Preprint: hep-ph/9412364 and UVA-INPP-94-7 [25 Nov 1994]
Presented at Int'l. Conf. on Physics with GeV Beams, Juelich, Aug. 22-25, 1994.
Proceedings to be published by World Scientific.

STUDY OF CHIRAL INVARIANCE USING PIONS*

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ABSTRACT

We review two manifestations of chiral symmetry breaking in the pion sector. First, we examine the $\pi\pi$ scattering amplitude at threshold, which is completely determined by the chiral symmetry breaking part of the strong interaction, and, thus, directly constrains the form of the low energy effective lagrangians. We review the current status of the $\pi\pi$ scattering lengths in view of the recent results on reactions $\pi N \to \pi\pi N$ near threshold, and discuss the needed improvements. Second, we focus on the rate of the pion beta decay, determined at the tree level by the CVC hypothesis, a remnant of chiral symmetry breaking. Precise knowledge of the $\pi\beta$ decay rate constrains Cabibbo–Kobayashi–Masakawa unitarity, and, thus, can be used to distinguish among certain extensions of the minimal Standard Model. Present status of CKM unitarity, and a new precise measurement of the $\pi\beta$ decay rate at PSI are discussed.

1. Introduction: Chiral Symmetry and Pions

Chiral transformations are the most general form of rotations in fermion flavor space; for free massless fermions they leave the lagrangian invariant. Thus, the simplest realization of chiral symmetry is in a massless two-fermion system. Pure "isospin" rotations in the two-dimensional flavor space of fermions "u" and "d",

$$\psi' = \begin{pmatrix} u' \\ d' \end{pmatrix} = \exp(i\vec{\alpha}\cdot\vec{\tau}) \begin{pmatrix} u \\ d \end{pmatrix}, \tag{1}$$

leave the free massless lagrangian unchanged, forming an SU(2)×U(1) symmetry group. The generators, $\vec{\tau}$, are the Pauli spin–1/2 matrices (the U(1) generator is the 2×2 unit matrix). This, however, is not the maximal symmetry of the massless two-fermion system. Independent rotations in flavor space of projected left– and righthanded states,

$$\begin{aligned} \psi_L &= \tfrac{1}{2}(1-\gamma_5)\psi = \begin{pmatrix} u_L \\ d_L \end{pmatrix} \qquad \psi'_L = \exp(i\vec{\alpha}_L\cdot\vec{\tau})\psi_L, \\ \psi_R &= \tfrac{1}{2}(1+\gamma_5)\psi = \begin{pmatrix} u_R \\ d_R \end{pmatrix} \qquad \psi'_R = \exp(i\vec{\alpha}_R\cdot\vec{\tau})\psi_R, \end{aligned} \tag{2}$$

*Work supported by the United States National Science Foundation.

form the full chiral symmetry group of the lagrangian, $\mathrm{U(1)_L{\times}SU(2)_L{\times}SU(2)_R{\times}U(1)_R}$; γ_5 is the familiar Dirac matrix. Usually, attention is focused on the nonabelian $\mathrm{SU(N)_L{\times}SU(N)_R}$ subgroup. Strong interactions leave the chiral $\mathrm{SU(2)_L{\times}SU(2)_R}$ symmetry for the u and d quarks nearly intact, whereas the $\mathrm{SU(3)_L{\times}SU(3)_R}$ group, obtained by including the s quark, is violated considerably more. This is not surprising given that fermion mass terms, $-m\bar{\psi}\psi$, break chiral invariance of the lagrangian by mixing the left– and righthanded states, and that $m_s \gg m_{u,d}$.

For hadrons, chiral symmetry is broken in a very particular way: (i) explicitly, through small non-zero quark masses, and (ii) "spontaneously", through the non-invariance of vacuum to chiral transformations, but only to their axial-vector part, leading to the well known "partially conserved axial-vector current" (PCAC) relation. The effect of spontaneous chiral symmetry breaking (ChSB) dominates and hides the full symmetry from the casual observer; at the same time, it is the source of a rich phenomenology, requiring the existence of an octet of pseudoscalar Goldstone bosons, π's, K's, and η. Clearly, chiral symmetry plays a major role in the strong dynamics at low energies, one that has been much studied throughout the past four decades.[1]

Chiral symmetry has recently attained a new significance due to the failure of the full QCD to describe nonperturbative strong interaction phenomena below about 1 GeV. At these energies chiral symmetry provides the basic framework or the essential constraints for all low energy effective lagrangian models, e.g., the chiral perturbation theory (ChPT),[2] and recent QCD formulations of the Nambu–Jona-Lasinio (NJL) model in SU(2) and SU(3) flavor space.[3]

We examine two aspects of chiral invariance in the pion sector: (i) the mechanism of explicit ChSB, as evidenced in threshold $\pi\pi$ scattering, and (ii) repercussions of the spontaneous ChSB on weak-interaction physics, in particular the pion beta decay and its implications regarding the Standard Model.

