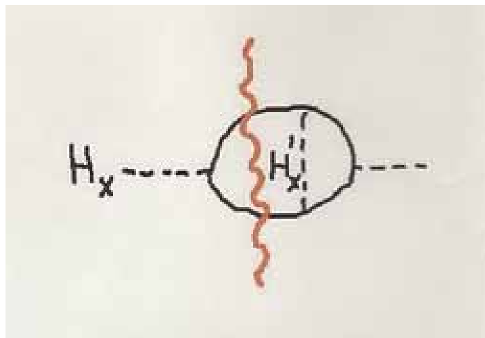


Dependence on dynamics

- 3 factors $$\frac{n_B}{s} = a \frac{n_X}{n_{th}} \varepsilon \qquad \varepsilon = \frac{n_{i \to f} - n_{\bar{i} - \bar{f}}}{n_{i \to f} + n_{\bar{i} - \bar{f}}}$$

- Dilution factor in late stages: $a = O[10^{-3}]\eta^{-1.2}$

  $O[10^{-3}]$ purely kinematical, determined by particle content

- When the out-of-equilibrium condition is partially satisfied, Result after integrating Boltzmann equations gives η-dependence. $$\eta \equiv \frac{\Gamma_X}{H_{T=m_X}}$$

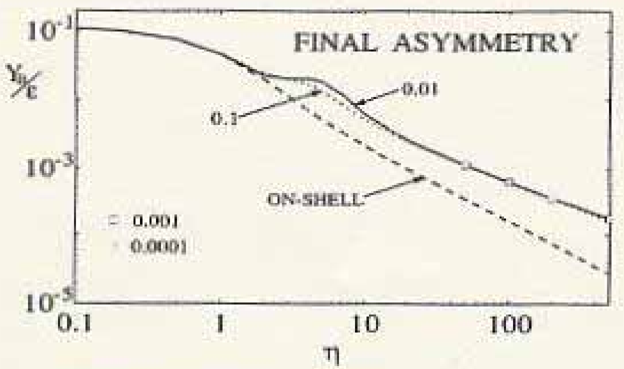


FIG. 8. Final amount of the baryon asymmetry plotted against η (= decay rate/Hubble rate at $T=M$). For comparison the result based on the on-shell contribution alone is shown by the dashed line. Those marked by open boxes and circles are results for smaller decay rates Γ/M.

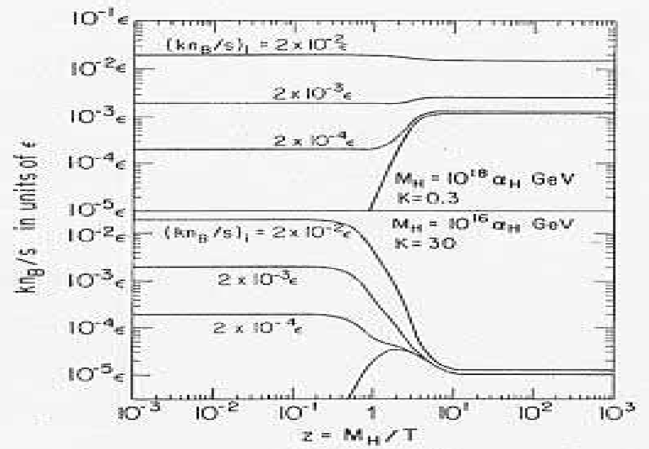


FIG. 6.(a) The time evolution ($z \sim t^{1/2}$) of the baryon asymmetry for $K=0.3$ (for $\alpha_H = 7\times10^{-5}$, $M_H = 7\times10^{13}$ GeV) is shown in units of ϵ for initial asymmetries of $2\times10^{-2}\epsilon$, $2\times10^{-3}\epsilon$, $2\times10^{-4}\epsilon$, and zero. Little damping occurs, and the final asymmetry is essentially the initial value plus $2\times10^{-3}\epsilon$ generated by delayed free decays. (b) The time evolution of kn_B/s is shown in units of ϵ for $K=30$ (for $\alpha_H = 7\times10^{-5}$, $M_H = 7\times10^{11}$ GeV) for the same initial conditions as above. The final asymmetry is essentially independent of the initial value, and equal to the value generated for zero initial asymmetry.

In GUT view,

- We are here,

  because matter that makes up us is ultimately unstable !

But,

lifetime of proton typically

10^{30} years >> age of universe

Constraints on and problems with GUT scenario

- Survival from combined effect of low T B-L-violation (e.g. $\Delta L \neq 0, and \Delta B = 0$) and electroweak damping, giving a constraint on L-breaking scale or neutrino mass

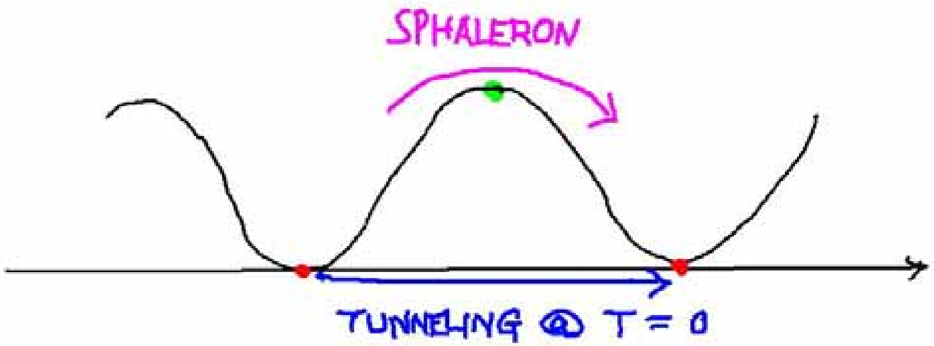
  e.g. Harvey-Turner

$$\langle m_\nu \rangle < \frac{4eV}{\sqrt{T_L / 10^{10} GeV}}$$

- Possible overproduction of gravitinos

  requiring a low reheat temperature after inflation

Electroweak baryon nonconservation



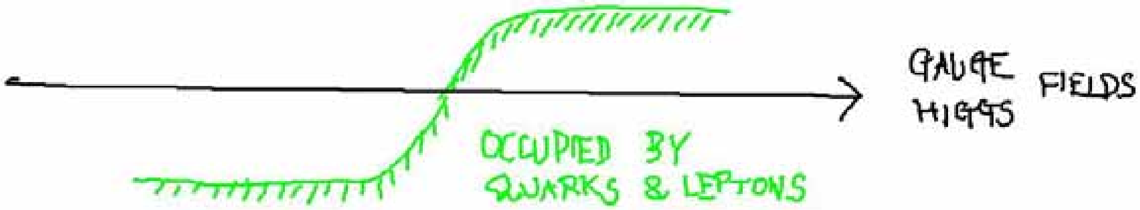
Gauge and Higgs

Electroweak baryon noncnoservation

suppressed at T=0 by e^{-137}

enhanced at finite T by barrier crossing

Can destroy preexisting B and L while keeping B-L



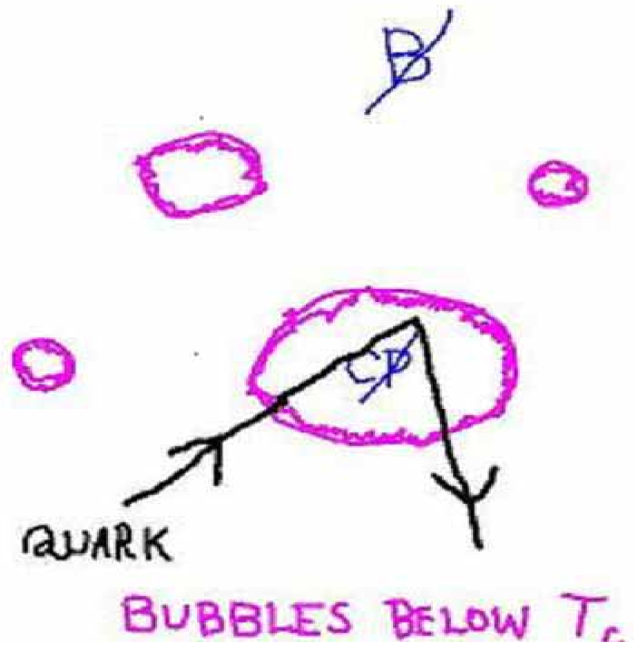
Mechanism due to level crossing of fermions caused by nontrivial gauge and higgs configuration of sphaleron and alike

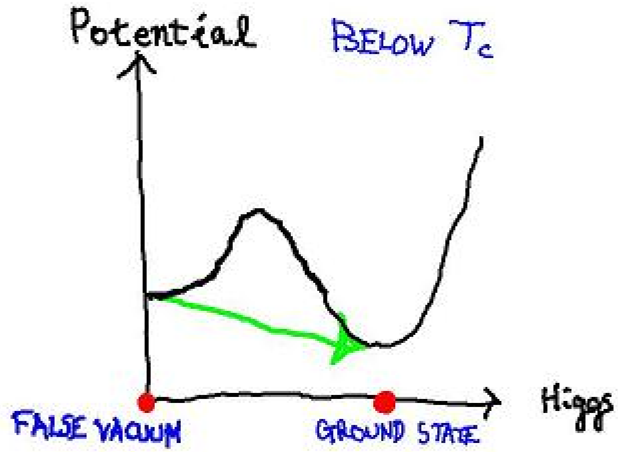
Baryogenesis in standard model

- $\not{B}$ unsuppressed $e^{-M_{sp}/T}$ at finite T

$$\gamma = o[1]\alpha_W{}^4 T \quad @T >> M_{sp} \approx O[TeV]$$

- $\not{CP}$ KM phase
- Out of equilibrium: 1st order phase transition via bubble formation





Difficulties of EW B-genesis

- No strong 1st order phase transition due to experimental Higgs mass bound

- Magnitude too small $$\frac{n_B}{n_\gamma} = o[10^{-21} - 10^{-25}]$$

Electroweak redistribution of B and L

$$B = a \cdot \Delta(B-L), \qquad a = \frac{8n_g + 4n_H}{22n_g + 13n_H} = \frac{28}{79}$$

For standard model of 3 generations

Damping effective @ $200 GeV < T < 10^{12} GeV$

e.g. Luty

B-L conserved and never washed out.

