NEUTRINO '72

I.

HUNGARY 1972

OMKDK-TECHNOINFORM

NEUTRINO '72

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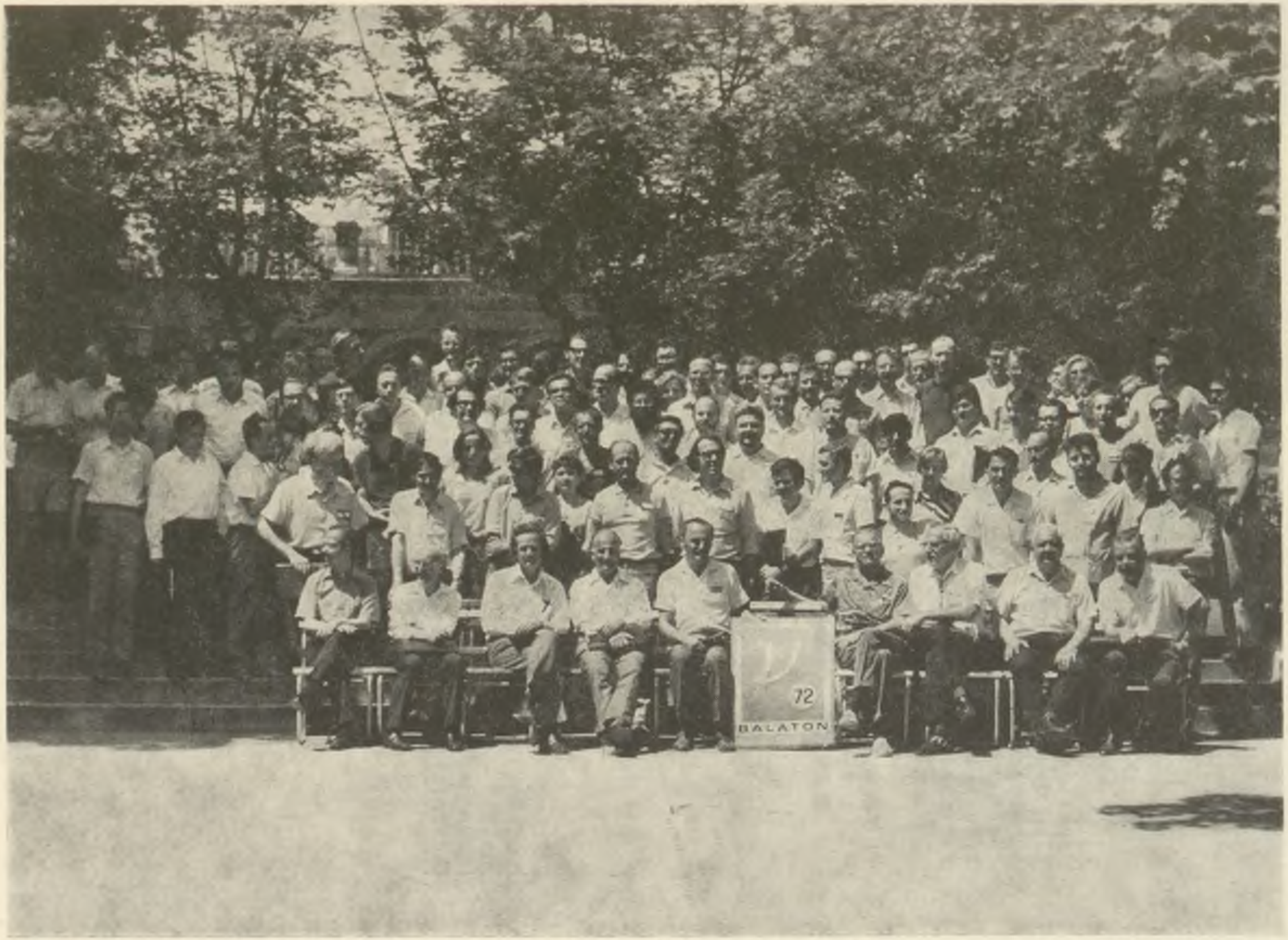

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VOLUME I

VOLUME II

NEUTRINO'72 (OPENING ADDRESS)

G.Marx, Department of Atomic Physics, Roland Eötvös University in Budapest

We may now celebrate the 40th anniversary of the discovery of neutrino. The understanding of weak forces acting in Nature was always directly connected with the investigation of this tiniest and simplest piece of matter. The four decades of neutrino research form a brilliant success story. Pauli told at the beginning of neutrino science that neutrinos would never be detected. But Fred Reines and Clyde Cowan taught us the way of experimenting with neutrinos. To-day sophisticated neutrino eyes are watching at the glowing of different sources in several laboratories.

NEUTRINO EXPERIMENTS

from reactors:	Hanford	1953
	Savannah River	1955
	Brookhaven	1956
from accelerators:	Brookhaven	1962
	CERN	1964
	Batavia	1973
from cosmic rays:	Home Stake	1965
	Kolar Gold Field	1967
	Utah	1969

Doing fundamental research is a hard job in our age. For tremendous efforts Nature pays only with very faint and very provisional results. But weak interactions helped us in obtaining valuable bits of exact information about the fundamental properties of matter. The last discovery of a strict conservation law was that of the leptonic charges /1952/. The C V C theorem built a bridge from electricity to beta decay /1957/. The weak interaction taught us, what was the essential difference between left and right /1956/, positive and negative /1957/, past and future /1964/.

The Universe is crowded by neutrinos. They are more plentyful, than protons or electrons. Neutrinos are not only the simplest, but also the most common forms of matter. We learned them to know earlier, than the neutrons or positrons or mesons. But neutrino science is far from being closed. Just the opposite statement is true: it is getting more and more puzzling in the recent years. We have several well-formulated questions, which are waiting for answer. We discovered the muon long ago, but we did not understood it up to now. /Its existence contradicts sharply all of our classic ideas about the origin of particle masses./ On the other hand, we tried hard, but we are unable to discover the W boson. /Consequently we still do not know, if field theory is valid anywhere but quantum electrodynamics, or not./ And we have also new puzzles: We observed the decay of the K_L^o meson into $\pi^+\pi^-$, what we did not want to observe. /We learned, that Nature is asymmetric with respect to time reversal, but the only statement we are able to formulate is that this asymmetry is superweak and almost unobservable./ We did not observe the decay of K_L^o into $\mu^+\mu^-$ what we needed desperately. /In this hide-and-seek game around the K_L^o meson we do not see, if there is any connection between the positive and negative surprises/. The neutrino eyes do not see the sunshine either.

We have these tricky problems, because we asked Nature on the sophisticated language offered by the modern experimental technique. But the systematic knowledge about weak interactions and the experimental ability to handle neutrinos will help us at the new frontiers of exact science, what we have reached: at the deep inelastic frontier inside the hadrons, at the supernova singularity in the late stellar evolution and at the Big Bang singularity at the beginning of time. There are good hopes, that the tiny neutrinos may show the way to the pioneers in these virgin lands.

The little old neutrinos may come again to the headlines of scientific journals. Feeling this, Prof Zatsepin called a specialized meeting to Moscow in 1969, in order to discuss the problems of cosmic neutrinos. Prof.Bernardini and Prof.Radicati called another meeting on stellar neutrinos to Cortona in 1970. On the Cortona conference we agreed to organize the third European neutrino conference in Hungary.

In our organizing work we enjoyed the sponsorship of the European Physical Society, the positive interest of the Joint Institute of Nuclear Research /Dubna/, CERN /Geneva/, International Centre for Theoretical Physics /Trieste/. The main sources of the financial support were the Hungarian Academy of Sciences, the Hungarian Physical Society, the Roland Eötvös University and the Central Research Institute in Budapest. But the Hungarian efforts were not enough, to overcome all the difficulties in the organizing work. We made a strong use of the Triangle Collaboration. This scientist's co-op started four years ago, according to the suggestion of Walter Thirring. The corners of the original Triangle were formed by the particle physics groups in Vienna, Bratislava and Budapest, but in the meantime also other nearby scientific centers joined us.

This is about the organization of this conference. Let the Moscow, Cortona and Balaton Meetings be followed by several Neutrino Conferences!

„THE TRIANGLE COLLABORATION"
IN PARTICLE PHYSICS

REPORT ON THE BROOKHAVEN SOLAR NEUTRINO EXPERIMENT*

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(Abstract)

A solar neutrino experiment has been in progress since 1967 to observe neutrinos from the sun by the neutrino capture reaction $^{37}Cl(\nu,e^-)^{37}Ar$. 380,000 liters of perchloroethylene, C_2Cl_4, are used as a neutrino capturing medium. The radioactive ^{37}Ar (35.1-day half-life) is removed from the liquid by purging with helium gas, purified, and counted in a small low-level proportional counter. Argon-37 decay events are characterized by the energy of the event, and the rise-time of the counter pulse. The results of six experiments performed during the last two years will be given. The ^{37}Ar production rate in the detector was 0.18 ± 0.10 per day, close to the production rate expected from cosmic ray muons. The experiment places an upper limit on the production of ^{37}Ar by solar neutrinos of 1×10^{-36} sec^{-1} $(^{37}Cl\ atom)^{-1}$. This limit can be compared to the currently predicted rate of 9×10^{-36} sec^{-1} $(^{37}Cl\ atom)^{-1}$ from solar model calculations of Bahcall and Ulrich, and Abraham and Iben.