2. Pion-Pion Scattering and the Explicit Breaking of Chiral Symmetry

It has been long established[4] that low energy $\pi\pi$ scattering provides a particularly sensitive tool in the study of the explicit breaking of chiral symmetry, since $a(\pi\pi)$, the $\pi\pi$ scattering lengths, vanish exactly in the chiral limit. To the extent that they differ from zero, $a(\pi\pi)$ provide a direct measure of the ChSB term in the pion sector, i.e., detailed knowledge of the $\pi\pi$ scattering lengths and phase shifts provides much needed input in fixing the parameters of ChPT and other effective models.

In the 1960's, before the advent of QCD, $\pi\pi$ scattering lengths were studied in order to distinguish between the various proposed dynamical mechanisms of ChSB in strong interactions. Weinberg's,[4] Schwinger's[5] and Chang and Gürsey's[6] effective lagrangians differ from each other in the form of the response of the pion field to chiral transformations, thus predicting different values for the s-wave $\pi\pi$ scattering lengths. (It was subsequently established that Weinberg's lagrangian has the leading-order ChSB term consistent with QCD.)

Weinberg's soft-pion prediction was improved in 1982 by Jacob and Scadron who introduced a correction due to the non-soft $S^* \to \pi\pi$ isobar background.[7] Soon thereafter, Gasser and Leutwyler applied chiral perturbation theory to the problem, and predicted the behavior of the $\pi\pi$ amplitude to order p^4.[8] Also inspired by QCD, Ivanov and Troitskaya used the model of dominance by quark loop anomalies (QLAD) to obtain $\pi\pi$ scattering lengths.[9] The Jülich group, however, constructed a meson exchange model of pseudoscalar-pseudoscalar meson scattering, and applied it to calculate a number of $\pi\pi$ and $K\pi$ scattering observables.[10] Calculations of $\pi\pi$ scattering lengths within the SU(2)$\times$SU(2) and SU(3)$\times$SU(3) realizations of the NJL model are found in Refs. [11] and [12], respectively. Kuramashi and coworkers successfully applied quenched lattice QCD on a $12^3\times 20$ lattice, and obtained $\pi\pi$ scattering length values in the neighborhood of the older current-algebra calculations.[13] Roberts *et al.* recently developed a model field theory, referred to as the global color-symmetry model (GCM), in which the interaction between quarks is mediated by dressed vector boson exchange.[14] Predictions of all these models are compared to experimental results below.

2.1. *Experimental Determination of $\pi\pi$ Scattering Lengths*

Experimental evaluation of $\pi\pi$ scattering observables is difficult, primarily because free pion targets are not available. Scattering lengths further require measurements close to the $\pi\pi$ threshold, where phase space strongly limits measurement rates. Several reactions have been studied or proposed as a means to obtain near-threshold $\pi\pi$ phase shifts: $\pi N \to \pi\pi N$, K_{e4} decays, $\pi^+\pi^-$ atoms, $e^+e^- \to \pi\pi$, etc. Only the first two reactions have so far proven useful in studying threshold $\pi\pi$ scattering. The three main methods are discussed below in the order of decreasing reliability.

2.1.1. *K_{e4} Decays*

The $K^+ \to \pi^+\pi^- e^+\nu$ decay (called K_{e4}) is in several respects the most suitable process for the study of near-threshold $\pi\pi$ interactions. The interaction takes place between two real pions on the mass shell, the only hadrons in the final state. The dipion invariant mass distribution in K_{e4} decay peaks close to the $\pi\pi$ threshold, and $l = I = 0$ and $l = I = 1$ are the only dipion quantum states contributing to the process. These factors, and the well understood $V - A$ nature of the weak decay, favor the K_{e4} process among all others in terms of theoretical uncertainties. On the other hand, measurements are impeded by the low branching ratio of the decay, $\sim 3.9 \times 10^{-5}$.