L genesis and B conversion

- L-genesis of amount ΔL first and electroweak conversion into B, via

$$B = -\frac{28}{79}\Delta L$$

For standard model of 3 generations

Interesting in view of possible connection to observed neutrino masses

Thermal L genesis

Fukugita-Yanagida

- Minimal extention of standard model with seesaw

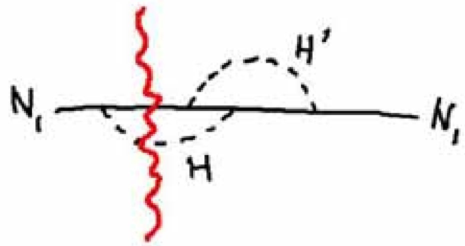
Right-handed Majorana decay $N_R \to lH, \overline{l}\overline{H}$

CP asymmetry with neutrino mass matrix $m_\nu = m_D M_N^{-1} m_D^T$

$$\varepsilon_1 = \frac{3}{16\pi}\frac{M_1}{v^2}\frac{\mathrm{Im}(m_D^\dagger m_\nu m_D^*)_{11}}{(m_D^\dagger m_D)_{11}} = O\left[\frac{M_1 \tilde{m}_\nu}{v^2}\delta\right]$$

Assuming mass hierarchy for 3 R-Majoranas N

$$\tilde{m}_\nu = \frac{(hh^+)_{11}}{M_1}v^2 = \frac{(m_D m_D^+)_{11}}{M_1}$$



δ = CP phase

Prejudices for simplification

- Completely general analysis meaningless due to many (18) parameters of m_ν, M_N matrices
- Constraints: known quantities

  $$\delta m^2{}_{23}, \delta m^2{}_{12}, \theta_{23}, \theta_{12}, \frac{n_B}{s}$$

- Some sort of mass hierarchy for heavy Majorana particles

  hierarchy for Dirac masses ?

- Symmetry

  GUT or flavor symmetry for Dirac term

  Effective parameters $\varepsilon_1, M_1, \tilde{m}_\nu$

  relation to theoretical models remote

Great impacts on neutrino masses and thermal history of universe

With hierarchy of masses, dependence on 3 parameters   Giudice et al

$$\varepsilon_1 , M_1 , \quad \tilde{m}_\nu$$

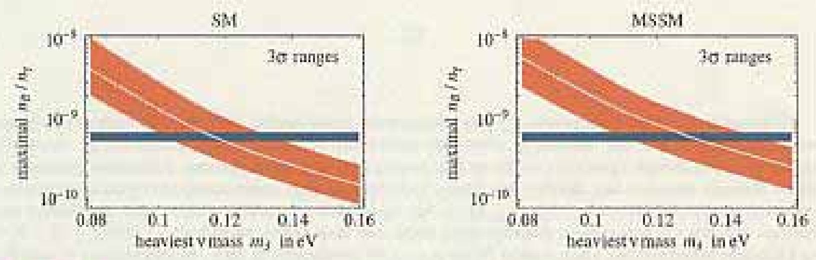


Figure 10: Leptogenesis bound on neutrino masses. *The plot shows the measur baryon asymmetry (horizontal line) compared with the maximal leptogenesis value as fun tion of the heaviest neutrino mass m_3, renormalized at low energy. Error bars are at 3σ*

- Connection to neutrino masses

$m_3 < 0.13 eV$ heaviest neutrino (WMAP,LSS 0.7eV)

$M_1 > 5 \cdot 10^8 GeV$ lightest R-neutrino

- Reheat temperature $$T_{RH} > M_1$$

Gravitino problem: a possible nightmare both for GUT B- and L-genesis

- Superpartner of graviton

mass $$m_{3/2} = O[TeV]$$

lifetime $$\Gamma = O[\frac{m_{3/2}{}^3}{m_{pl}{}^2}] = O[(10^5 \sec)^{-1}(\frac{m_{3/2}}{TeV})^3]$$

- Usual estimate of gravitino abundance and constraint from nucleosynthesis

$$\frac{n_{3/2}}{s} = O[10^{-2}]\frac{T_{RH}}{m_{pl}}$$

$$T_{RH} < 10^8 - 10^{10} GeV$$

Possible to produce GUT H_X **?**

A possible resolution, using preheating after inflation

- Important new element for particle production and B-genesis after inflation

Non-perturbative effect of parametric resonance, leading to

Complicated high energy phase of reheating, i.e. preheating

including dilution of gravitino abundance

Common to copious non-thermal production of R-Majorana neutrino for L-genesis and GUT Higgs

宇宙の熱史のはじまり

- 宇宙初期と起源論

  軽元素 $0.1 MeV$

  バリオン物質 $10^{15} GeV$ $\frac{n_B}{n_\gamma} \approx 10^{-10}$

  暗黒物質

  時空、構造形成の種 $10^{19} GeV$ $\frac{\delta\rho}{\rho} \approx 10^{-5}$

- 宇宙のエントロピー $n_\gamma \approx 400 cm^{-3}$

  からっぽの宇宙から 熱い宇宙へ

  インフレーション後の再加熱

起源論への答え

- 時空間 — インフレーション
- エントロピー — 周期振動するスカラー場からの粒子生成
- バリオン — バリオン数非保存とCPの破れ
- 構造形成の種 — インフラトン場の量子揺らぎ

暗黒物質　？
宇宙項　　？

Theory of particle production with chaotic potential

- Inflaton field oscillation given by

$\xi(t) = \xi_0 \cos(m_\xi t)$ (spatially homogeneous, periodic)

$\xi_0 >> m_{pl} \quad m_\xi \approx 10^{13} GeV$

Interaction by $g\xi\,\varphi^2$

Producing a pair of φ particles

For each momentum mode of massive particle

$$\ddot{\phi}_k + 3\frac{\dot{a}}{a}\dot{\phi}_k + (k^2 + m_\phi^{\ 2} + g\xi_0 \cos(m_\xi t))\phi_k = 0$$

$$h = \frac{k^2 + g\xi_0}{m_\xi^{\ 2}} \qquad \theta = \frac{g\xi_0}{m_\xi^{\ 2}} >> 1$$

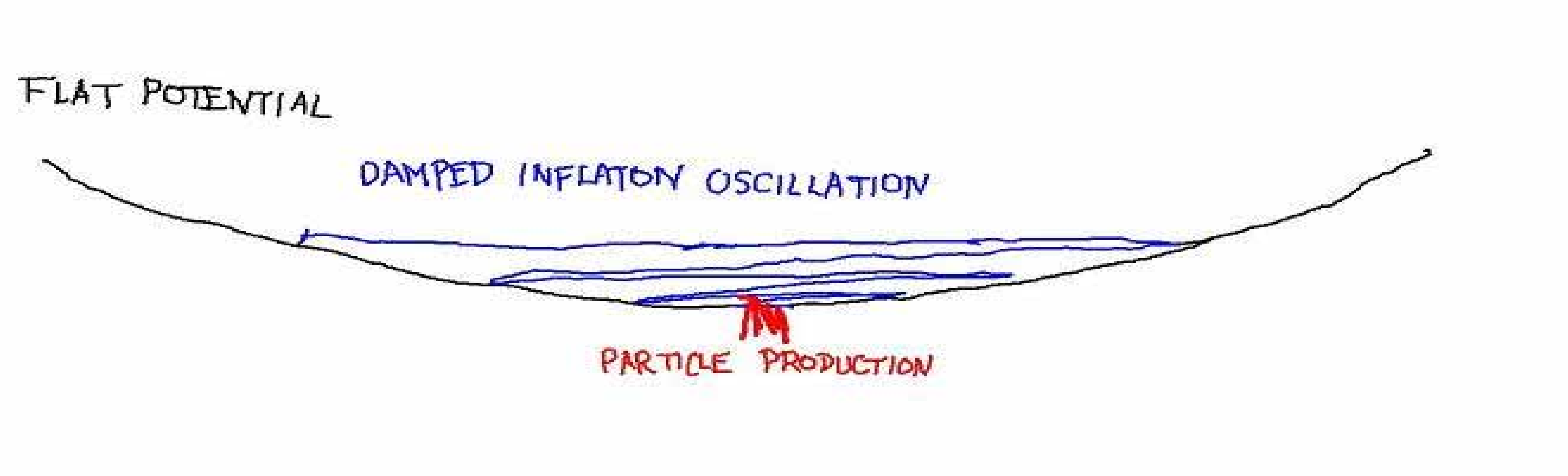
Model of inflation: Chaotic inflation

- Damped inflaton oscillation wth its mass

$$m_\xi = O\left[10^{13}\right] GeV$$

and initial dimensionless amplitude

$$\xi_0 = O\left[\frac{m_{pl}{}^2}{m_\xi}\right]$$



Theory of reheating

- Old view

Coherent inflaton oscillation = aggregate of 0-momentum particles

Independent particle decay $\xi \to \phi\phi$

$$\Downarrow$$

Instantaneous thermalization due to fast interaction

$$T_{RH}{}^4 \approx \rho_\xi$$

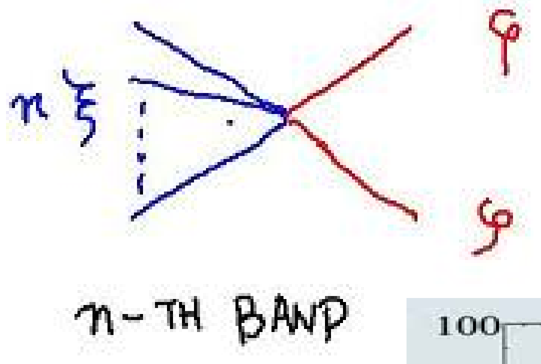
leading to reheat temperature $T_{RH} \approx \sqrt{\Gamma_\xi m_{pl}}$