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A solar neutrino detector based upon the neutrino capture reaction $^{37}Cl(\nu,e^-)^{37}Ar$ was built in 1964-1967 and has been operating since early 1967. The initial experiments [1] showed that the total flux-cross section product was less than 3×10^{-36} sec^{-1} $(^{37}Cl\ atom)^{-1}$. This conclusion was

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*Research performed under the auspices of the U.S. Atomic Energy Commission.

substantiated by additional experiments performed during the period 1968-1970, and reported in a series of reports [2]. In these experiments the ^{37}Ar decay events were characterized in the proportional counter by observing the energy of the Auger electrons emitted following the decay. The recent experiments that will be reported here, the ^{37}Ar decay event in the counter is characterized by the energy (pulse height) and rise-time of the pulse [3]. This added discrimination for ^{37}Ar events has lowered the counter background for ^{37}Ar-like events permitting a more sensitive search to be made for the neutrino radiation from the sun.

This report will present a brief summary of the critical experimental tests that have been performed to demonstrate the efficiency of recovering ^{37}Ar from 380,000 liters of perchloroethylene, the rise-time counting system, and the observations that have been made over the last two years. The detector was arranged so that it could be shielded with a water shield with an average thickness of about 2 meters to eliminate the effect of fast neutrons from the surrounding rock. Three experiments will be reported with this water shield in place.

The Detector

The detector consists of a 380,000 liter tank of perchloroethylene (C_2Cl_4), a pumping system for circulating the liquid and purging the liquid with helium gas, and a gas absorption system for removing argon from the circulating helium gas. The detector is located at a depth of 4400 hg/cm^2 in the Homestake Gold mine at Lead, South Dakota. The arrangement of the apparatus underground is shown in Fig. 1, and a schematic diagram of the extraction system is shown in Fig. 2.

Argon-37 produced in the liquid is removed by purging with helium gas. Two 1700 liter/min pumps circulate liquid from the tank through a series of

40 eductors that draw helium from the gas space at the top of the tank, and distribute it through the liquid in the form of small bubbles. The 20,000 liters of helium in the tank at 1.3 atmospheres pressure is circulated through the liquid every 2 minutes by the pumping system. Two eductors are used to draw the helium from the tank through a condenser at -35°C, a molecular sieve adsorption trap, a charcoal trap cooled to 77°K, and finally returning the helium to the tank. Helium flows through the extraction system at a rate of 310 liters/min. The helium flow rates depend upon the gas pumping rates of the eductors, and these rates have remained constant at the values given above.

The efficiency for extracting argon from the 380,000 liters of C_2Cl_4 was tested in several ways.

(1) In each experiment a measured volume (0.05 to 0.10 cm^3) of isotopically pure ^{36}Ar is introduced into the tank to serve as a carrier gas and to measure the argon recovery achieved in each experiment [1]. The ^{36}Ar carrier is introduced at the beginning of the period of irradiation and is recovered several months later by the helium purge. The ^{36}Ar content of the helium gas in the tank was analyzed, and from the relative volumes of the gas and liquid phases we calculate that approximately 95 percent of the ^{36}Ar introduced dissolves in the liquid. The recovery of ^{36}Ar as a function of the volume of helium circulated through the system was measured. The volume of argon remaining in the tank V_R decreases exponentially with the volume of helium circulated [1]. With the fixed flow rates, a 95 percent recovery can be obtained in a 22 hour period.

(2) An analogous test was made with ^{37}Ar. A sample of ^{37}Ar activity in helium gas was prepared, introduced into the tank, recovered during 22 hours of operation extending over a 2 day period, and counted. The ^{37}Ar

activity introduced was 10.0 dis ^{37}Ar/day, and the recovered sample contained 11.0 $\pm$ 1.8 dis ^{37}Ar/day corrected to the time it was introduced. This experiment was performed May 19, 1972, and the counting of this sample and the analysis is not yet complete. However, this preliminary result does demonstrate the recovery of ^{37}Ar is consistent with recovery based on the ^{36}Ar carrier method.

(3) The 380,000 liter tank is provided with a reentrant pipe that extends to the center of the tank. A Ra-Be neutron source was placed in this source tube to produce ^{37}Ar in the liquid by the successive nuclear reactions $^{35}Cl(n,p)^{35}S$ followed by $^{37}Cl(p,n)^{37}Ar$. The yield for this neutron source test was 7.4×10^{-7} ^{37}Ar atoms per neutron, and it was shown that the ^{37}Ar produced was removed in successive purges along with the ^{36}Ar carrier [1]. The yield of ^{37}Ar activity obtained in this experiment was compared to similar neutron irradiations of containers of perchloroethylene with smaller diameters. These yields (^{37}Ar atoms/neutron) and diameters were: 3.0×10^{-7}, 29 cm; 6.4×10^{-7}, 120 cm; 7.4×10^{-7}, 600 cm.

These experimental tests demonstrate that ^{37}Ar produced in the 380,000 liters of perchloroethylene is efficiently removed by the helium purge. This result is entirely reasonable based upon the chemical behavior of argon. The neutrino capture process forms an argon ion that recoils from its parent C_2Cl_4 molecule. The argon ion rapidly reaches thermal energy in the liquid, capturing an electron from adjacent C_2Cl_4 molecules becoming a neutral ^{37}Ar atom. This neutral ^{37}Ar atom is identical chemically with the dissolved ^{36}Ar carrier atoms present in the liquid at a concentration of 5×10^{9} atoms/cm^3. The recovery of ^{36}Ar from the liquid is therefore a valid measure of the ^{37}Ar recovery. Our colleagues at Brookhaven [4] have made a search for the formation of complex argon-C_2Cl_4 ions in a high pressure mass spectrometer source. No evidence was obtained in these experiments for the formation of stable $C_mCl_nAr^+$ ions. Evidence was found for a charge transfer process that proceeds with a rate coefficient at least two orders of magnitude greater than an upper limit for the rate coefficient for the reactions producing $Ar\text{-}C_mCl_n^+$. This is evidence for rapid charge transfer in the gas phase and supports the assumptions made above concerning the charge transfer process in the liquid phase.

<u>Counting</u>

The entire argon sample recovered from the tank was placed in a small proportional counter that has already been described [1]. In the experiments reported here the rise time of the pulse and the pulse height was measured for each pulse. The 2.8 keV Auger electron produced in the electron-capture decay of ^{37}Ar has a range in the counter gas of approximately 0.1 mm. The ion pairs produced in the gas are localized in a small volume of the counter, and therefore they reach the center wire nearly at the same time giving a fast rising pulse. Background events from more energetic Compton-electrons, betas or muons in passing through the counter gas, produce ion-pairs more widely spaced, and hence reach the center wire over longer periods of time giving a slower rising pulse. By measuring the rise-time of the pulse one can distinguish ^{37}Ar decay events, X-ray absorption events, and tritium decays from background events produced by gamma radiation, betas, and muons.

A fast charge-sensitive preamplifier with a 6 nsec pulse rise-time response feeds an Ortec model 410 timing filter amplifier that integrates (10 nsec) and differentiates (10 nsec) the charge pulse. The shaped pulse is then fed to a pulse-stretching amplifier for acceptance by an analog to digital converter. This digital information is plotted on a y-axis of a two parameter plot and is designated the amplitude of the differentiated pulse (ADP). The pulse from the preamplifier is also separately routed to a conventional RC shaping spectroscopy amplifier and subsequently analyzed by a second analog to digital converter. This digital information is displayed on the X-axis as the energy. The arrangement is shown schematically in Fig. 3. Each time the counter functions the ADP and pulse height (energy) is written on a paper tape along with the clock time, and a sequence of test pulses for ADP and pulse height (energy). The system is automatically checked once every 1000 minutes with the test pulses, and at this time the anti-coincidence counter rate is recorded. The counting system is capable of recording events from four separate counters.

The counters are calibrated with an ^{55}Fe X-ray source. The location of region of fast-rising X-ray pulses is determined by recording on a separate pulse height analyzer the ADP pulses gated in coincidence with a narrow energy region selected by a single channel analyzer. This measurement is carried out at several different energies to define the region of X-ray pulses. The lines drawn on the two parameter plot define the X-ray region. The indicated region contains approximately 95 percent of the fast X-ray pulses. The location of the position of ^{37}Ar on the linear energy scale was determined from the ^{55}Fe X-ray peak position. The ^{37}Ar resolution (full width at half maximum, fwhm) was taken as 1.45 times the ^{55}Fe resolution.

The discrimination against pulses produced by gammas, was measured by placing a ^{60}Co source near the counter and recording the ADP pulse distribution corresponding to particular energies. The fraction of the area under the ADP spectrum exterior to the X-ray region is a measure of the effective rejection of pulses arising from energetic betas, gammas, and muons. In the experiments reported here there was a variation in the rejection ratio that depended upon counter filling (argon pressure, percent CH_4) and the dimensions of the counter.

The lowest rejection ratio was 82 percent, and the highest achieved was 99 percent.

Results

Three experimental runs (Nos. 18, 19, and 20) were made with the unshielded tank, as in all previously reported experiments. The two-parameter plots for these three runs are shown in Figs. 4 and 5. The periods of counting are indicated on the plots, and correspond approximately to a half-life of ^{37}Ar (35.1 days). It may be observed that 5, 9, and 3 counts, respectively, were observed with the correct ADP and an energy within the half-width of the ^{37}Ar peak position. The samples were counted for additional periods to see if these ^{37}Ar-like events decayed away with the 35.1-day half-life period. The data are summarized in Table 1 where the number of fast counts observed in the energy intervals 1 to 2.4 keV, 2.4 to 3.2 keV (fwhm), and 3.2 to 7 keV are given. It may be seen that there was satisfactory decay in runs 18 and 19 to a background counting rate of ~ 2 and ~ 3 counts per 36 days, respectively, but little evidence for decay in run 20.