The most recent K_{e4} measurement was made by a Geneva–Saclay collaboration in the mid-1970's who detected some 30,000 K_{e4} decays.[15] Analysis of this low-background, high-statistics data illustrates well the difficulties encountered in extracting experimental $\pi\pi$ scattering lengths. By itself, the Geneva–Saclay experiment determines $a_{\ell=0}^{I=0}$ with a $\sim$35% uncertainty, and constrains the slope parameter b:[15]

$$a_0^0 = 0.31 \pm 0.11\ \mu^{-1}, \qquad b = b_0^0 - a_1^1 = 0.11 \pm 0.16\ \mu^{-1}, \tag{3}$$

reflecting the relative insensitivity of K_{e4} data to $a(\pi\pi)$. Pion-pion phase shifts are, however, constrained by unitarity, analyticity, crossing and Bose symmetry, which were extensively studied by Roy,[16] and Basdevant *et al.*,[17, 18] resulting in a set of relations known as the "Roy equations". By using the Roy equation constraints of Basdevant, *et al.*,[18] more accurate results for a_0^0 and b_0^0 were obtained,[15] as well as scattering lengths and slope parameters for $(l = 0,\ I = 2)$ and $(l = 1,\ I = 1)$.[19]

2.1.2. Peripheral $\pi N \to \pi\pi N$ Reactions: the Chew–Low Method

Particle production in peripheral collisions can be used to extract information on low-energy scattering of two of the particles in the final state, as shown by Chew and Low in 1959.[20] For the $\pi N \to \pi\pi N$ reaction, the Chew–Low formula,

$$\sigma_{\pi\pi}(m_{\pi\pi}) = \lim_{t\to\mu^2} \left[\frac{\partial^2 \sigma_{\pi\pi N}}{\partial t \partial m_{\pi\pi}} \cdot \frac{\pi}{\alpha f_\pi^2} \cdot \frac{p^2 (t-\mu^2)^2}{t m_{\pi\pi} k} \right], \tag{4}$$

relates $\sigma_{\pi\pi}(m_{\pi\pi})$, the cross section for pion-pion scattering, to double differential $\pi N \to \pi\pi N$ cross section and kinematical factors: p, momentum of the incident pion, $m_{\pi\pi}$, the dipion invariant mass, t, the Mandelstam square of the 4-momentum transfer to the nucleon, $k = (m_{\pi\pi}^2/4 - \mu^2)^{1/2}$, momentum of the secondary pion in the rest frame of the dipion, $f_\pi \approx 93$ MeV, the pion decay constant, and $\alpha = 1$ or 2, a statistical factor involving the pion and nucleon charge states. The method relies on an accurate extrapolation of the double differential cross section to the pion pole in order to isolate the one pion exchange (OPE) pole term contribution. Since the exchanged pion is off-shell in the physical region ($t < 0$), this method requires measurements under conditions which maximize the OPE contribution and minimize all background contributions, i.e., peripheral pion production with t as close to zero as possible, and incident momenta typically above ~ 3 GeV/c.

The Chew–Low method has been refined considerably over time, particularly by Baton and coworkers,[21] whose approach enables extraction of $\pi\pi$ phase shifts through appropriate treatment of the angular dependence of the $\pi N \to \pi\pi N$ exclusive cross sections. The "Roy equation"[16, 17, 18] dispersion constraints enable the use of accurate high energy peripheral pion production data, whereas such data are not available below $m_{\pi\pi} \approx 500$ MeV. The peripheral pion production data base has not essentially changed since the early 1970's, and is dominated by two experiments, performed by the Berkeley[22] and CERN-Munich[23] groups. These data were ultimately combined with the Geneva–Saclay K_{e4} data in a comprehensive dispersion-relation analysis,[19] as discussed in the preceding section.

There have been independent Chew–Low type analyses since the 1970's; the last published one, performed by the Kurchatov Institute group in 1982, was based on a set of some 35,000 events recorded in bubble chambers.[24] The same group has recently updated their analysis.[25]

2.1.3. $\pi N \to \pi\pi N$ Reactions near Threshold

Weinberg pointed out early on that the OPE graph plays an important role in the $\pi N \to \pi\pi N$ reactions near threshold.[4] Olsson and Turner's soft-pion model incorporates only the OPE and contact terms in the lagrangian at threshold,[26] leading to a straightforward relation between the $\pi\pi$ and $\pi N \to \pi\pi N$ threshold amplitudes with only one parameter, ξ. Thus, soft-pion theory has provided the main inspiration for the $\pi N \to \pi\pi N$ near-threshold measurements. In Olsson–Turner's model, it is sufficient to measure total $\pi N \to \pi\pi N$ cross sections, from which quasi-amplitudes can be calculated and extrapolated to threshold. Unfortunately, the Olsson–Turner parameter ξ is incompatible with QCD.