with Γ_ξ Born decay rate

Non-perturbative effect of parametric resonance, producing large mass particles

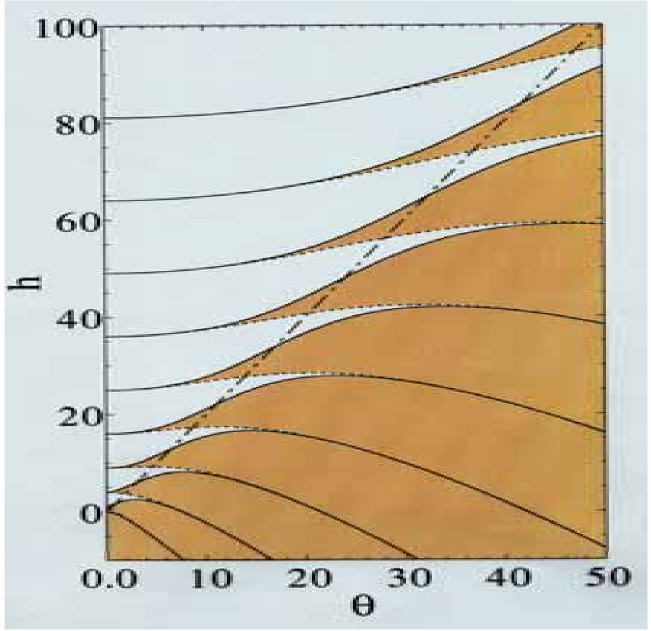
- n-th band contribution like

$$\xi^n \to \phi\phi$$



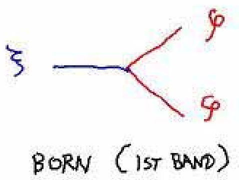
- Large mass production possible if with large n

$$\frac{m_\xi}{2} < m_\phi < \frac{n}{2} m_\xi$$



- Perturbative Born decay; from E-conservation

$$m_\phi < \frac{m_\xi}{2}$$

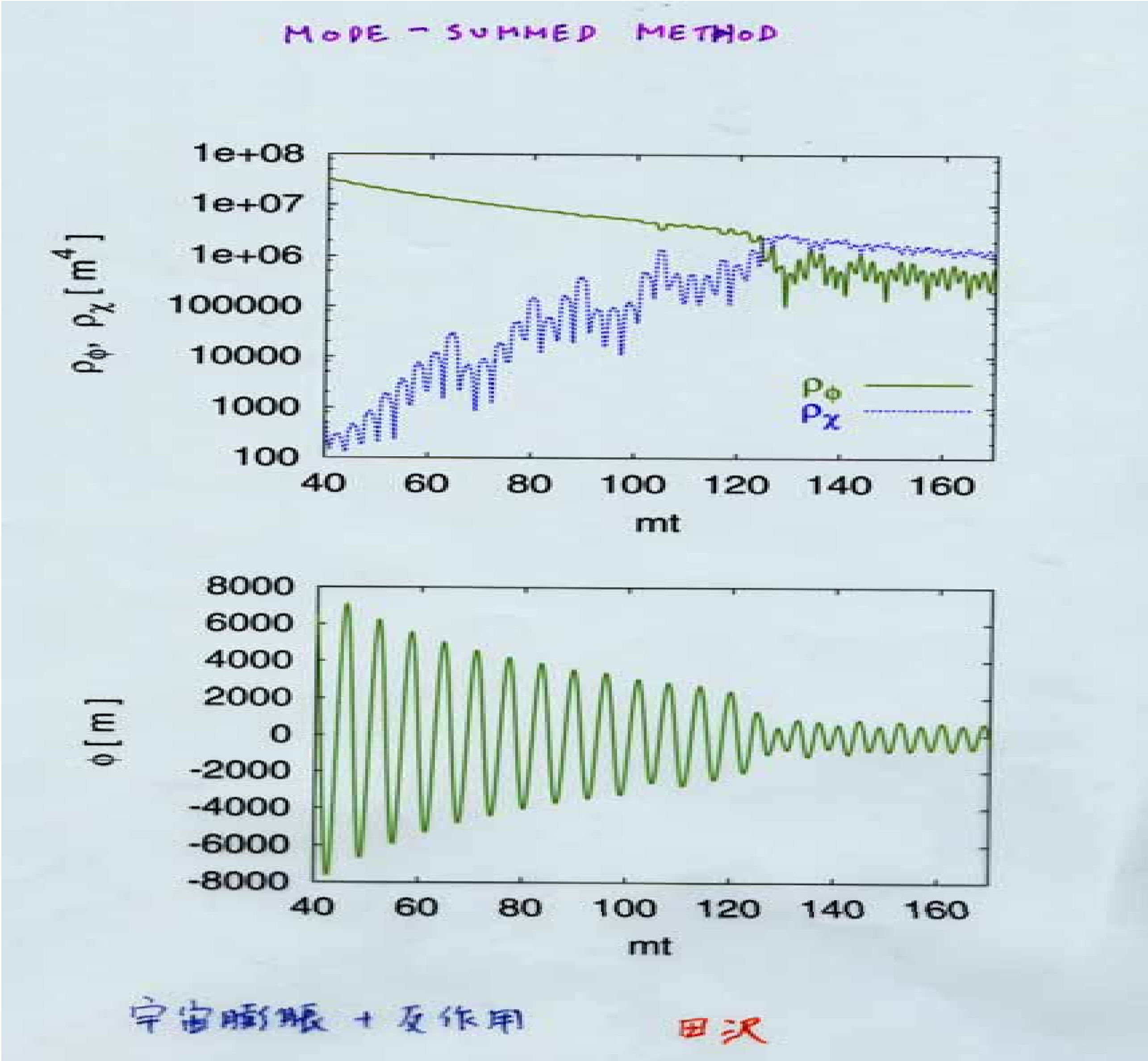


Problem of **parametric resonance**

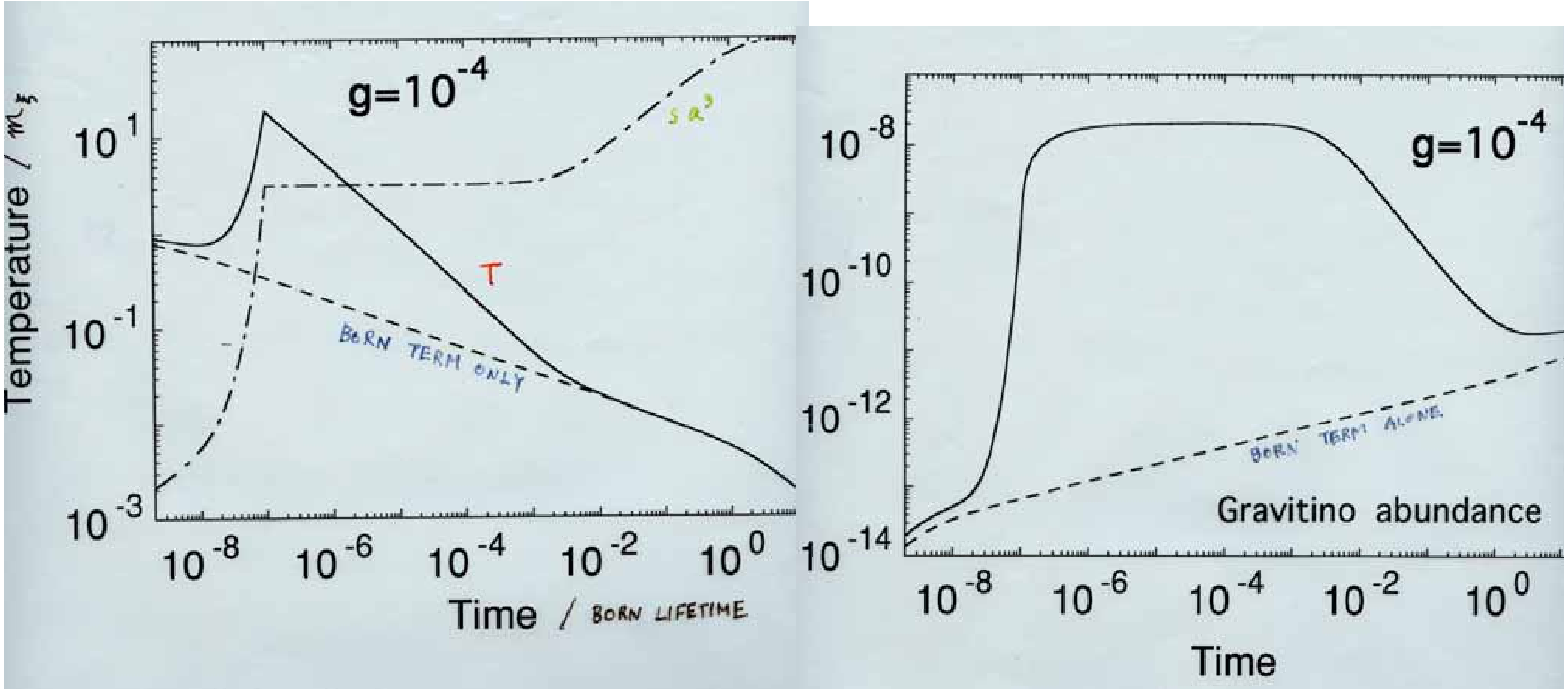
for large amplitude oscillation

How to swing： Need to vary center of your body periodically

Integration, with back-reaction and Einstein equation



Preheating stage and gravitino abundance



• e.g. B-generation during preheating and
gravitino abundance lowed by perturbative estimate is possible