Following these experiments the tank chamber was flooded with water to provide a neutron shield. Fast neutrons above 1 MeV energy can produce ^{37}Ar in the liquid by the successive reactions mentioned earlier. A small flux of fast neutrons arises from spontaneous fission of ^{238}U and (α,n) reactions from alpha decays of the uranium and thorium and their daughters present in the rock. The fast neutron flux was measured by a sensitive radiochemical technique based upon the $^{40}Ca(n,\alpha)^{37}Ar$ reaction. These measurements indicated that the ^{37}Ar production in the 380,000 liter tank was about 0.04 ^{37}Ar atom per day, too small to be observed. However, the water shield should completely eliminate any possible fast neutron background from the rock wall.

The number of fast counts observed in the indicated energy intervals for three experiments with the water shield in place, nos. 21, 22, and 23, are listed in Table 2. It is clear from these results that the number of counts that could be attributed to ^{37}Ar is low, less than or equal to 1 per experiment. In Run 23 a large counter (5 cm^3) with a correspondingly high background was used. However, this particular counter had an exceedingly high rejection ratio, and thus a low background for fast counts. The three fast counts observed in the ^{37}Ar peak position correspond approximately to the expected background of this counter.

A detailed statistical analysis of the data has not been made, though it is evident that the counting rates are indistinguishable from counter background for experiment nos. 20 (without water shield), 21, 22, and 23 (incomplete decay information), and an indication of a positive amount of ^{37}Ar in the two runs nos. 18 and 19 without the water shield. The data can be used to derive an upper limit to the solar neutrino flux. We have taken the experiments with the water shield in place and the observed number of fast counts in the energy region 2.4 to 3.2 keV as ^{37}Ar decays. Accepting all of these counts as ^{37}Ar without allowing for counter background should be a safe procedure for setting upper limits. Correcting for growth factor, decay factor, counter efficiency, and recovery yield we derive an ^{37}Ar production rate in 380,000 liters of C_2Cl_4 of 0.18 $\pm$ 0.10 per day.

Background Effects

Cosmic ray muons and cosmic ray produced muon-neutrinos can produce ^{37}Ar in the liquid. The ^{37}Ar production from muons was evaluated by exposing 6800 liters of C_2Cl_4 arranged in 3 tank cars at various levels in the mine from 285 to 1080 hg/cm^2 depth. An analysis of this data is given in another paper presented at this Conference [5]. It is concluded that the ^{37}Ar

production rate from muons in 380,000 liters of C_2Cl_4 at a depth of $4400\ hg/cm^2$ is approximately 0.12 ± 0.04 per day. Argon-37 can also be produced by cosmic ray produced muon-neutrinos, ν_μ, by the reaction $^{37}Cl(\nu_\mu,\mu^-)^{37}Ar$. The production rate from this process was estimated to be 0.024 per day [5].

Some attention was devoted to the possibility that there is ^{37}Ar in the earth's atmosphere. In the processing of the 380,000 liters of C_2Cl_4 an excess of ^{40}Ar is found in each experiment, and this is attributed to the inleakage of atmospheric argon. Quantities of ^{37}Ar have been found in the atmosphere that are two orders of magnitude higher than the amount produced by cosmic ray interactions [6]. The level was high during the period corresponding to our run no. 19, however the specific activity was a factor of ten below that observed in the run 19 argon. The air in the Homestake mine was monitored during the period March (1971) to March (1972) but the ^{37}Ar content was in general low, <0.15 dis/min liter of argon.

Conclusions

The cosmic ray muon background must be subtracted from the ^{37}Ar production rate limit of 0.18 ± 0.10 given above. The resulting rate that could be attributed to solar neutrinos is then 0.06 ± 0.11 ^{37}Ar atoms per day. We then conclude that the solar neutrino production rate is less than 0.2 ^{37}Ar atoms/day. The corresponding flux-cross section product limit is then,

$$\Sigma\phi\sigma < 1 \times 10^{-36}\ sec^{-1}\ (^{37}Cl\ atom)^{-1}$$

This limit can be compared to the currently calculated flux-product from standard solar models of

$$\Sigma\phi\sigma\ (theoretical) = 9.1 \times 10^{-36}\ sec^{-1}\ (^{37}Cl\ atom)^{-1}$$

The significance of this discrepancy between observations will be discussed by Dr. J. N. Bahcall in a paper presented at this conference.

Our conclusions may be summarized as follows:

(1) The total flux-cross section product for solar neutrino capture in ^{37}Cl is a factor of 9 below the total (B^8, Be^7, and PeP) rate predicted by standard solar model calculations, and below the rate predicted from ^7Be decay and PeP reaction neutrinos

(2) The flux of ^8B solar neutrinos is less than 10^6 cm^{-2} sec^{-1}

(3) The sun produces less than 3 percent of its energy by the CNO cycle.

Acknowledgement

The authors would like to acknowledge the skillful technical assistance of Mr. John Galvin during the entire period of these investigations.

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Table 1

Summary of Results

Without Water Shield

Run No.	Exposure Period	Counting Period	Fast Counts 1.0-2.4 keV	Ar^{37} FWHM 2.4-3.2 keV	3.2-7 keV
18	Apr 12 to Nov 14 (1970)	I 39 days	10	5	5
		II 35 "	9	5	9
		III 34 "	12	3	6
		IV 36 "	6	1	9
		V 39 "	11	2	11
19	Nov 14 (1970) to Mar 6 (1971)	I 36 days	6	9	3
		II 39 "	5	5	13
		III 37 "	7	5	12
		IV 41 "	20	4	8
		V 37 "	20	3	8
		VI 36 "	7	3	5
		VII 34 "	5	4	8
20	Mar 6 to June 17 (1971)	I 40 days	3	3	2
		II 30 "	7	1	2
		III 37 "	15	3	7
		IV 38 "	8	0	3
		V 51 "	13	1	6
		VI 36 "	7	3	3

Table 2

Summary of Results

With Water Shield

Run No.	Exposure Period	Counting Period	Fast Counts 1.0-2.4 keV	Ar^{37} FWHM 2.4-3.2 keV	3.2-7 keV
21	Jan 17 to Oct 2	I 39 days	5	0	7
	(1971)	II 34 "	8	1	8
22	Oct 2 to Dec 13	I 35 days	6	2	3
	(1971)	II 36 "	9	2	5
		III 40 "	9	3	6
23	Dec 13 (1971) to Mar 2 (1972)	I 40 days	4	3	5

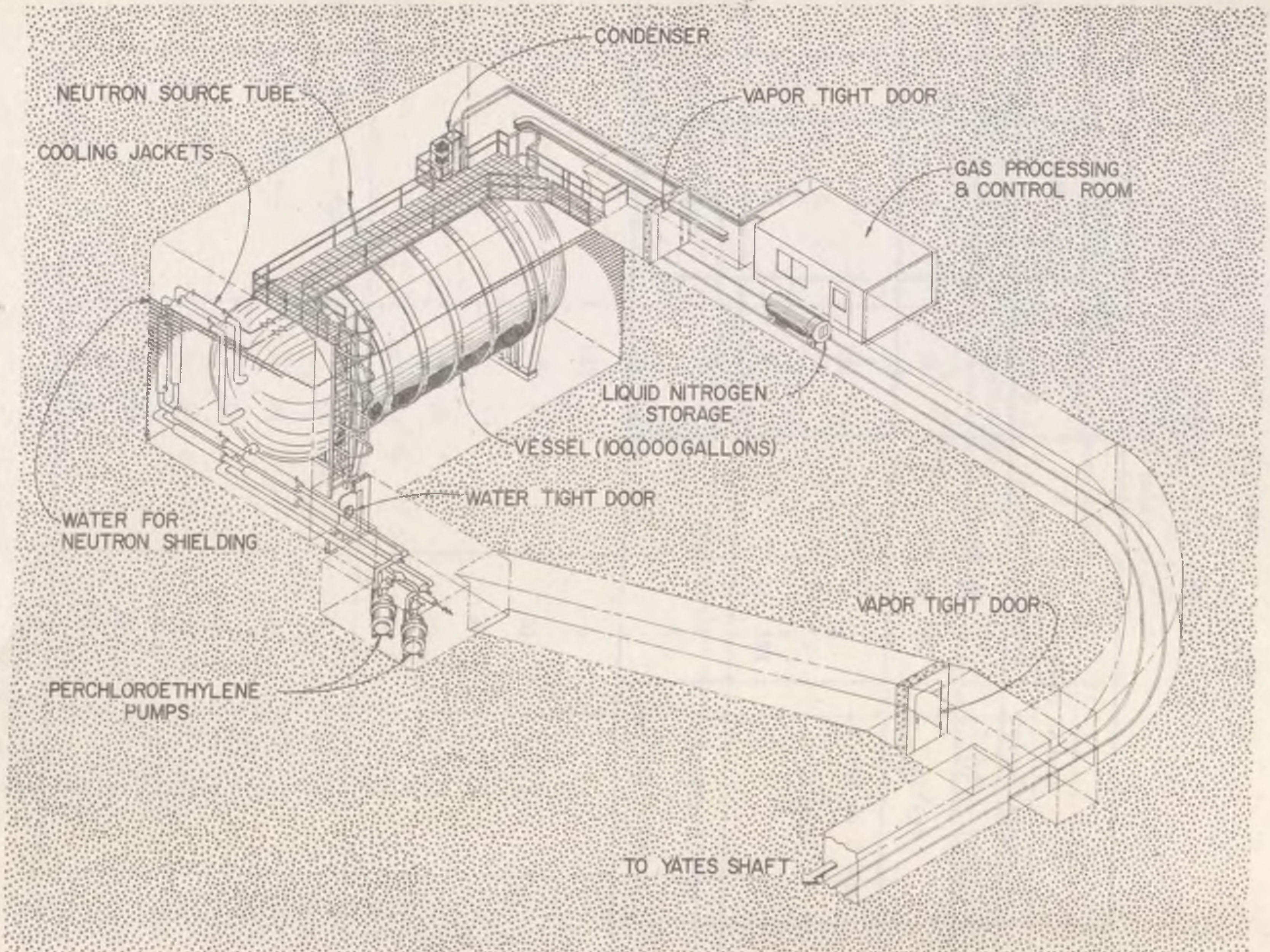

CONDENSER
NEUTRON SOURCE TUBE
VAPOR TIGHT DOOR
COOLING JACKETS
GAS PROCESSING
& CONTROL ROOM
LIQUID NITROGEN
STORAGE
VESSEL (100,000 GALLONS)
WATER TIGHT DOOR
WATER FOR
NEUTRON SHIELDING
VAPOR TIGHT DOOR
PERCHLOROETHYLENE
PUMPS
TO YATES SHAFT

Figure 1.