Unlike the two preceding methods, there has recently been a great deal of experimental activity on near-threshold $\pi N \to \pi\pi N$ measurements. There are 5 charge channels accessible to measurement, $\pi^- p \to \pi^-\pi^+ n$, $\pi^- p \to \pi^0\pi^0 n$, $\pi^- p \to \pi^-\pi^0 p$, $\pi^+ p \to \pi^+\pi^0 p$, and $\pi^+ p \to \pi^+\pi^+ n$. We label them for brevity with their final state charges as ($-$+n), ..., (++n), respectively. Recent experiments reporting total $\pi N \to \pi\pi N$ cross sections are summarized below, while data available before 1984 is reviewed in Ref. [27]

The OMICRON group at CERN has measured cross sections in the (+$-$n), ($-$0p), (++n) and (+0p) channels by detecting coincident charged particle pairs in a two-sided in-plane magnetic spectrometer.[28] Sevior and coworkers at TRIUMF measured inclusive cross sections for the reaction (++n) down to 5 MeV above threshold, using a novel technique involving an active plastic scintillator target and neutron detection.[29] J. Lowe and coworkers at Brookhaven measured the channel (00n)down to about 5 MeV above the threshold, using the Crystal Box detector, in a kinematically complete measurement.[30] Finally, a Virginia-Stanford-LAMPF team studied the (+0p) channel in an exclusive measurement using the LAMPF π^0 spectrometer and an array of plastic scintillation telescopes for π^+ and p detection.[31, 32] Several other experiments are currently under way at TRIUMF and LAMPF, aiming to bring the existing cross section data base even closer to the $\pi\pi$ threshold.

Apart from a discrepancy in the (++n) channel between the TRIUMF[29] and OMICRON[28] data, the entire body of data appears globally consistent within the quoted uncertainties.[34, 31, 32] We also note the high-statistics angular correlation measurements in the (+$-$n) channel performed by the Erlangen group at somewhat higher energies than the bulk of the other work in this Section.[33] Closer analysis of the exclusive cross sections is only now beginning, as high-statistics data were not available until recently. Therefore, the methods of analysis are still being refined. Recent work of interest has been reported by the Erlangen[35] and St. Petersburg[36] groups, who also discuss earlier work of other groups not included in this brief review.

2.2. Comparison of Experimental Results and Predictions for $a(\pi\pi)$

Figure 1. summarizes quantitatively the theoretical model predictions and the experimental determinations of the s-wave $\pi\pi$ scattering lengths.