New features : preheating

Violent process of particle production after O[10-100]oscillations

Initially highly non-thermal

Possibility of producing high mass GUT particles

Gravitino and B or L abundance to be comuted simultaneously, considering preheating

Conclusion on baryogenesis

- (B-L) genesis is a great hint on physics beyond the standard model, linking the micro and the macro worlds
- B-genesis still alive, waiting for nucleon decay
- L-genesis interesting due to its possible connection to the neutrino sector and lepton flavor violation
- Watch out gravitino overproduction
- Some new idea necessary for relation to low energy CP violation in K and B systems

カミオカンデの原点

企画課発行

15年 2月 21日（金） 日経 新聞（朝刊・夕刊）

担当： 坂井

私の履歴書

小柴（こしば）昌俊（まさとし） ⑳

ドイツ電子シンクロトロン研究所（DESY）での共同研究は一九七三年（昭和四十八年）八月から本格的な実験が始まった。須田英博君、山田作衛君（現高エネルギー加速器研究機構・素粒子原子核研究所長）、戸塚洋二君ら研究室のメンバーを送り込んだ。

陽子崩壊

その年の十一月に米スタンフォード大学と米国立ブルックヘブン研究所が電子・陽電子衝突で新粒子「J／プサイ」をほぼ同時に発見した。この発見は、それまで理論上で想定されてきた究極の素粒子「クォーク」に実在の確証を与えた。物理学の「十一月革命」と呼ばれる大事件だ。DESYチームは同じ電子・陽電子実験ですぐさまこの結果を検証する成果をあげた。

このころ日本には欧米に比肩できる加速器はなく、欧米の研究に参加して実績を積みたかった。しかし新素粒子探しに血眼の欧米勢は実績のない連中を仲間には入れない。ジレンマの中でDESYでの成果が突破口になった。

「大統一理論」巡り予測

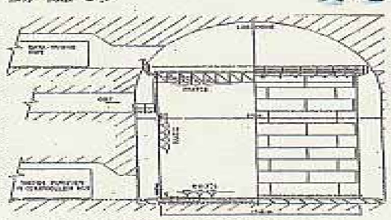
カミオカンデ実験へ動く

七八年にDESYで、東京大学は独ハイデルベルク大学、英マンチェスター大学と対等に組んで「JADE」と呼ぶ大型実験を始めた。日本のJA、ドイツのD、英国のEをとった愛称だ。折戸周治君（故人、元東大教授）を助教授にして、日本チームの指揮をとらせた。

二正面作戦のうち加速器実験が成果をあげる一方で、宇宙線研究にも転機が訪れた。七八年暮れ、当時の高エネルギー物理学研究所教授だった菅原寛孝君（現高エネルギー加速器研究機構長）から電話があった。「陽子崩壊の実験を考えてくれ」と言う。

自然界には四つの力がある。日常なじみの重力と電磁気力のほか、原子の内部で働く「強い力」と「弱い力」だ。この四つの力は宇宙の始まった時には一つの力だったが、時間とともに分かれたと考えられる。四つの力を統一的に説明する理論を築くことが物理学の大目標だ。重力を除く三つの力を説明する「大統一理論」は「非常にまれだが陽子が壊れることがある」と予測する。これを実験で確認できないかというのが菅原君の要望だ。

すぐに一つのアイデアが浮かんだ。宇宙線の届かない地下深くに大量の水をためてセンサーで監視する。水分子を構成する水素と酸素の原子核内で陽子崩壊が起きれば、それをキャッチできるかもしれない。こうしてカミオカンデ実験の構想が生まれた。



学会で発表したカミオカンデの概念図（1982年）

実はこのアイデアには根っこがあった。

話はシカゴ時代にさかのぼる。マルセル・シャイン教授の死後、気球実験の指揮官になった私がまずしなければならなかったことが、未使用の原子核乾板を宇宙線の来ない場所に保管することだった。そこで探したのがクリーブランドにある岩塩坑。ガイガーカウンターで測ったら宇宙線はほとんど来ない。

ここに無事保管した後、シカゴの自宅で顧問役のジュゼッペ・オッキャリーニ教授とビールを飲みながら話した。

「あの岩塩坑に穴を掘って水を張ったら、菌も繁殖しない澄み切った飽和食塩水のプールができる。そこに光電子増倍管（センサー）を設置したら何が見えるだろう」

この会話が頭の隅に残っていた。何が巡り巡って日の目を見るかわかったものではない。

（東京大学名誉教授）

1979年2月研究会

乞掲示

研究会の御案内

下記のような要領で研究会を開きます。

1) テーマ　宇宙のバリオン数と統一理論

2) 日時　2月13日　1:30 PM ～
　　　　2月14日　夕刻

3) 場所　高エネルギー物理学研究所

4) 別紙のようなプログラム（案）で講演を予定しています。

5) 興味をお持ちの方で参加を希望される方はその旨に宿泊の予定を添えて今月末日迄に下記へ申し込んで下さい。

〒300-32
茨城県筑波郡大穂町上原1-1
高エネルギー物理学研究所
荒船次郎

昭和54年1月25日

世話人代表
菅原寛孝

宇宙バリオン数と統一理論研究会プログラム（案）

13日

- 1:30　吉村　宇宙のバリオン数
- 2:30　Comments and Discussions
- 3:00　tea time
- 3:15　沢田 and/or 高岩 +Discussions　陽子の寿命の計算
- 4:00　~~井上~~ 中野　統一理論
- 5:00　藤川，柳田　comment

14日

- 9:30　木口義勝　クォークプラズマ
- 10:15　~~佐藤(勝)~~ 佐々木
- 11:15　佐藤(文)
- 12:00　昼休み
- 1:30　細谷　Hawking effect
- 2:15　木舟 and/or 渡辺　陽子の寿命の観測
- 3:00　tea time
- 3:15　Comments and Discussions

Proceedings of the Workshop on Unified Theories and Baryon Number in the Universe

National Laboratory for High Energy Physics (KEK)
February 13 - 14, 1979

$$m_\zeta \sim \frac{G_\phi^{\nu}\langle\phi\rangle}{G_\chi^{\nu}\langle\chi\rangle}\cdot G_\phi^{\nu}\langle\phi\rangle$$

Edited by
Osamu SAWADA and Akio SUGAMOTO

KEK, High Energy Accerelator Research Organization

Our life is finite.

- Should we explain everything now ?
- Richness of physics is fully explored ?
- What is interesting to me is useful to many others ?

宇宙環境での素過程

- 従来の手法

  on-shell S-行列の始状態、終状態についての平均化 (しばしば熱平衡状態)

  Boltzmann 方程式の熱平衡平均

- 不十分な点

  量子力学のoff-shell 効果が入らない

  粒子伝播への熱環境効果は？

応用

- 粒子崩壊

  X粒子, N粒子崩壊によるバリオン、レプトン生成

- 安定粒子の対消滅

  Dark matter 粒子の残存量

- トンネル効果

  宇宙相転移への環境効果

w. 松本

Off shell effects on B-asymmetry

Prolonged decay and CP asymmetry

I. Joichi, Sh. Matsumoto, and M. Yoshimura
Department of Physics, Tohoku University, Sendai 980-8578, Japan

(Received 2 March 1998; published 22 July 1998)

The time evolution of unstable particles that occur in the expanding universe is investigated. The off-shell effect not included in the Boltzmann-like equation is important for the decay process when the temperature becomes much below the mass of unstable particle. When the off-shell effect is taken into account, the thermal abundance of unstable particles at low temperatures has a power law behavior of temperature T, $(\Gamma/M)(T/M)^{\alpha+1}$ unlike the Boltzmann suppressed $e^{-M/T}$, with the power α related to the spectral rise near the threshold of the decay and with Γ the decay rate. Moreover, the relaxation time towards the thermal value is not governed by the exponential law; instead, it is the power law of time. The evolution equation for the occupation number and the number density of the unstable particle is derived, when both of these effects, along with the cosmic expansion, are included. We also critically examine how the scattering of thermal particles may affect the off-shell effect to the unstable particle. As an application showing the importance of the off-shell effect we compute the time evolution of the baryon asymmetry generated by the heavy X boson decay. It is shown that the out-of equilibrium kinematics previously discussed is changed; this change becomes considerable for large values of $\Gamma/H \gg 1$ where H is the Hubble rate at the temperature equal to the X-boson mass, while we confirm the previous result for small values of $\Gamma/H \ll 1$. [S0556-2821(98)02516-8]

PACS number(s): 98.80.Cq, 05.70.Ln, 11.30.Er

I. INTRODUCTION

There are many short-lived particles that have existed in abundance in the early universe whose temporary presence did not leave behind any measurable effect. Important exceptions to this exist, such as the neutron which certainly is the key for the explanation of the element abundance of the present universe.