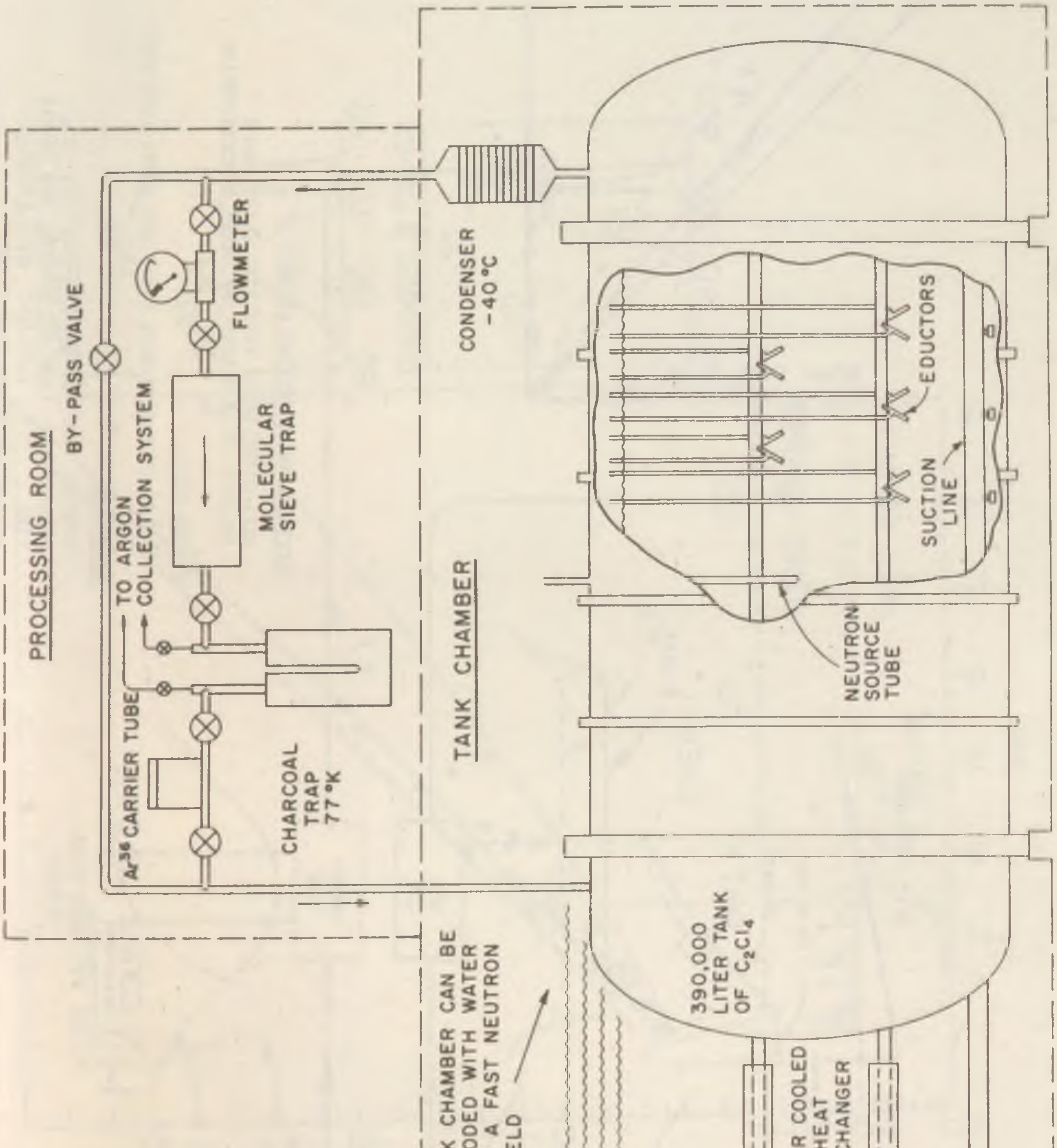
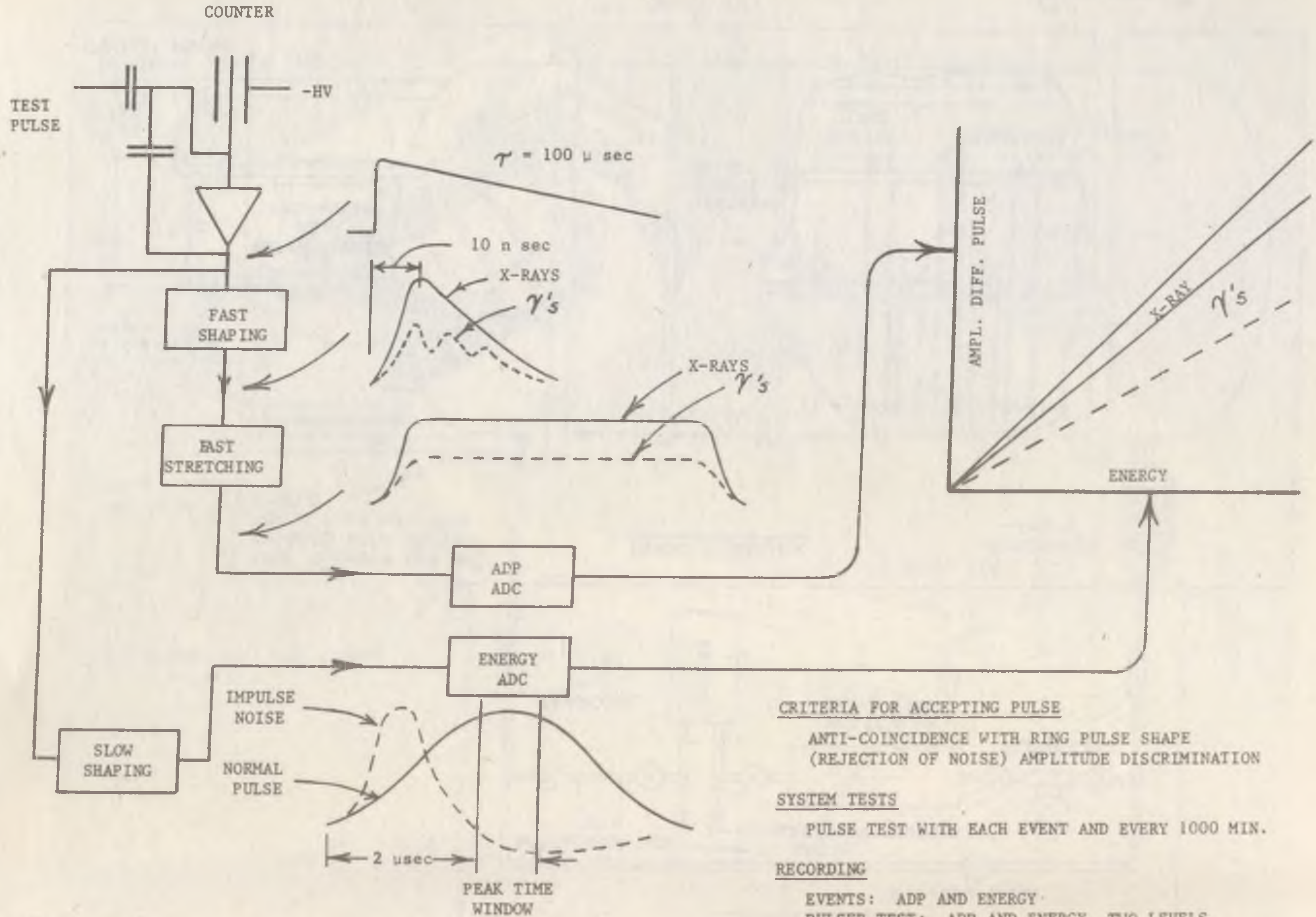

PROCESSING ROOM
BY-PASS VALVE
TO ARGON COLLECTION SYSTEM
MOLECULAR SIEVE TRAP
FLOWMETER
CONDENSER -40°C
Ar^{36} CARRIER TUBE
CHARCOAL TRAP 77°K
TANK CHAMBER
EDUCTORS
SUCTION LINE
NEUTRON SOURCE TUBE
390,000 LITER TANK OF C_2Cl_4
TANK CHAMBER CAN BE FLOODED WITH WATER FOR A FAST NEUTRON SHIELD
WATER COOLED HEAT EXCHANGER
PUMP ROOM
PUMPS 1700 LITERS/min

Figure 2.



CRITERIA FOR ACCEPTING PULSE

ANTI-COINCIDENCE WITH RING PULSE SHAPE
(REJECTION OF NOISE) AMPLITUDE DISCRIMINATION

SYSTEM TESTS

PULSE TEST WITH EACH EVENT AND EVERY 1000 MIN.

RECORDING

EVENTS: ADP AND ENERGY
PULSER TEST: ADP AND ENERGY, TWO LEVELS
TIME
ANTI-COINCIDENCE RING COUNTS
DEAD TIME

Figure 3.