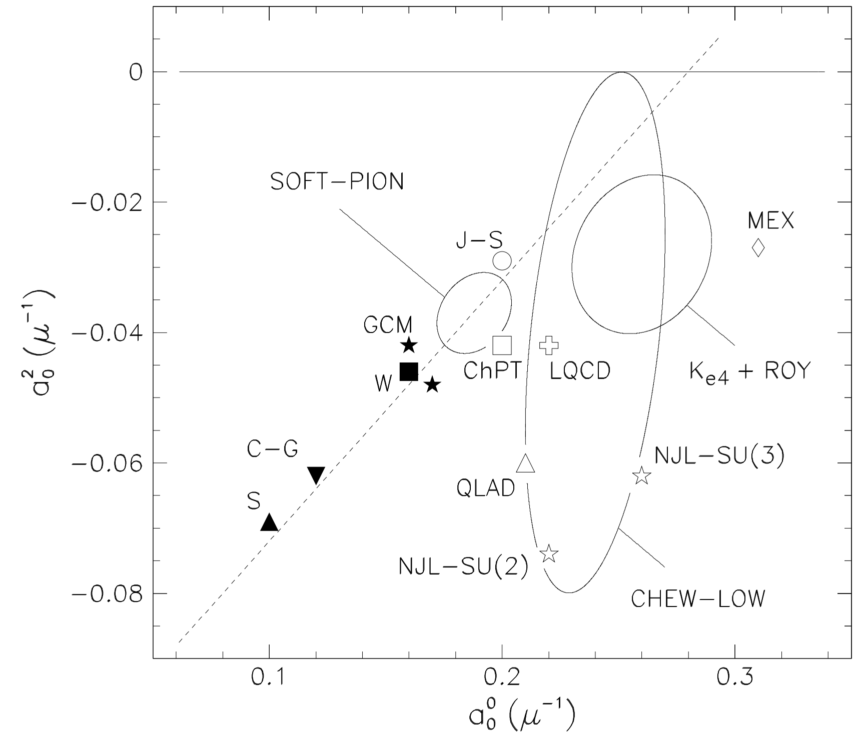

Figure 1: Summary of the $\pi\pi$ scattering length predictions (symbols) and experimental results (contour limits). Model calculations: Weinberg (full square), Schwinger (filled triangle), Chang and Gürsey (filled inverted triangle), Jacob and Scadron (open circle), Gasser and Leutwyler – ChPT (open square), Ivanov and Troitskaya – QLAD (open triangle), Lohse *et al.* – Meson Exchange (open rhomb), Ruivo *et al.* and Bernard *et al.* – NJL (open stars), Kuramashi *et al.* – quenched lattice gauge QCD (open cross), Roberts *et al.* – GCM (filled stars).

Considerable scatter of predicted values of a_0^0 and a_0^2 is evident in Fig. 1. Even disregarding the 1960's calculations of Schwinger[5] and Chang and Gürsey[6] later superseded by QCD, a significant range of predicted values remains. However, in view of the present experimental uncertainties (see Fig. 1.), most authors claim that their results are supported by the data. This is not a satisfactory state of affairs. Progress is required on two fronts in order to improve the present uncertainties in the experimentally derived $\pi\pi$ scattering lengths.

First, more accurate data on K_{e4} decays are needed, in order to reduce the current error limits in the analysis. However, K_{e4} data alone will not suffice because of their insensitivity to the $I = 2$ $\pi\pi$ channel. Measurements of the $\pi^+\pi^-$ atom proposed at several laboratories, if feasible with reasonable statistics, could provide the needed additional information. The quantity to be measured is the decay rate of the $\pi^+\pi^-$ ground state into the $\pi^0\pi^0$ channel, which is proportional to $|a_0^0 - a_0^2|^2$.[37]

Second, the mounting volume and accuracy of the exclusive and inclusive near-threshold $\pi N \rightarrow \pi\pi N$ data contain, in principle, valuable information regarding threshold $\pi\pi$ scattering. Theoretical interpretation of this data requires much improvement in order to make full use of this information. The work of Bernard, Kaiser and Meißner[38] appears to be a promising step in that direction.

3. CVC, the Pion Beta Decay, and Extensions to the Standard Model

Numerous hadronic processes, besides $\pi\pi$ scattering, are dominated by chiral symmetry at low energies; however, to appreciate fully the implications of chiral symmetry, we turn our attention to the weak interactions of hadrons. Since weak interactions do not conserve hadronic flavor, ChSB effects are pronounced in weak decays and transitions.

A linear transformation changes the (L,R) fermion states to the (V,A) basis,

$$\begin{array}{lllll} \psi_V = \frac{1}{2}(\psi_R + \psi_L) & , & \psi'_V = \exp(i\vec{\alpha}_V \cdot \vec{\tau})\psi_V & , & \alpha_V^a = \frac{1}{2}(\alpha_R^a + \alpha_L^a) \\ \psi_A = \frac{1}{2}(\psi_R - \psi_L) & , & \psi'_A = \exp(i\vec{\alpha}_A \cdot \vec{\tau})\psi_A & , & \alpha_A^a = \frac{1}{2}(\alpha_R^a - \alpha_L^a). \end{array} \quad (5)$$

As mentioned in Sect. 1, unlike axial-vector transformations, vector rotations in hadronic flavor space leave the vacuum invariant, preserving a chiral symmetry subgroup $\mathrm{SU(N)}_V \times \mathrm{U(1)}_V$ which is associated, through Noether's theorem, with a conserved vector current (CVC). Early recognition of this symmetry, along with the observation that the coupling constants in leptonic and hadronic beta decays are nearly the same, led to the concept of quark-lepton (Cabibbo) universality. Study of pure vector (Fermi) beta decays to date has confirmed the CVC hypothesis to a high degree. An interesting aspect of pure vector semileptonic transitions is their sensitivity to the physics beyond the Standard Model through the Cabibbo–Kobayashi–Masakawa (CKM) unitarity.