A theoretical estimate of the abundance of these unstable particles after the cosmic temperature drops below the mass of the unstable particle is very important for subsequent time evolution. Most works in the past [1] were based on the Boltzmann equation that takes into account relevant reactions in the expanding universe. The use of the Boltzmann equation has however been questioned recently [2]; a more precise quantum mechanical description of the decay process in a thermal medium should contain important off-shell contributions not properly treated in the Boltzmann approach. These off-shell effects are eminent in the low temperature region. Low temperature effects are clearly important in this problem, since unstable particles are typically very nonrelativistic when they disappear in the early universe.

In the present work we shall develop a general formalism of computing the time evolution of the net number density of unstable particles and clarify the off-shell effect. The off-shell effect appears in two ways: first, in a slower relaxation towards the equilibrium abundance and second, in a larger equilibrium value not suppressed by the Boltzmann factor such as $e^{-\Delta M/T}$ where ΔM is the mass difference of the parent and the daughter particles. It is shown below that the off-shell effect becomes dominant below some temperature T_{eq}. The abundance of unstable particles then follows the power law: $n/T^3 \approx (\Gamma/M)(T/M)^{\alpha+1}$, where α is a parameter related to the threshold behavior of the spectral function for the decay and Γ is the decay rate. Thus, unstable particles do not disappear suddenly. Instead, their abundance gradually decreases with a power of decreasing temperature as the universe expands. Physical processes that follow after the decay are then prolonged. The off-shell effect turns out to be more prominent for a larger decay rate.

We next consider as an illustrative application of this general result the hypothetical X boson decay that may have created the matter-antimatter asymmetry when they decay [3,4]. We find that the time evolution of the baryon asymmetry is substantially changed and the severe lower bound of the X boson mass is considerably relaxed by the off-shell effect. For the first time we find that some mode of the X boson decay for baryogenesis is excluded due to the off-shell effect. This is the S-wave decay mode into a boson-pair.

This paper is organized as follows. In Sec. II the theoretical model of unstable particle decay is explained. This is a field theoretical extention of the harmonic model for the quantum dissipation in thermal medium discussed in [2]. We first present and formally solve the quantum mechanical model of the decay of excited levels in thermal medium. A great virtue of this model is that its integrability leads to explicit formulas for many quantities of interest. One can clearly see how the off-shell effect arises in these formulas. Extention to the unstable particle decay in field theory models can be made, but it is in general complicated and not readily solvable. But fortunately, in a thermal medium far away from the degeneracy limit which is relevant in the early universe the decay process is approximately described by this class of solvable quantum mechanical models extended to infinitely many decay channels. In Sec. III the occupation number and the number density of a species of unstable particles is calculated and its time evolution equation is derived in the expanding universe. The stationary abundance when the cosmic expansion is switched off is worked out, and its behavior at both high and low temperatures is studied in detail. In Sec. IV we pay special attention to the off-shell effect and its role in cosmology. We also discuss a possible

$$\frac{dY_+}{d\tau} = -\gamma(Y_+ - Y_0 - S_0), \tag{5.28a}$$

$$\frac{dY_-}{d\tau} = -\gamma\left(Y_- - \frac{\bar{\delta}}{2}(Y_1 + S_1)Y_B\right), \tag{5.28b}$$

$$\frac{dY_B}{d\tau} = \gamma\left(\epsilon Y_+ + \delta Y_- - \epsilon(Y_0 + S_0) - \frac{\delta\bar{\delta}}{2}(Y_1 + S_1)Y_B\right) - c(Y_1 + S_1)Y_B, \tag{5.28c}$$

$$\gamma = \frac{e^{-\Gamma t} + (\alpha+2)B(\alpha)\left(\frac{\Gamma}{M}\right)^{\alpha+4}\left(\frac{T}{M}\right)^{\alpha}(\Gamma t)^{-\alpha-3}}{e^{-\Gamma t} + B(\alpha)\left(\frac{\Gamma}{M}\right)^{\alpha+4}\left(\frac{T}{M}\right)^{\alpha}(\Gamma t)^{-\alpha-2}},$$

$$B(\alpha) = \frac{\Gamma(\frac{3}{2}\alpha + 3)}{\Gamma(\frac{1}{2}\alpha + 3)}, \tag{5.29}$$

$$S_0 = \frac{\zeta(\alpha+4)\Gamma(\alpha+4)\Gamma\left(\frac{\alpha}{2}+1\right)}{16\pi^2\sqrt{\pi}\,\Gamma\left(\frac{\alpha}{2}+\frac{5}{2}\right)}\frac{\Gamma}{M}\left(\frac{T}{M}\right)^{\alpha+1}, \tag{5.30}$$

$$c = \frac{3\left(\sum_j \alpha_j^2(\gamma_j + \bar{\gamma}_j) - \frac{\delta^2}{2}\right)}{\sum_j \alpha_j^2} > 0, \tag{5.31}$$

$$S_1 = S_1(\alpha) = \frac{\zeta(\alpha+3)}{8\pi^2\sqrt{\pi}}\frac{\Gamma\left(\frac{\alpha}{2}+1\right)\Gamma(\alpha+3)}{\Gamma\left(\frac{\alpha}{2}+\frac{3}{2}\right)}\frac{\Gamma}{M}\left(\frac{T}{M}\right)^{\alpha+1}. \tag{5.32}$$

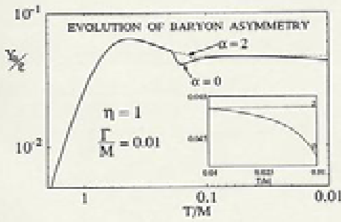


FIG. 6. Comparison of the time evolving baryon asymmetry. The case of $\alpha=0$ shown by the solid line and enlarged in the inset gives the vanishing value for the final asymmetry, unlike the $\alpha=2$ case shown by the dotted line.

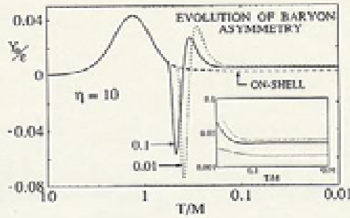


FIG. 7. Time evolution of the baryon asymmetry. Two cases of different decay rates, $\Gamma/M = 0.1, 0.01$, are compared to the evolution given by the on-shell contribution alone (the broken line). In the inset detailed behaviors are stressed.

$$S_2 = S_1(2\alpha) = \frac{\zeta(2\alpha+3)}{8\pi^2\sqrt{\pi}}\frac{\Gamma(\alpha+1)\Gamma(2\alpha+3)}{\Gamma(\alpha+\frac{3}{2})}\frac{\Gamma}{M}\left(\frac{T}{M}\right)^{2\alpha+1}, \tag{5.33}$$

$$Y_1 = \frac{1}{2\pi^2T^2}\int_0^\infty dk\,\frac{2k^2+M^2}{\sqrt{k^2+M^2}}\frac{1}{e^{\sqrt{k^2+M^2}/T}-1}. \tag{5.34}$$

The low temperature approximation was not assumed here for $f^{\mathrm{th}}(\omega)$, hence

$$Y_0 = \frac{1}{2\pi^2T^3}\int_0^\infty dk\,\frac{k^2}{e^{\sqrt{k^2+M^2}/T}-1}. \tag{5.35}$$

It can be readily proved by a rescaling argument that both Y_- and the baryon asymmetry Y_B is in direct proportion to the fundamental CP parameter ϵ. We assume that $\alpha \le 2$ as required for any renormalizable decay interaction.