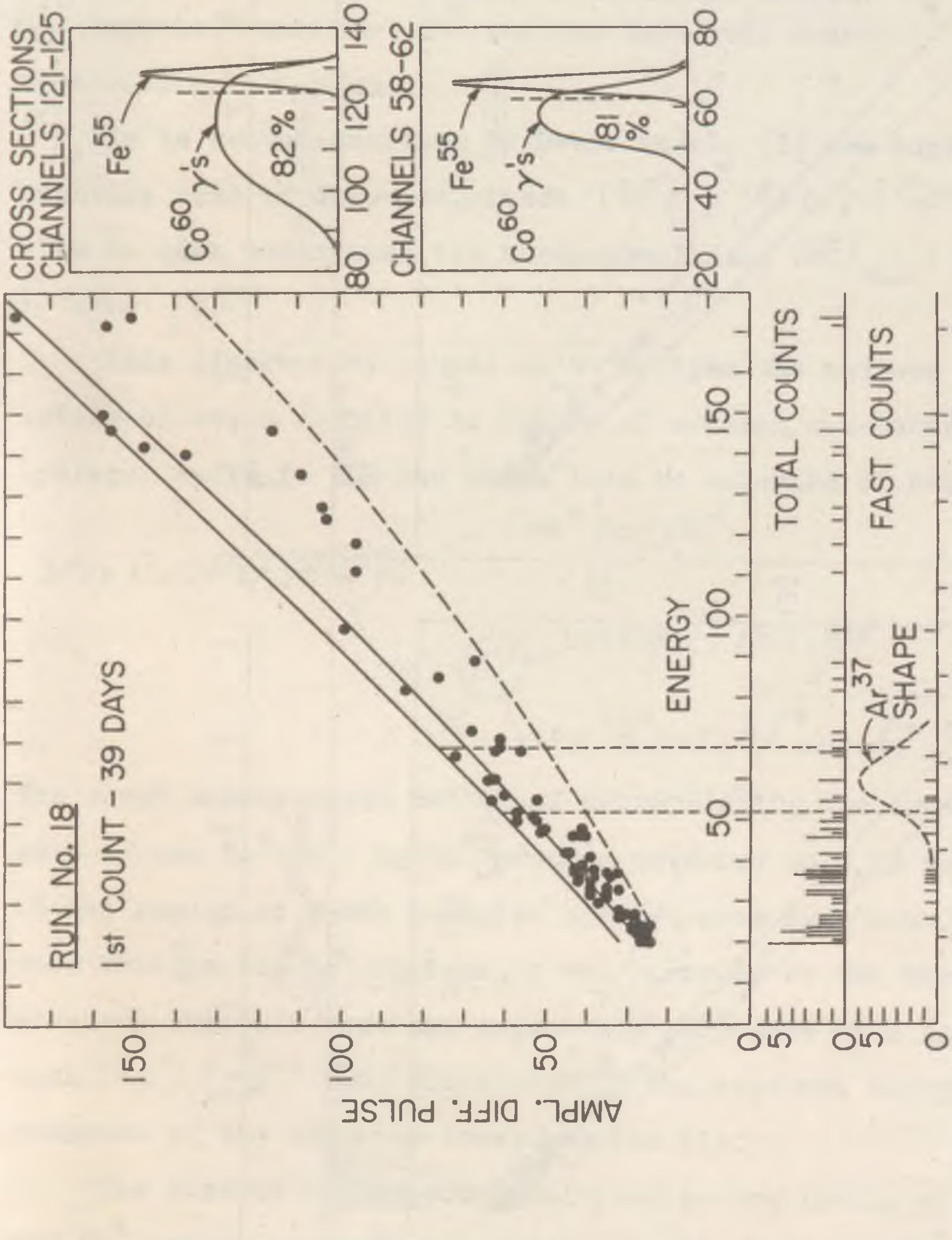

RUN No. 18
1st COUNT - 39 DAYS
AMPL. DIFF. PULSE
150
100
50
ENERGY
TOTAL COUNTS
FAST COUNTS
Ar37 SHAPE
keV
1
2
3
4
5
6
7
8
CROSS SECTIONS
CHANNELS 121-125
Fe55
Co60 γ's
82 %
80
100
120
140
CHANNELS 58-62
81 %
20
40
60
80

Figure 4.

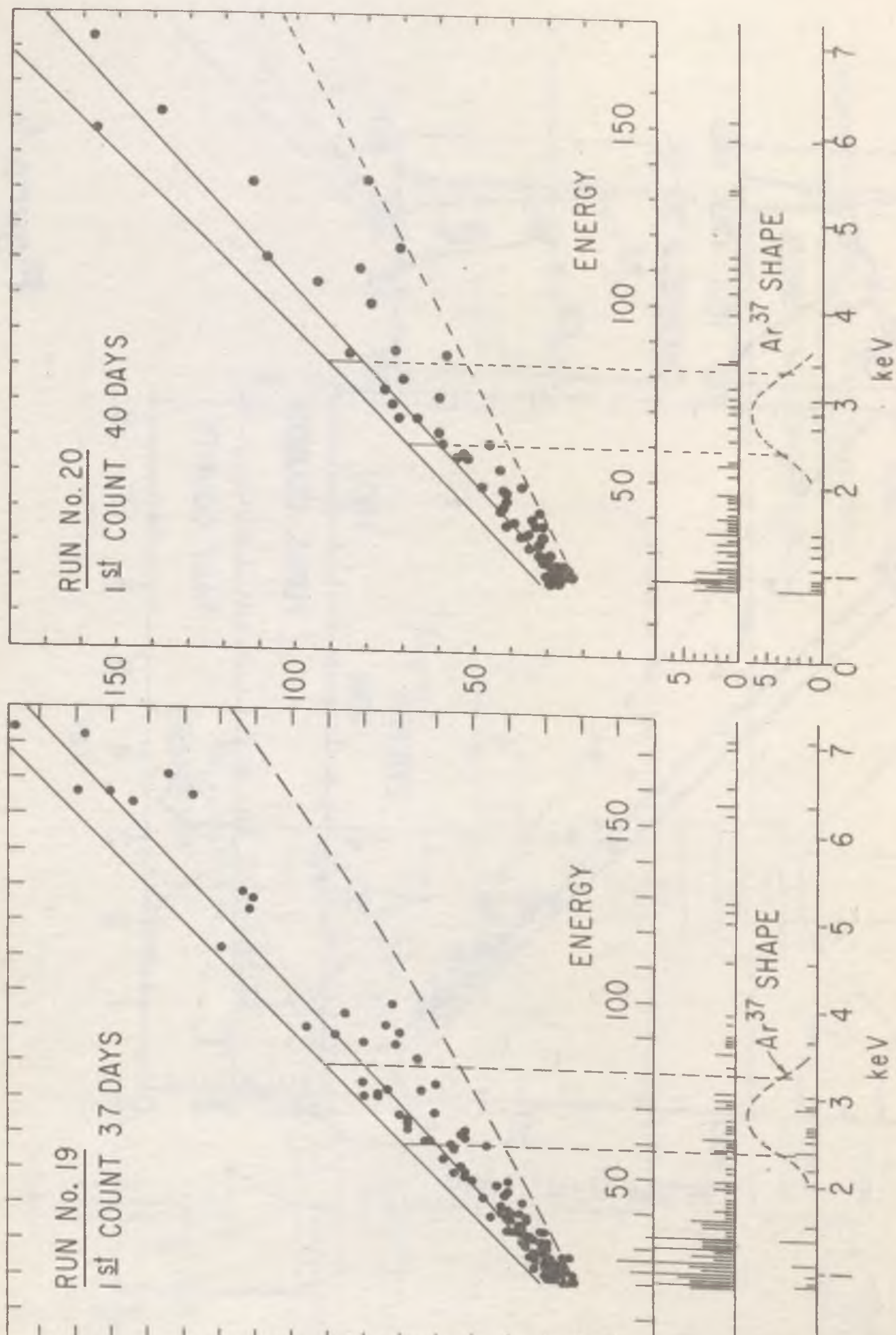

RUN No. 19
1st COUNT 37 DAYS
RUN No. 20
1st COUNT 40 DAYS
AMPL. DIFF. PULSE
ENERGY
Ar37 SHAPE
keV
0
50
100
150
0
5
0
1
2
3
4
5
6
7

Figure 5.

ARE THE SOLAR-NEUTRINO EXPERIMENTS SUGGESTIVE OF THE EXISTENCE OF A RESONANCE IN THE He^3+He^3 SYSTEM ? (DISCUSSION REMARK)

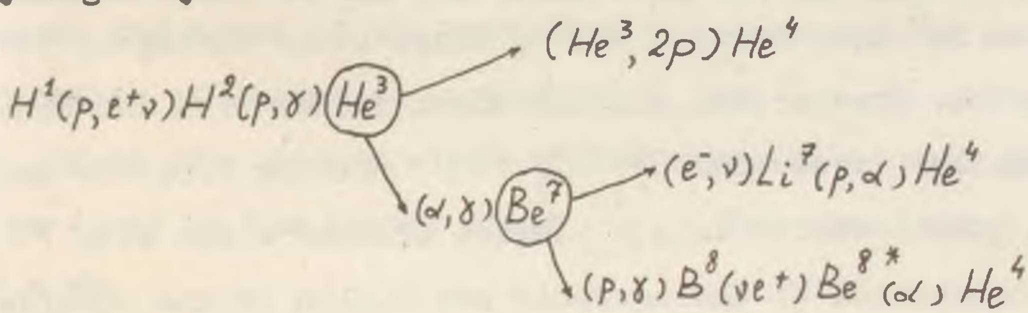
V.N.Fetisov, P.N.Lebedev Physical Institute, Moscow,

Yu.S.Kopysov, Insititute for Nuclear Research, Moscow

/Presented by A.E.Chudakov/

As is known according to Davis at al. [1] the experimental counting rate of solar neutrinos $(\Phi\sigma)_{exp.} = (1.5 \pm 1)\,10^{-36}\,sec^{-1}$ per Cl^{37} atom is much lower than the theoretical one: $(\Phi\sigma)_{theor.} = (9 \pm 5) \times 10^{-36} \times sec^{-1}\,atom^{-1}$ [2].