3.1. Motivation for a Precise Determination of the $\pi\beta$ Decay Rate

Pion beta decay, $\pi^\pm \rightarrow \pi^0 e^\pm \nu$, is one of the fundamental semileptonic weak interaction processes. It is a transition between two spin-zero members of an isospin triplet, and is therefore analogous to superallowed pure Fermi transitions in nuclear beta decay. Due to its accuracy of theoretical description, Fermi beta decay stands today among the prominent successes of the Standard Model (SM) of electroweak interactions.

The rate of pure Fermi nuclear beta decays is directly related to the vector coupling constant G_V^2 through the CVC hypothesis. Two types of nucleus-dependent corrections are needed in order to extract G_V from measured decay rates: (a) phase space dependent radiative corrections δ_R due to electroweak and QCD loop diagrams, and (b) nuclear structure dependent corrections δ_C.

Radiative corrections of $O(\alpha)$ and the leading terms of $O(Z\alpha^2)$ have been calculated, while contributions of $O(Z^2\alpha^3)$ have been estimated.[39] The resulting δ_R's amount to 3–4%, typically, with a theoretical uncertainty of 0.2% arising mainly from short-distance axial-vector-induced contributions.[40].

On the other hand, the nuclear corrections δ_C are smaller, ranging from 0.2 to 1%, but their evaluation appears to be accompanied by significant ambiguities.[41] Past evaluations of δ_C have yielded inconsistent results[42, 43, 41] The discrepancy was recently removed by Hardy *et al.*[44] Using their result, one obtains $|V_{ud}|^2+|V_{us}|^2+|V_{ub}|^2 = 0.9962 \pm 0.0016$, i.e., a value 2.4σ lower than the SM prediction.

V_{ud} can also be determined from neutron beta decay using τ_n and G_A/G_V. The current world average of the two quantities yields $V_{ud} = 0.9804 \pm 0.0022$, bringing $|V_{ud}|^2+|V_{us}|^2+|V_{ub}|^2$ up to 1.0096 ± 0.0044, which also differs from the SM prediction by 2.2σ, but in the *opposite* direction.

Thus, the determination of V_{ud} and the test of universality is presently not in a satisfactory state at the level of a few tenths of a percent. Determination of the $\pi^\pm \to \pi^0 e^\pm \nu$ rate at the same level is, therefore, important. Already at about 0.5%, such a measurement would provide an independent test of universality in a meson. It is of particular interest to find out where the pion beta decay result will fall, whether it will be consistent with the nuclear Fermi decays, or with the neutron decay data.

Pion beta decay is, furthermore, free of nuclear screening and overlap corrections δ_C, and involves a simpler axial-vector contribution to the radiative correction δ_R.[40] The overall theoretical uncertainty involved in evaluating G_V^2 or $|V_{ud}|^2$ from an experimental $\pi\beta$ decay rate is thus $\leq 0.1\%$ [40, 45] The only other significant source of uncertainty in such an evaluation involves the pion mass difference $\Delta = m_{\pi^+} - m_{\pi^0}$. Δ has been recently determined[46] to an accuracy of 10^{-4}, and thus contributes just 0.05% towards a total uncertainty of $\leq 0.11\%$. The theoretical advantage is, however, offset by experimental difficulties due to the small $\pi\beta$ branching ratio of $\sim 10^{-8}$; hence the present experimental accuracy of 4%.[47] The following physics becomes accessible through a precise constraint on the CKM unitarity:

(a) A possible fourth fermion generation could manifest itself through u–b' mixing inferred from an experimental violation of three-generation CKM unitarity.[40] Present data give only an upper bound of $V_{ub'} \leq 0.07$ at the 90% confidence limit.

(b) Another possible source of quark-lepton universality violation involves supersymmetric particles, as postulated in SUSY extensions to the minimal SM. Barbieri *et al.* have pointed out that superparticle exchanges in μ and β decay give rise to a violation of quark-lepton universality, unless their masses are suitably constrained.[48] Significant regions in the $(m_{\tilde{q}}, m_{\tilde{\ell}}, m_{\widetilde{W}})$ space can be excluded on the basis of presently available superallowed Fermi nuclear β decay data. For example, given reasonable assumptions on the higgsino mass parameter range, i.e., $m_{\tilde{h}} \geq O(100\ \mathrm{GeV})$, the combination $m_{\tilde{q}} = 80$ GeV, $m_{\tilde{\ell}} = 25$ GeV, $m_{\widetilde{W}} = 30$ GeV is excluded.