Some results of numerical integration of the time evolution equation are presented in Fig. 6 and Fig. 7. The time evolution for $\alpha=0$ and $\alpha=2$ is evidently different, as seen in Fig. 6. Notably, the final Y_B vanishes for $\alpha=0$. This difference will be understandable analytically, as will be dis-

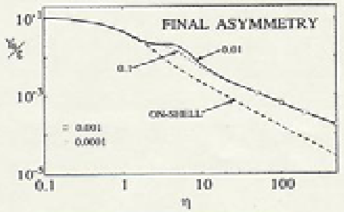
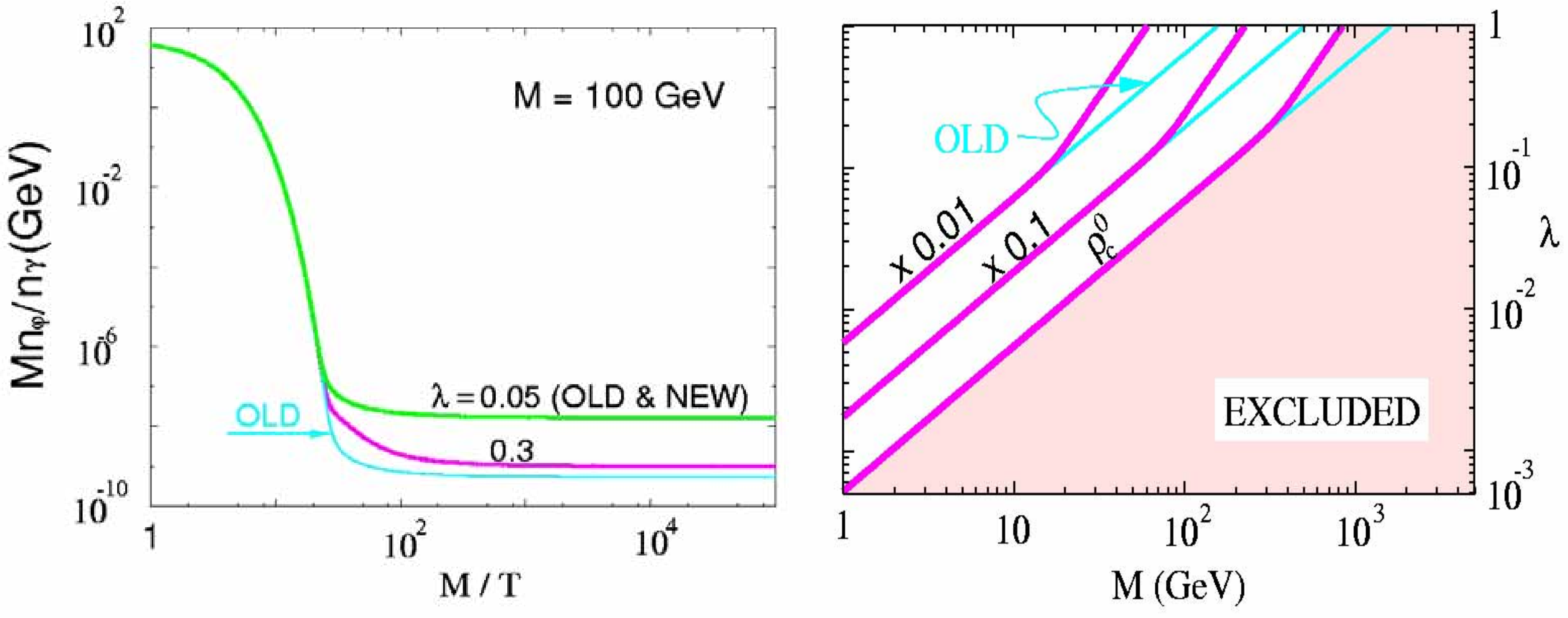


FIG. 8. Final amount of the baryon asymmetry plotted against η (= decay rate/Hubble rate at $T=M$). For comparison the result based on the on-shell contribution alone is shown by the dashed line. Those marked by open boxes and circles are results for smaller decay rates Γ/M.

Dark matter pair annihilation



Towards dynamical theory of 1st order phase transition

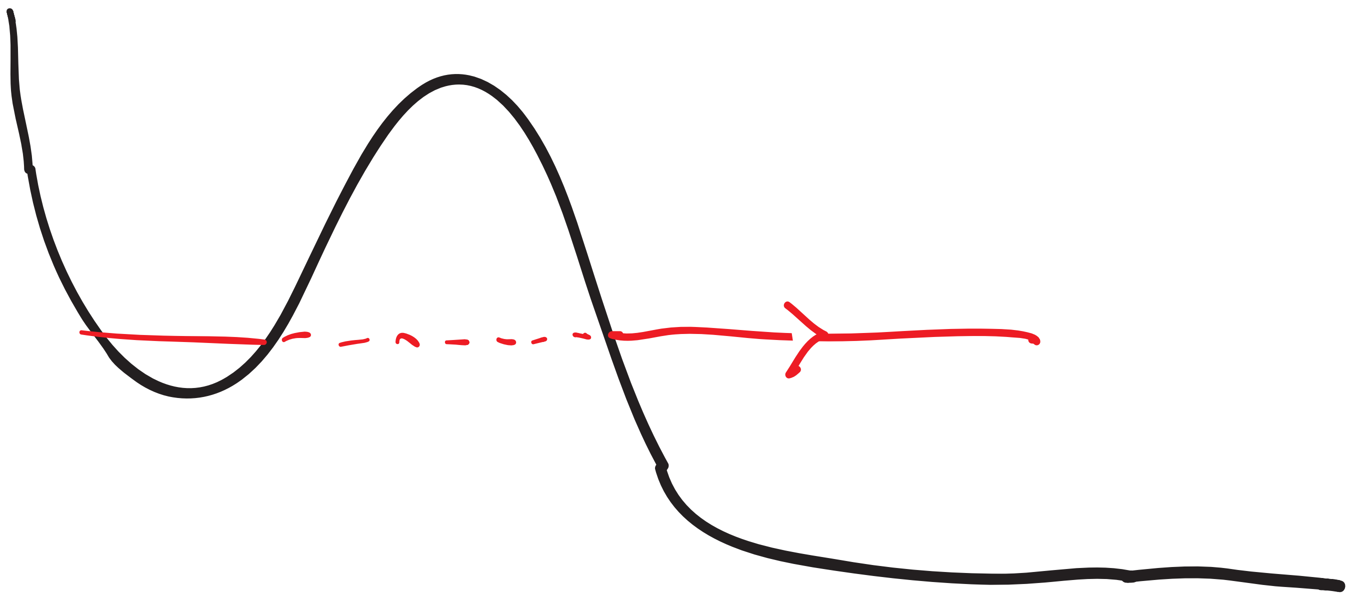
Real-time description of quantum tunneling

with effects of cosmological environment

New time scale caused by resonant enhanced tunneling

Semiclassical plus quantum picture

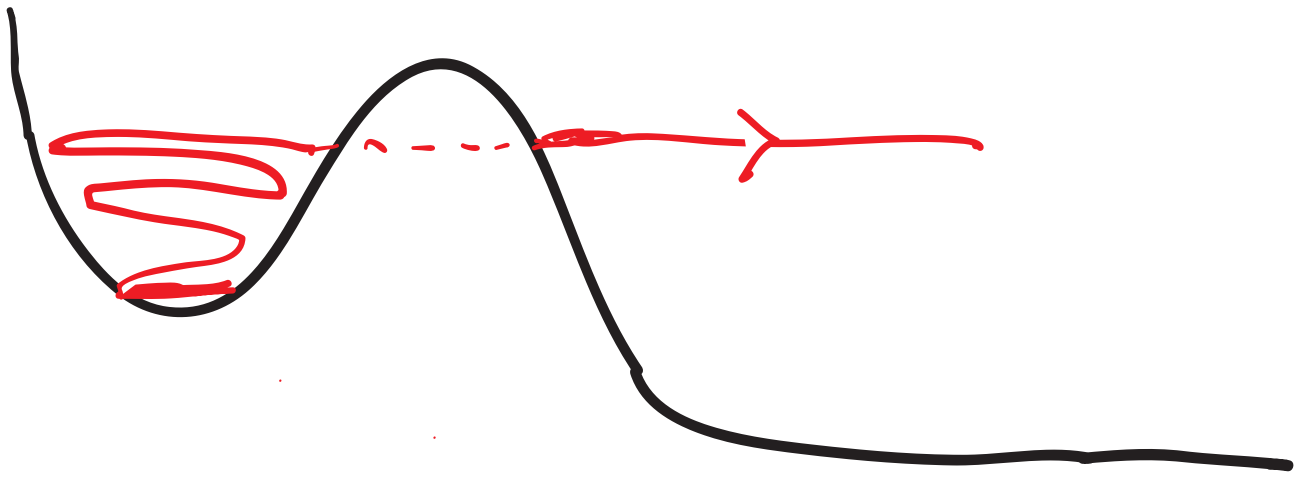
- Semiclassical approximation for metastable potential well, combined with quantum penetration formula beyond the barrier