This discrepancy causes us to analyse the methods of calculation of cross sections in chains of nuclear reactions of the hydrogen cycle in the Sun which lead to emission of neutrinos:



(1)

(2)

(3)

The usual nonresonance method of extrapolating the cross section of the $He^3(He^3, 2p)He^4$ process measured only up to 80 keV [3] to the region of lower energies arouses some suspition. We will show that in the Be^6 nucleus in the vicinity of the threshold of break-up into $2He^3$ one can expect the existence of a narrow level with $J^\pi = 0^+$, $T = 1$ and, consequently, the resonant energy dependence of the reaction cross section (1).

Now turn to the experimental data on the levels of the Li^6 and He^4 nuclei. According to the analysis [4] in He^4 there exist dipole states with $J^\pi = 1^-$, $T = 1$ having the shell configuration $|s^3 p>$ in the excitation energy interval 27÷30 MeV and a lower monopole excited state with $J^{\pi}=0^{+}$, $T=0$ at $E^{*}=20.5$ MeV. From the experimental data on photonuclear reactions in Li^6 [5] and from the experiments [6] on quasifree scattering of protons by Li^7 it is known that at $E^{*}\simeq 18\div20$ MeV in the Li^6 nucleus there are excitations with the $|s^3p^3\rangle$ configuration corresponding in cluster terms to the dipole excitation of an α-particle in Li^6 [7]. Using the data on the spectrum of He^4 we find that when going from the dipole inner excitation of the α-cluster to the monopole one the corresponding level with $J^{\pi}=1^{+}$, $T=0$ in Li^6 falls within the region located $7\div10$ MeV below the group of dipole levels, i.e. somewhere around $E^{*}(Li^6)=12$ MeV. Since Be^6 and He^6 have isospin $T=1$, there is no analogous state in these nuclei. However, in all the three nuclei with A=6 there must be a similar level with $J^{\pi}=0^{+}$, $T=1$ arising from the excitation of a quasideuteron into a singlet cluster state with $T=1$ with the conservation of the monopole excitation of the α-cluster. In the Li^6 nucleus the first 0^{+}-level with $T=1$ and $E^{*}(Li^6)=$ 3.5 MeV arises just from the similar excitation of the quasideuteron [8]. Adding ~ 3.5 MeV to the energy $E^{*}(Li^6)$ of the monopole excitation of the α-cluster we find that in Be^6 and He^6 the 0^{+}-level falls within the region of the threshold for the formation of 2 He^3 or $2H^3$ clusters. The reduced width of decay of this state via the He^4+2N channel due to break-up of the α-cluster must be much smaller than the reduced width of decay via the $2He^3$ and $2H^3$ channels. In addition, we note that there are indications of the existance of a positive-parity level in Li^6 at $E^{*}\simeq 15.8$ MeV [9].

Returning to the astrophysical aspect of the developed arguments favouring the existence of new levels in nuclei with A=6, we may assume that the 0^{+}-level of the indicated nature falls within the region of the Gamov peak located in the region of 20 KeV above the threshold of break-up of Be^6 into $2He^3$. In such a case, if the position of the resonance E_r and its total width Γ satisfy the conditions $E_g-\gamma_g \leq E_r \leq E_g+\gamma_g$, $\Gamma \leq 2\gamma_g$ (E_g, γ_g are the position and the half-width of the Gamov peak) one can expect a considerable increase in the reaction rate (1) and, as a consequence, a decrease in the neutrino yield via the channels (2) and (3).

Fig. 1 shows the ratio $\lambda = \frac{\langle\sigma v\rangle_{res.}}{\langle\sigma v\rangle_{nonres.}}$ ($\langle\sigma v\rangle$ is the product of the cross section by the velocity averaged over the Maxwell distribution at the temperature of the Sun $T_\odot = 1.5 \cdot 10^7\,^\circ K$ [2]) between the rates of the resonance and nonresonance reactions as a function of resonance parameters E_r and Γ . To obtain the maximum effect the reduced width of the entrance channel was taken to be equal to the Wigner limit $\frac{3\hbar^2}{2mR^2}$ for the channel radius R=3.4 fm.

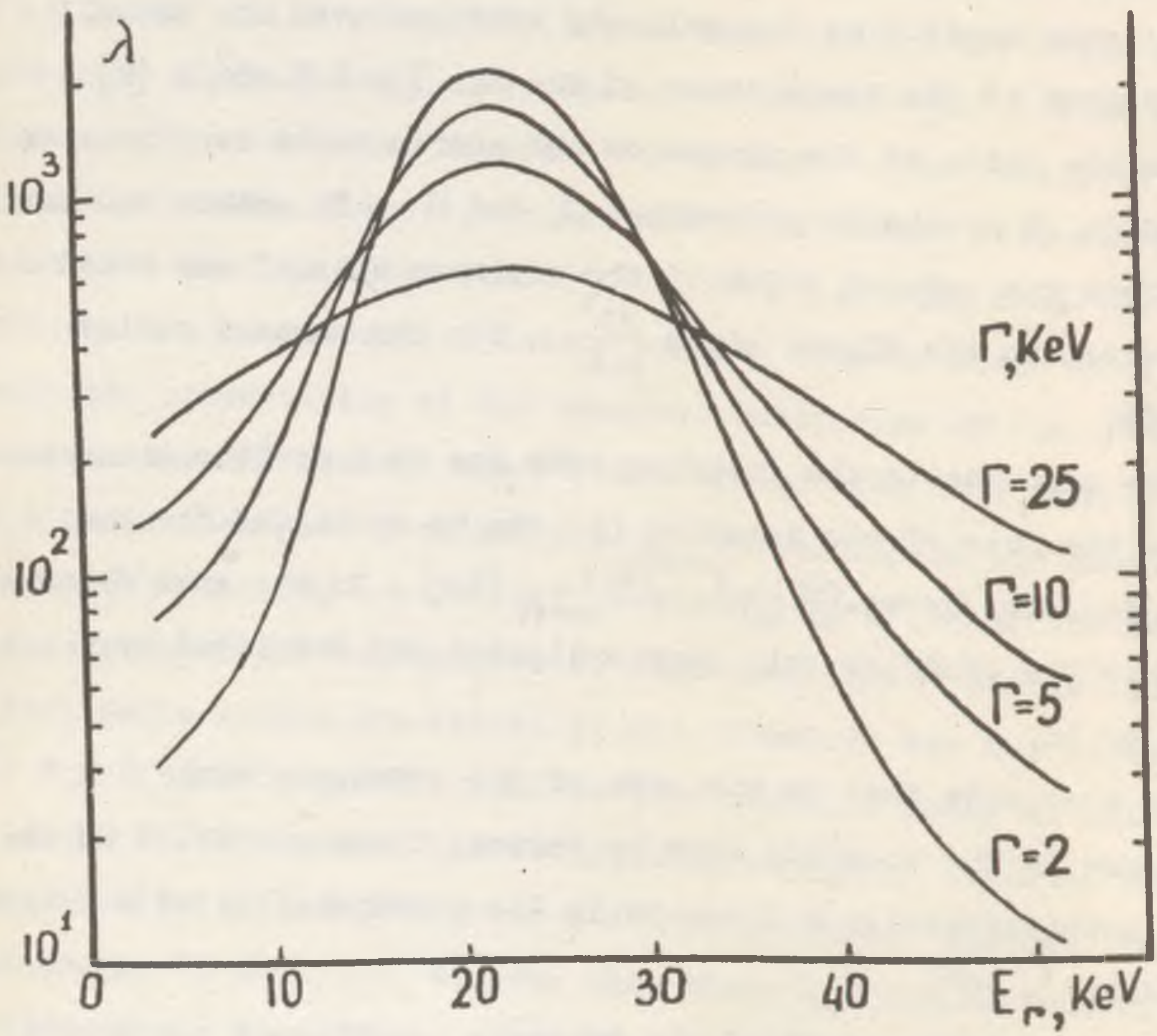
The decrease in the counting rate due to a drastic enhancement in the rate of the reaction (1) can be estimated for the Davis' detector by eq. $(\Phi\sigma)_{res.} \simeq \lambda^{-0.37}(\Phi\sigma)_{nonres.}$ [10] . In the most favourable case the counting rate thus evaluated may decreased by a factor of 16.

We also note that in the case of the resonance with $\ell = 1$ a decrease in the counting rate by several times can still be obtained notwithstanding a decrease in the penetrability of a P-wave by a factor $\sim 10^2$.

One may attempt to find the accurate position of new levels in the nuclei with A=6 for example using inelastic electron scattering by Li^6 or the reactions $He^3(H^3) + He^4 \rightarrow n(p) + Be^6(He^6)$.

The authors wish to thank Drs. R.A.Eramzhyan and Yu.F.Smirnov for valuable discussion of the Li^6 structure.

After this conference professor A.E.Chudakov kindly informed us that in an unpublished paper by Fowler [11] a possible explanation of the Davis' experiments by the influence of a resonance in the He^3+He^3 system was also discussed.



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SOLAR NEUTRINOS: THEORY

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I. INTRODUCTION

The most important fact about the subject I am reviewing is that there is a large discrepancy between calculation and observation, the latter being represented by the Brookhaven experiment Dr. Davis has just described. The origin of this discrepancy is unknown. I have attempted to organize this review so that you can see for yourself what is being tested and how. The first two sections concerning cross sections and solar models are based on a review article by myself and R. L. Sears that will appear in the Annual Review of Astronomy and Astrophysics, Vol. 10. The last two sections include my personal opinion as of May 1972 on the questions of whether or not solar neutrinos reach the earth and on what one should do next to understand the origin of the discrepancy between theory and experiment.