(c) Possible additional neutral gauge bosons Z', besides the known γ and Z of the standard electroweak theory, would also affect CKM unitarity. Marciano and Sirlin have computed quantum loop corrections to the quark mixing angles resulting from several types of Z'.[49] Present data infer lower bounds on $m_{Z'}$ ranging from 254 to 413 GeV, depending on the choice of the GUT model for the Z'.[45, 49]

(d) Compositeness of elementary fermions and vector gauge bosons presents a possible way to extend the SM. Sirlin has examined the consequences of several models of compositeness on electroweak radiative loop corrections.[50] In this framework, present experimental limits on the validity of quark-lepton universality can be used to constrain the mass scale of compositeness to the range above $\Lambda \geq 7$ TeV.[45]

Above constraints are comparably restrictive to those extracted so far from direct high-energy measurements. Thus, the present level of theoretical accuracy in the interpretation of vector β decay processes enables a selective exploration of certain phenomena on a mass scale largely inaccessible to existing accelerators.[49, 50, 45] A new, more precise measurement of the pion beta decays is particularly well suited to extending further the above limits on new physics.

3.2. The PIBETA experiment at PSI

A new project is under way at the Paul Scherrer Institute (PSI) with the aim of improving the present accuracy of the pion beta decay rate by an order of magnitude. This objective is being pursued in stages; the goal of the first phase is a 0.5% measurement. The experiment is a collaboration of seven institutions: the University of Virginia, PSI, Arizona State University, INS Swierk, JINR Dubna, Tbilisi State University and the Rudjer Bošković Institute.[51]

The PIBETA experiment relies on detection of pion decays at rest. An intense low energy π^+ beam $(1 - 2 \times 10^6 \pi/\mathrm{s})$ is stopped in a fragmented active target made up of thin scintillating fibers. The signature of $\pi\beta$ decays is given by the pair of back-to-back photons following π^0 decay. The photons are detected in a segmented 240-module pure CsI electromagnetic shower calorimeter subtending 3/4 of 4π. The calorimeter is spherical in shape, made up of truncated irregular hexagonal and regular pentagonal pyramids. The thickness of the pure CsI modules is 22 cm, or 12 radiation lengths. The inner radius of the CsI sphere is 26 cm, and the outer 48 cm. Light produced by showers is read out by photomultiplier tubes with quartz windows. The associated positron in $\pi\beta$ decay is also detected inside the active target.

Measurement of pion beta decays at rest allows the normalization to $\pi \to e\nu$ decays with nearly the same systematics (acceptance, efficiency, energy resolution). Present accuracy of the $\pi \to e\nu$ decay branching ratio, $\simeq$0.3%,[52] suffices for the first phase of the PIBETA project. The $e\nu$ signal involves only one monochromatic positron-induced shower and, therefore, requires a more demanding trigger. By far the most important backgroud in both $\pi\beta$ and $e\nu$ detection is presented by the numerous Michel positron events coming from the decay of muons in the active target, $\mu \to e\nu\bar{\nu}$.

In order to separate the various classes of events cleanly, there are two concentric cylindrical MWPCs for tracking in the central detector region. Also, an array of thin plastic scintillator hodoscope detectors, placed concentrically around the MWPCs serve to identify charged particles and provide accurate timing. Figure 2 shows the schematic cross section of the PIBETA detector with all of its main components. The entire detector system will be enclosed inside a temperature and humidity stabilized container, surrounded by and array of cosmic muon veto counters.