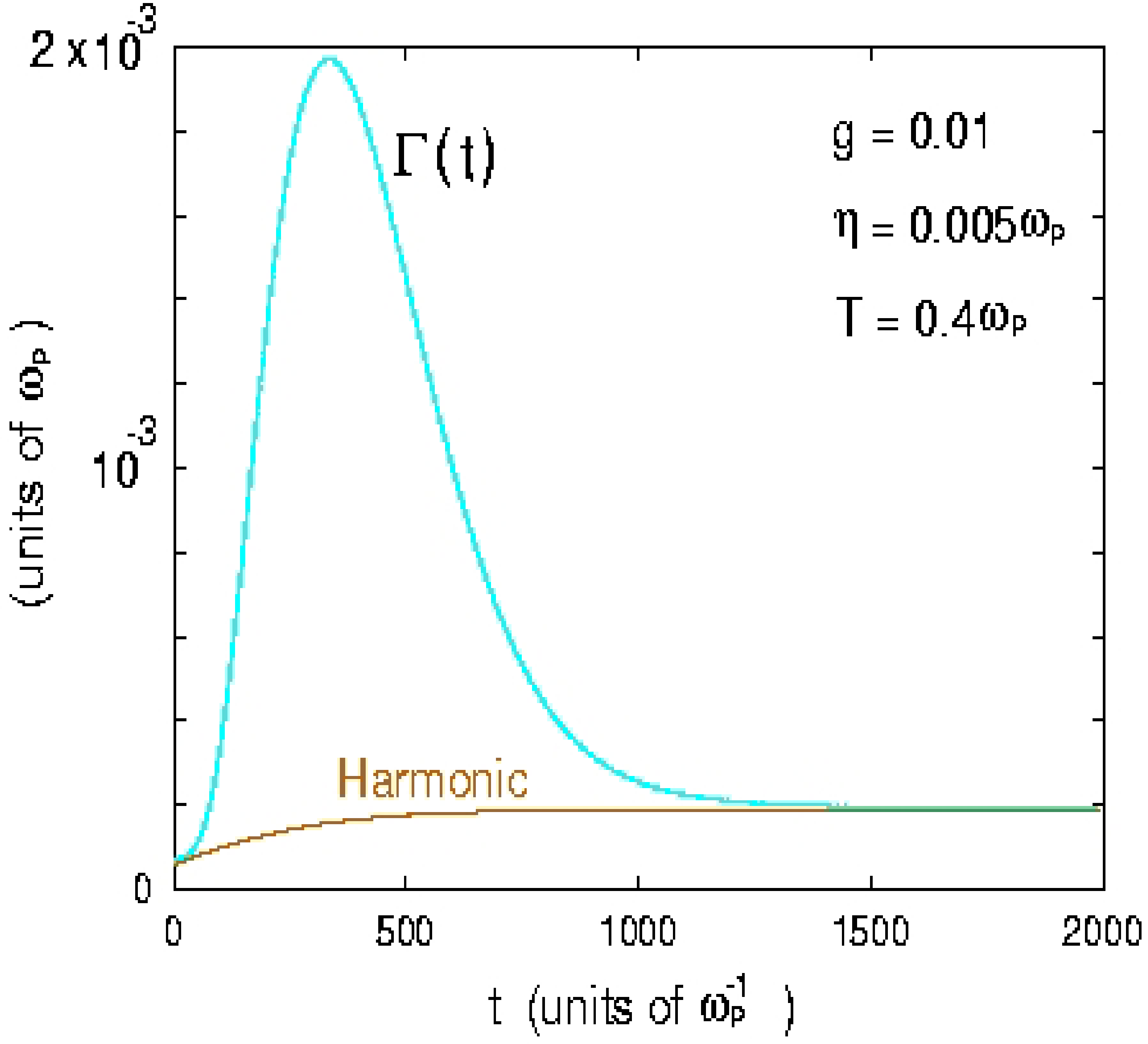


Resonant enhanced tunneling

During semiclassical motion,

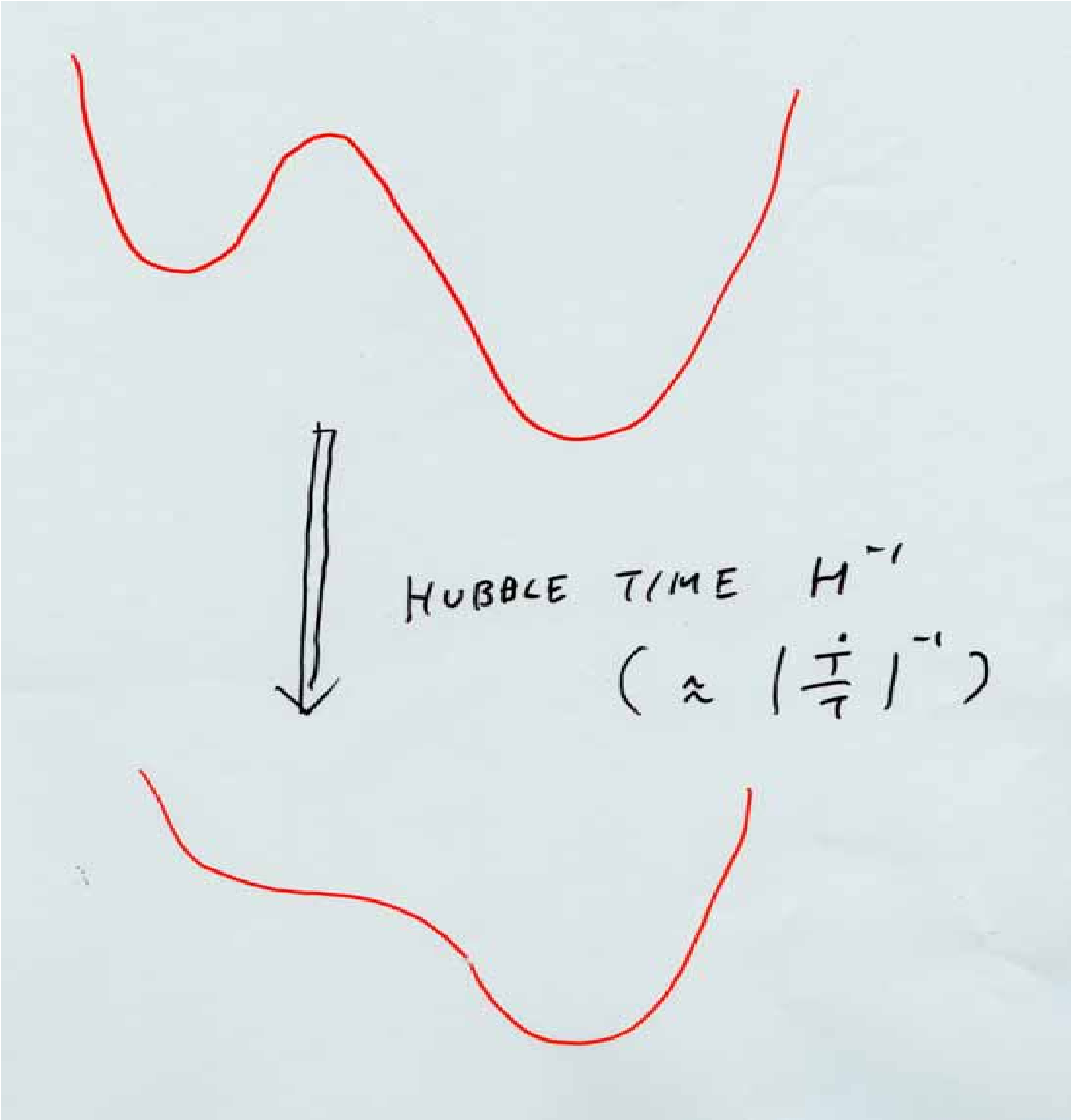
jump to higher levels by environment interaction



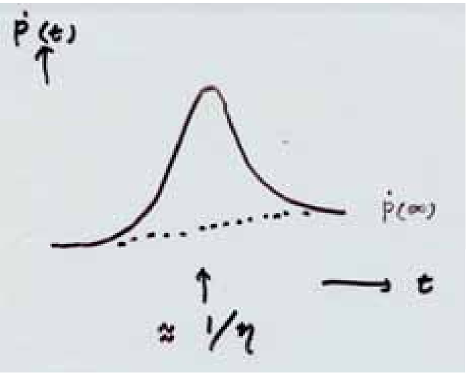


2 x10^-3
10^-3
0
(units of ωp)
Γ(t)
g = 0.01
η = 0.005ωp
T = 0.4ωp
Harmonic
0
500
1000
1500
2000
t (units of ωp^-1)

Usual picture of potential change



Time scale of resonant enhanced tunneling



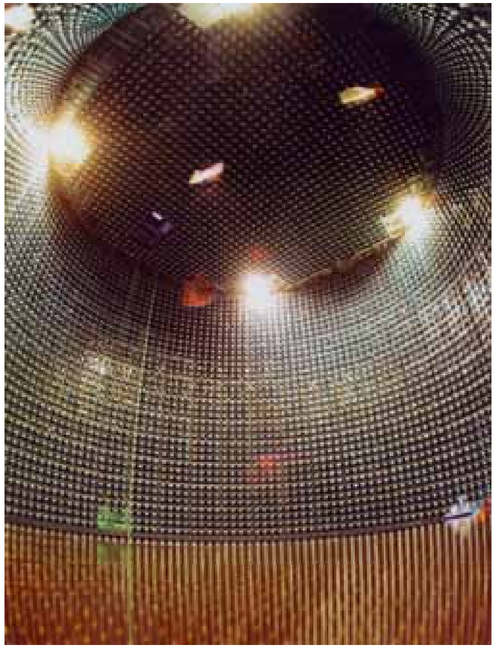
$$\approx 1 / \text{friction from enviroment}$$