The basic reason for doing solar-neutrino experiments (and the associated theoretical calculations) is to test quantitatively the theories of nuclear energy generation in stars and of stellar evolution. Because of the small photon mean free path in stellar material, photons (which are the subject of conventional astronomy) come to us from the outermost layers of a star. Neutrinos, because of their large mean free paths (>> a stellar radius under normal astronomical conditions), can reach us from the deep interior of a star where the temperature is highest and the nuclear reactions, which are believed to be responsible for energy generation and stellar evolution, occur. We of course know

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more about the sun than any other star because of its proximity and because of the fact that it is in the simplest stage of stellar evolution (the quiescent main sequence phase). The Brookhaven Solar Neutrino experiment of R. Davis, Jr., and his colleagues which is designed to observe solar neutrinos is therefore a critical test of the theories of stellar evolution and nuclear energy generation in stars.

II. CROSS SECTIONS FOR NEUTRINO CAPTURE

The observed counting rate in any neutrino experiment is the integral of neutrino flux per unit energy interval, $\varphi_E\, dE$ (in $cm^{-2}\, sec^{-1}$), times absorption cross section (in cm^2) of the target. The counting rate is $\int \varphi_E \sigma(E) dE$. The calculated absorption cross sections for some of the most frequently discussed targets are listed in Table 1. For the convenience of the user, we follow the standard practice of giving the cross sections averaged over the appropriate neutrino energy spectrum associated with a particular source (e.g., ^8B or ^{13}N decays). Scattering cross sections for the reaction $\nu + e^- \rightarrow \nu' + e^-$ are available as a function of energy in several sources (e.g., Bahcall 1964d, Reines & Kropp 1964, Pontecorvo & Zatsepin 1970).

Because of the surprising results of the ^{37}Cl experiment, the question has arisen as to whether or not the calculated absorption cross sections (Bahcall 1964a,b, 1966) for this experiment could be the cause of the discrepancy between theory and observation. In our opinion, this is not possible. The cross sections for the neutrinos from $p + e + p \rightarrow {}^{2}H + \nu$, ${}^{7}Be$, ${}^{13}N$, and ${}^{15}O$ all depend only on the ground state transition from ${}^{37}Cl$ to ${}^{37}A$. This ground state transition is the inverse of the well-studied laboratory decay of ${}^{37}A$ (i.e., $e^{-} + {}^{37}A \rightarrow {}^{37}Cl + \nu$) and a negligible error (< 10%) is introduced by the translation of the laboratory experiments into neutrino absorption cross sections using the established theory of nuclear beta decay (e.g., Konopinski 1966).

The evaluation by Bahcall (1964a,b) of transitions from the ground state of ${}^{37}Cl$ to various excited states of ${}^{37}A$ increased the calculated cross section for this reaction by a factor of eighteen. This theoretical step caused considerable discussion when it was first made but subsequent experimental studies have confirmed the accuracy of the original calculation. In fact, the discovery (by Hardy & Verrall 1964, Reeder, Poskanzer & Esterlund 1964) of the isotope ${}^{37}_{20}Ca_{17}$ (whose existence and decay lifetime were predicted by Bahcall 1964a,b) has made possible the calculation of the cross sections for ^{37}Cl neutrino absorption from essentially experimental data (see Bahcall 1966). The essential point to recognize is that $^{37}_{20}Ca_{17}$ is the mirror nucleus to $^{37}_{17}Cl_{20}$, i.e., the nuclear structure of ^{37}Ca is the same as that of ^{37}Cl because of the well established symmetry between neutrons and protons in nuclear physics. Beta decays from ^{37}Ca to the excited states of $^{37}_{19}K_{18}$, the mirror nucleus of $^{37}_{18}A_{19}$, are analogous to neutrino-induced transitions from ^{37}Cl to the excited states of ^{37}A. Thus, the measurements (Poskanzer, McPherson, Esterlund & Reeder 1966) of the decay rates of $^{37}Ca \rightarrow {}^{37}K + e^- + \nu$ allow one to largely determine experimentally the required ^{37}Cl absorption cross sections. It should be noted also that about five percent of the total ^{8}B cross section arises from the ground state transition between ^{37}Cl and ^{37}A and about sixty-five percent from a superallowed transition whose absolute value is known on the basis of general nuclear physics considerations (the precise location of the analogue level, which is difficult to calculate or measure to better than two hundred kev, is unimportant since the cross section must be averaged over the neutrino spectrum of ^{8}B which is 14 MeV broad). The remaining twenty-five percent is parcelled out among various excited state transitions; the parcelling is constrained experimentally by a number of measured parameters including the measured (Poskanzer et al 1966) branching ratios in the ^{37}Ca decay, the spins and parities of the ^{37}K states (Goosman & Kavanagh 1967), and the observed total ^{37}Ca and ^{37}A decay lifetimes. Bahcall (1966) estimated an overall uncertainty of the order of ten percent in the calculated cross section for ^{8}B neutrino absorption by ^{37}Cl and this estimate still seems reasonable.

III. SOLAR MODELS

III.1. CONSTRUCTION OF STANDARD SOLAR MODELS

One does not need a detailed model of the internal constitution of the sun to predict the total number of solar neutrinos reaching the earth per cm^2 per sec. In fact, this total number is $2\ L_\odot/(25\ \mathrm{MeV} \cdot 4\pi\ (1\ \mathrm{au})^2) \sim 10^{11}\ \mathrm{cm}^{-2}\ \mathrm{sec}^{-1}$, since on the average the basic fusion reaction $4p \rightarrow \alpha + 2\ e^+ + 2\nu$ releases ~ 25 MeV. However, the proposed methods of detection record only a portion of the neutrino energy spectrum, which ranges from zero to more than 10 MeV, and the cross section for detection depends strongly on the neutrino energy. Moreover the form of the neutrino energy spectrum is a sensitive function of temperature, density, and composition (e.g., the production rate of the important high-energy neutrinos from ^8B decay, first discussed by Fowler (1958) and Cameron (1958), varies roughly as T^{13}). Thus in the past ten years a considerable effort has gone into the construction of accurate solar models and the computation of the neutrino energy spectra they imply.

Table 2 illustrates the main process currently believed to convert H to He in the sun, the proton-proton chain. The form of the neutrino energy spectrum depends on the relative frequencies of the three branches by which the chain is completed and a detailed solar model is required to determine the branching ratios.

We shall briefly describe the construction of solar models because the present disagreement between calculations and observation may depend on one or more assumptions or parameters of the models. The standard theory of main-sequence stars (Schwarzschild 1958, Cox & Giuli 1968, Chiu 1968) assumes, at each point in the star, (i) hydrostatic equilibrium between gravitational force and pressure gradient (radial, for a spherically symmetric star), (ii) energy transport by radiation and convection, and (iii) energy production by hydrogen burning. These assumptions may be stated as four first-order ordinary differential equations, together with appropriate boundary conditions and constitutive relations. The latter include (i) the equation of state which connects pressure, density, and temperature; (ii) the radiative opacity (which is a function of chemical composition); and (iii) the nuclear energy-generation rate, involving the cross sections for the several reactions (see Clayton 1968 for a clear and a thorough treatment of

the nuclear reaction processes). As described by Schwarzschild (1958, pp. 96-8) the only data required to obtain a solar model are the total mass and the distribution of chemical composition throughout the star. We forego a description of the standard numerical techniques (see, e.g., Sears & Brownlee 1965, Kippenhahn, Weigert & Hofmeister 1967); the results that emerge from the computer include the march of physical variables throughout the star, the total radius and luminosity (ergs per second), the distribution of the energy-production processes, and the central temperature. The calculated neutrino spectrum is determined by numerical integration over mass shells.

In the case of the sun the total mass is accurately known (1.989×10^{33} g) but the distribution of chemical composition is not known a priori. The conventional assumption in most recent work has been to adopt a homogeneous initial composition and then to construct an evolutionary sequence of models (Schwarzschild 1958, pp. 98-100). In any one evolved model the helium/hydrogen ratio is greatest at the center and falls off to the primordial value some distance away from the center, with the usual assumption of no mixing. A further datum required, then, is the age of the sun, usually taken as about 4.7×10^9 years (see Bahcall & Shaviv 1968 for references). One constructs an evolutionary sequence of a half-dozen or so models, ending with "the" solar model, which is required to have the present-day luminosity after 4.7×10^9 years. In the case of the sun a strict requirement on the computed radius is not useful because the calculated value depends on the uncertain structure of the convective envelope; fortunately this uncertainty does not significantly affect the calculated deep interior structure (Schwarzschild 1958, pp. 144-5; Sears 1964).

The question of which initial homogeneous composition to adopt is a critical one. The composition is conventionally described by X, Y, and Z, respectively the mass fractions of hydrogen, helium, and heavier elements (Z = 1-X-Y); thus only two free parameters appear. Once values of the composition parameters are adopted, a solar model can be obtained. If the luminosity of the model differs from the observed luminosity then one usually changes slightly one of the composition parameters--because the calculated luminosity is sensitive to Z via the opacity (roughly, $L \sim Z^{-1}$) and is sensitive to X and Y via the opacity and the mean atomic weight per free particle in the perfect-gas equation of state (roughly, $L \sim \mu^{7.5}$). Hence one approach (Sears 1964, 1966, Bahcall & Shaviv 1968, Bahcall, Bahcall & Ulrich 1969) has been to adopt the value of Z/X given by spectroscopic observations of the solar photosphere and to pick successive values of Y, computing an evolutionary sequence for each, until a sequence is obtained in which the final model reproduces the observed luminosity. Note that this approach yields Y, the primordial helium content of the sun, as a theoretical result which can be compared with chromospheric determinations, solar cosmic rays, and the primordial abundance from a variety of cosmologies. Another approach, advocated by Iben (1968, 1969a), has been to adopt ab initio a value of Y that one believes in on the basis of other astronomical evidence and to vary Z until one obtains a luminosity fit.