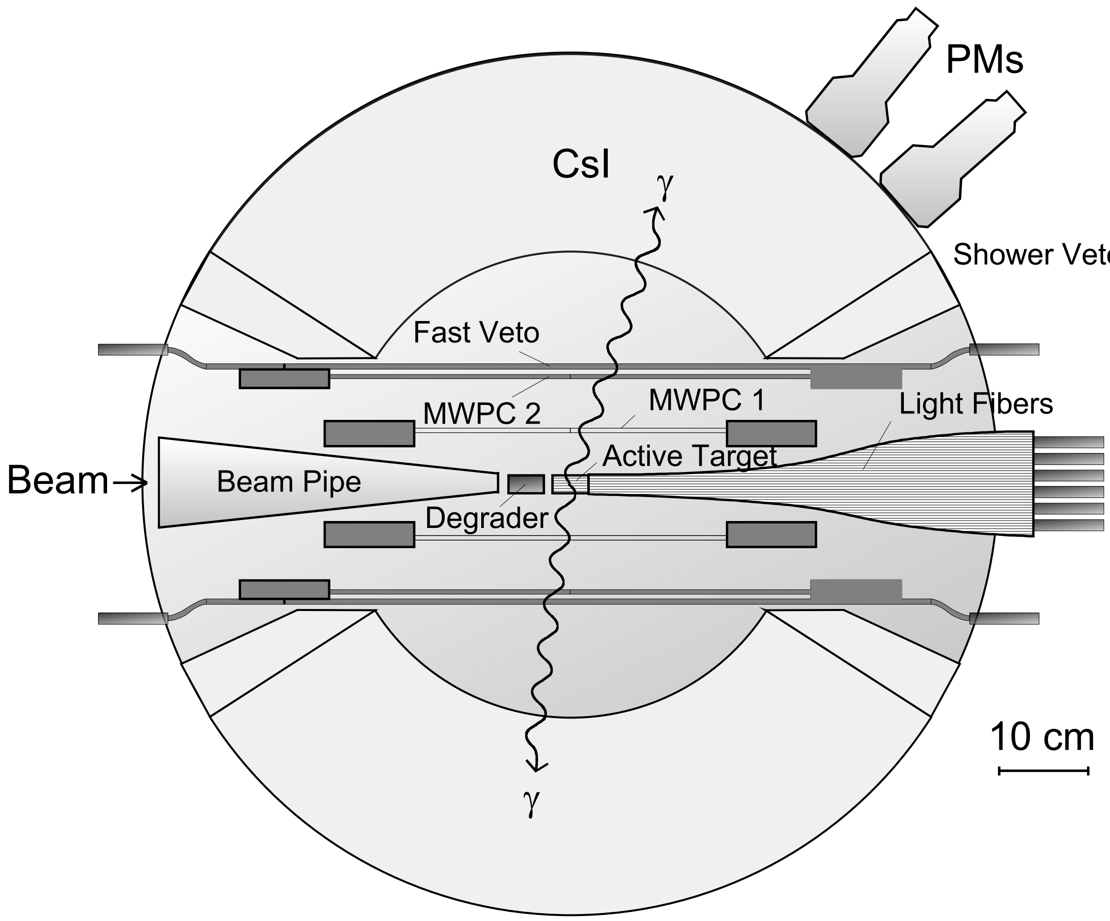

Figure 2: Schematic cross section of the PIBETA apparatus at PSI.

In a series of development runs since 1992, the PIBETA collaboration has addressed and solved the outstanding problems: pion stopping density, muon and positron contaminations in the beam (held to below 15% each), pure CsI surface treatment and energy resolution (currently about 4.2 MeV FWHM for 70 MeV showers), the fast analog shower clustering required for the main trigger modes, the $\pi \to e\nu$ trigger, etc. Quality control for the pure CsI detector modules (uniformity and percentage of fast component in the light output) is performed in a dedicated 3-dimensional cosmic muon tomography setup.

Plans call for the detector system to be assembled and completed by the end of 1995, while data taking runs are scheduled to begin in 1996. Due to the high required measurement precision, accumulating the necessary number of events will take about 10^7 s of running, or most of a calendar year in real time. With all the necessary systematic checks and calibration measurements, it will take several years after detector commissioning to complete the first, 0.5% phase. The second phase, which is to increase the accuracy by another factor of 2–3, will require instrumental improvements, and will follow the completion of the first phase.

4. Conclusions

The brief discussion of just two aspects of chiral symmetry in the pion sector serves to illustrate several points pertinent to this subject of study. First, chiral invariance is intimately related to the pion dynamics at low energies; indeed, the pion's very existence is tied to chiral symmetry trough the Goldstone mechanism. Given the central role chiral dynamics plays at low energies, the pion sector offers a clean testing ground for effective strong lagrangians. Second, details of the spontaneous ChSB also affect fundamentally the weak interactions of hadrons, making accessible certain high energy aspects of the Standard Model physics through very precise low energy measurements. All of these considerations continue to make the pion a rich source of information about some of the most basic aspects of the strong and weak interactions, even after more than three decades of research.

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