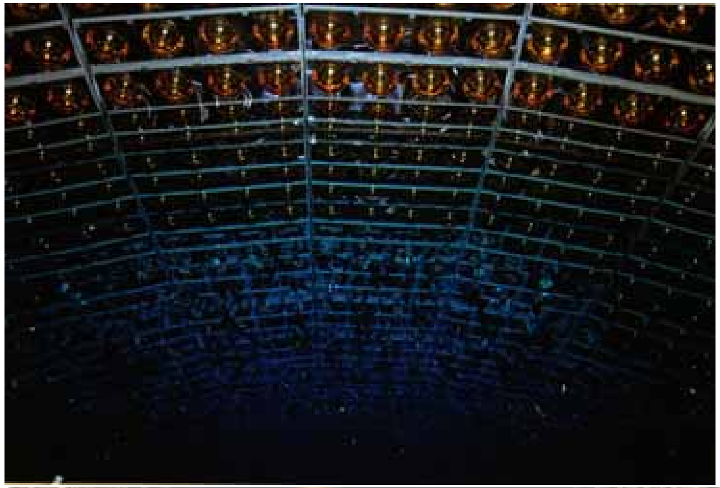
much shorter 1/Hubble time

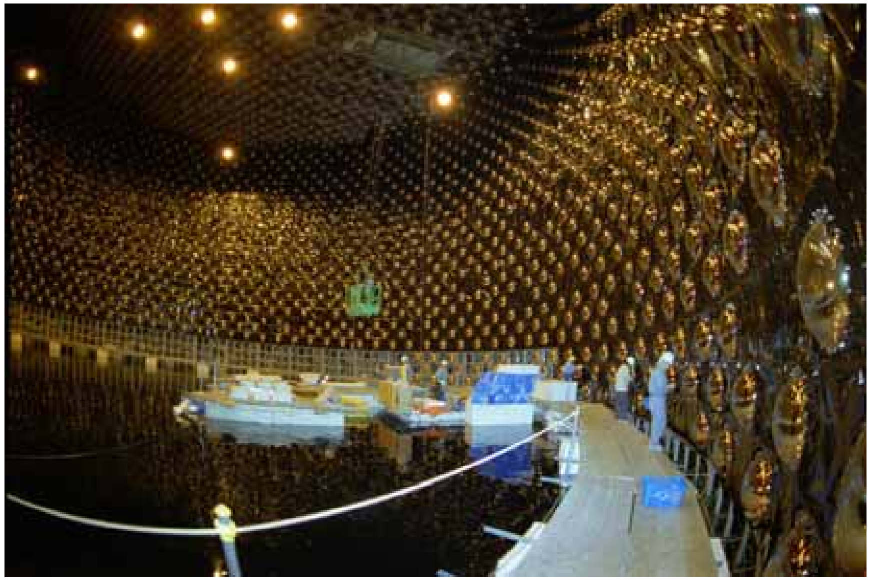
Thus,

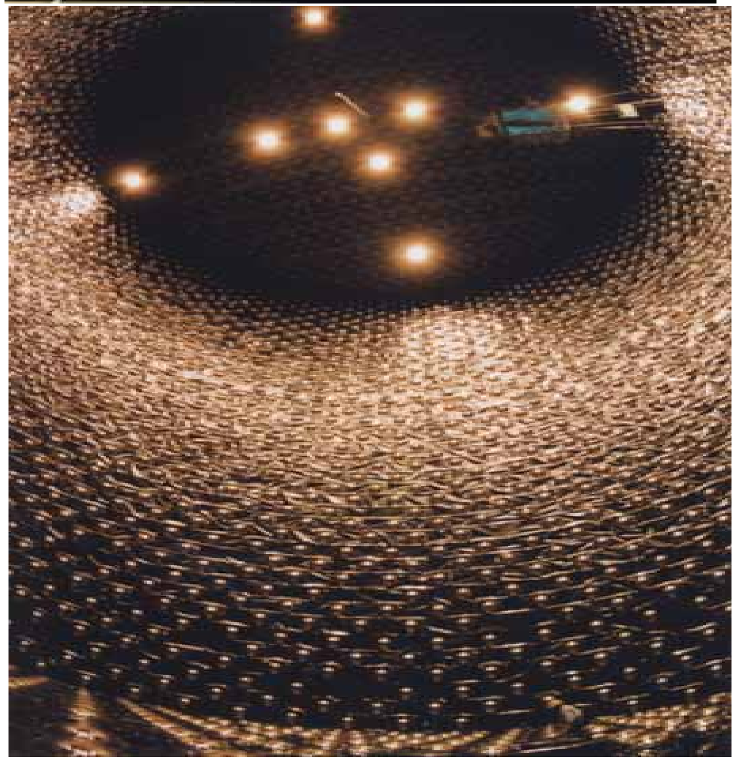
possibility of changing picture of 1st order PT

SK 復旧









近況

目標と自己評価

- 目標

時代を先取りするアイディアを出したい。
“できる限り”現象、実験との関連を追及し，
検証可能にしたい。

- 評価

成功したか、疑わしい。　まだ将来への課
題を残した。（これからも仕事がしたい）

好きな言葉
ウィトゲンシュタインとアインシュタイン

四・一一二　哲学の目的は思想の論理的解明である。
哲学は教説ではなく、活動である。
哲学の仕事は本質的に註解から成る。
哲学の成果は「哲学的な諸命題」ではなく、諸命題の明確化である。
哲学は、そのままではいわば濁っていて輪郭のはっきりしない諸思想を清澄にし、はっきりと限界づけなくてはならない。

（「哲学」なる語は、自然諸科学の上とか下とかに位する何ごとかを意味しているのでなくてはならず、自然諸科学と並ぶものではない。）

四・一一四　哲学は思考可能なものごとの限界を定め、そのことによって思考不可能なものごとの限界を定めるべきである。
哲学は思考不可能なものごとを思考可能なものごとによって内側から限界づけなくてはならない。

四・一一六　およそ考えうるものごとは、すべて明晰に考えうる。いい表わしうるものごとは、すべて明晰にいい表わしうる。

四・一二　命題は全現実を描出することができるが、それを描出しうるために現実と共有していなくてはならないもの――論理形式――を描出することはできない。
論理形式を描出しうるためには、われわれは当の命題と共に論理の外側に立つことができるのでなくてはならない。すなわち、世界の外側に。

四・一二七一　変項がそれぞれ形式的諸概念の符号になる。
なぜなら、いかなる変項も、その値すべてが所有する一定の形式を描出しており、その

ウィトゲンシュタイン「論理哲学論考」

'Science without religion is lame, religion without science is blind.' So Einstein once wrote to explain his personal creed: 'A religious person is devout in the sense that he has no doubt of the significance of those super-personal objects and goals which neither require nor are capable of rational foundation.' His was not a life of prayer and worship. Yet he lived by a deep faith—a faith not capable of rational foundation—that there are laws of Nature to be discovered. His lifelong pursuit was to discover them. His realism and his optimism are illuminated by his remark: 'Subtle is the Lord, but malicious He is not' ('Raffiniert ist der Herrgott aber boshaft ist er nicht.'). When asked by a colleague what he meant by that, he replied: 'Nature hides her secret because of her essential loftiness, but not by means of ruse' ('Die Natur verbirgt ihr Geheimnis durch die Erhabenheit ihres Wesens, aber nicht durch List.').

A, PAIS " SUBTLE IS THE LORD ... "

新たな旅立ち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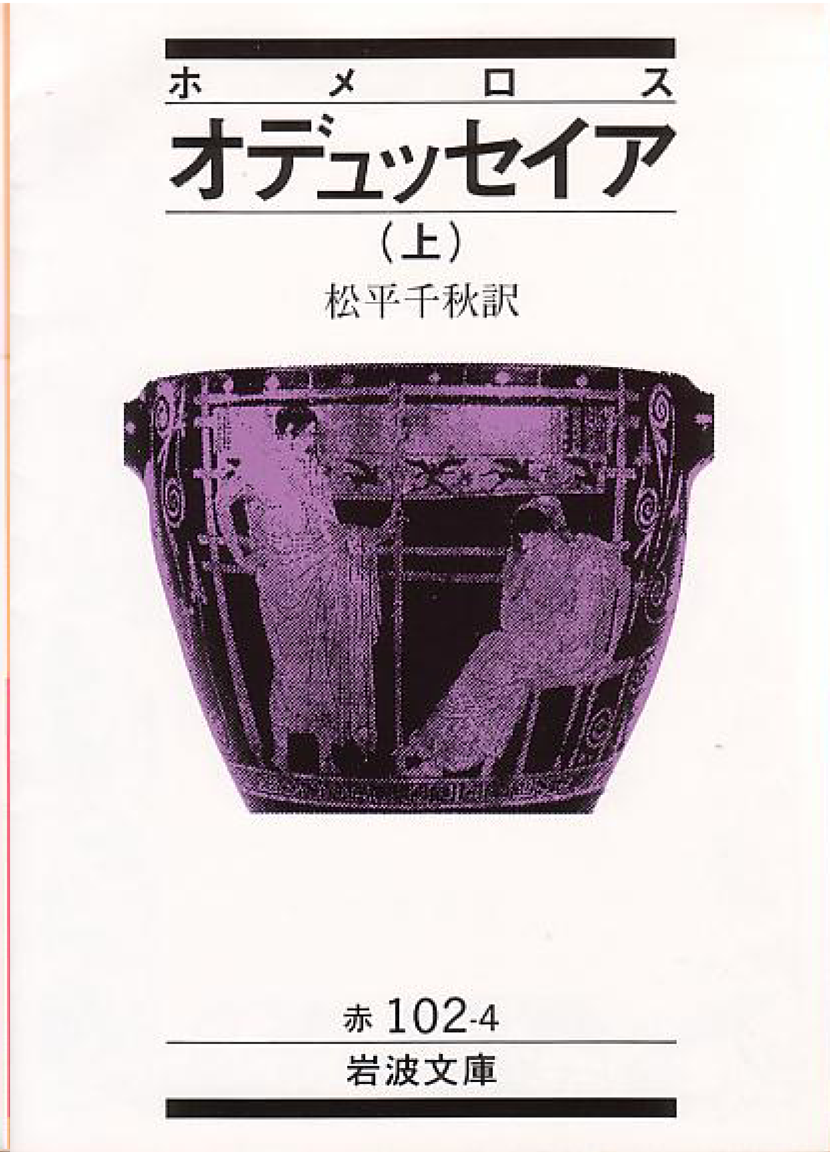
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