III. 2 RESULTS OF STANDARD SOLAR MODELS

III. 2.1 Review of Recent Models

The earliest detailed solar model used to calculate solar neutrino fluxes was described by Bahcall, Fowler, Iben & Sears (1963). Another early model was used by Pochoda & Reeves (1964) to estimate a neutrino energy spectrum. The first systematic study of solar models for the purpose of predicting neutrino fluxes and their uncertainties was that of Sears (1964). The basic parameters of his preferred Model J were an initial composition parameter $Z/X = 0.028$; the mass (1.989×10^{33} g); and the age of 4.5×10^9 years. The equation of state was that of an ideal gas, including electron-degeneracy pressure; the opacity was based on the tables of Keller & Meyerott (1955); and the energy-generation rate [including the small gravitational contraction term (Schwarzschild

1958, Eq. 12.10)] involved the pp-chain and the CN cycle, with the cross-section factors in the former from Parker, Bahcall & Fowler (1964). As described in the preceding section the model was fitted to the observed solar luminosity, 3.90×10^{33} erg sec^{-1}. The resulting initial helium content was $Y = 0.27$ and the resulting ^8B neutrino flux was $\varphi(^8B) = 1.9 \times 10^7$ cm^{-2} sec^{-1}. The principal parameters of this (1964) model are listed in Table 3. Bahcall (1964b)adopted an average over several of Sears's models which gave $\varphi(^8B) = 2.5 \times 10^7$, and this, plus a small contribution from the ^7Be electron-capture neutrinos, led to a predicted number of neutrino captures at the earth of $\Sigma(\varphi\sigma) = 36$ SNU.

Solar models obtained by Ezer & Cameron (1965) and by Weymann & Sears (1965) contained minor improvements, primarily in the use of Los Alamos opacities (Cox & Stewart 1966, Cox, Stewart & Eilers 1966). Since these are lower than the Keller-Meyerott opacities in the deep interior the resulting helium contents and neutrino fluxes were slightly reduced compared to those of Sears (1964). Bahcall (1966) combined fluxes he calculated with the aid of the improved solar models and new experimental information on the mass-37 system and on the $^7Be(p,\gamma)\ ^8B$

cross section to find a neutrino capture rate of $\Sigma(\varphi\sigma) = 30$ SNU.

The next major development was a substantial increase (a factor of five) in the reported low-energy experimental cross-section factor for the ^3He (^3He, 2p) ^4He reaction in the period 1965-67 (Dwarakanath & Winkler 1971, Bacher & Tombrello 1968; see Tombrello 1967, Barnes 1971, and Kavanagh 1972 for reviews). This enhanced the branching through termination I of the pp-chain (see Table 2) and lowered the predicted ^8B neutrino flux; the first models incorporating the new experimental results were reported by Shaviv, Bahcall & Fowler (1967). General uncertainty in the low-energy cross sections at that time led Bahcall, Bahcall, Fowler & Shaviv (1968) to construct two models based on assumed extreme values for four important cross-section factors in the pp-chain. A preferred model, incorporating the best parameters then available, was constructed by Bahcall & Shaviv (1968), who found $\Sigma(\varphi\sigma) = 22$ SNU. They also discussed the effects of several parameter variations.

Shortly afterwards, in early 1968, two new developments stimulated further model building. A redetermination of the half life of the free neutron by Christensen et al (1967, 1971) and a refinement in the calculation of various nuclear parameters led to an increase in the cross-section factor of the proton-proton reaction (Bahcall & May 1968, 1969). A redetermination of the solar photospheric chemical composition, by Lambert (1967a,b, 1968), Lambert & Warner (1968a,b,c), and Warner (1968), gave a new lower value of $Z/X = 0.019$, thus suggesting a decrease in the solar interior opacity. Both of these developments implied a further reduction in the predicted neutrino flux and new models were obtained by Bahcall, Bahcall & Shaviv (1968) (see also Torres-Peimbert, Ulrich & Simpson 1969). At the same time the first results from the Brookhaven experiment became available: Davis, Harmer & Hoffman (1968) found an observational upper limit to the solar neutrino capture rate of $\Sigma(\varphi\sigma) = 3$ SNU. This upper limit was less than that predicted by the most probable model of Bahcall, Bahcall & Shaviv (1968): $\Sigma(\varphi\sigma) = 7.5$ SNU.

The model work of the latter authors was refined by Bahcall, Bahcall & Ulrich (1969), who incorporated revised values for the $^{7}Be(p,\gamma)\,^{8}B$ cross section and for the $^{7}Be(e^{-},\nu)\,^{7}Li$ capture rate, both of which tended to lower the predicted $\varphi(^8B)$ value. Their (1969) model (see Table 3) gave a capture rate of $\Sigma(\varphi\sigma)$ = 6 SNU, only a factor of two above the observational limit. Their paper gave an extensive discussion of parameter variations, including effects of perturbations on the oapcity and on the equation of state.

Another discussion of the first results of Davis et al was given by Iben (1968, 1969a) who constructed a large number of models to illustrate the parameter dependence of neutrino fluxes. He utilized Y, the primordial helium content, as the composition parameter rather than Z/X, the observed heavy-element-to-hydrogen ratio derived from the photosphere. Iben emphasized that consistency with the experimental results could almost be achieved with values of Y considerably below those determined in other objects in the galaxy. This conclusion is still possible at the present time although later developments have modified the arguments in Iben (1969a). His paper is particularly useful because of the extensive illustrations and physical descriptions of the effects of parameter changes.

The next two developments were a substantial increase in the Fe abundance in the photosphere from revised f-value determinations (Garz, Holweger, Kock & Richter 1969, Martinez-Garcia, Whaling, Mickey & Lawrence 1971, Wolnik, Berthel & Wares 1971; cf. Ross 1970, Cowley 1970, 1971) and the inclusion of electron-correlation effects in the opacity (Diesendorf 1970). Watson (1969, 1970) pointed out that an increase by a factor of 10 in the iron number abundance would lead to a substantial (~ 30%) increase in the opacity of the solar core and thus an increase in the predicted ^{8}B neutrino flux. The electron-correlation effects, as Diesendorf noted, tended to decrease the opacity. A new model incorporating these effects was computed by Bahcall & Ulrich (1971); see Table 3. At the same time new results from the Brookhaven experiment became available (Davis, Rogers & Radeka 1971): a positive detection of ^{37}A, $\Sigma(\varphi\sigma) = 1.5 \pm 1.0$ SNU was suggested. The model prediction by Bahcall & Ulrich (1971), $\Sigma(\varphi\sigma) = 9.0$ SNU, was greater by a factor of six; similar results were reported by Abraham & Iben (1971) and by Ezer & Cameron (1971).

III.2.2 Dependence of Predicted Neutrino Fluxes on Parameter Values

As noted above, the effects of parameter changes have been explored in many papers. We shall not give a detailed discussion, partly to avoid repetition and partly because of the possibility that a qualitatively new development may change the present situation.

The dependence of the predicted counting rate in the ^{37}Cl experiment on most of the relevant parameters is succinctly illustrated by the relation

$$\Sigma(\phi\sigma) = 1.35 \times 10^{-36}\ \mathrm{sec}^{-1}\ \mathrm{per}\ {}^{37}\mathrm{Cl\ atom}$$
$$\times \left(\frac{S_{11}}{S_{11}(0)}\right)^{-2.5} \left(\frac{S_{33}}{S_{33}(0)}\right)^{-0.37} \left(\frac{S_{34}}{S_{34}(0)}\right)^{+0.8}$$
$$\times \left[1 + 3.47 \left(\frac{S_{17}}{S_{17}(0)}\right)^{+1.0} \left(\frac{\lambda_{e7}}{\lambda_{e7}(0)}\right)^{-1.0}\right] \tag{1}$$
$$\times \left(\frac{\mathrm{Age}}{4.7 \times 10^{9}\,\mathrm{yr}}\right)^{+1.4} \left(\frac{Z}{0.015}\right)^{+1.1} \left(\frac{L_\odot}{3.83 \times 10^{33}\ \mathrm{erg\ sec}^{-1}}\right)^{+7.3}$$

The denominators are the parameter values used by Bahcall, Bahcall & Ulrich (1969) (see Table 3); $\lambda_{e7}(0)$ is the ^{7}Be electron-capture rate from Bahcall & Moeller (1969). Equation(1) was derived numerically by Bahcall, Bahcall & Ulrich (1969) from a comparison of many models. Setting the ratios equal to unity gives their standard value, 6.0 SNU (Table 3) in 1969. The signs of the exponents give the directions of the dependences which are not obvious in all cases (usually because of the strong influence of the boundary condition that the calculated present luminosity equal the observed luminosity; see Section III.1). More detailed discussions of individual dependences are given in a number of papers (Sears 1964, Bahcall, Cooper & Demarque 1967, Shaviv, Bahcall & Fowler 1967, Bahcall, Bahcall, Fowler & Shaviv 1968, Bahcall & Shaviv 1968, Bahcall, Bahcall & Shaviv 1968, Iben 1968, 1969a, Bahcall, Bahcall & Ulrich 1969, Bahcall & Ulrich 1971). Most of these papers have graphs illustrating the dependences.

The various factors in Equation (1) enable one to find the individual effect of each of the input parameters on the predicted counting rate. For example, the effect of the five-fold increase in S_{33}, the $^3He(^3He, 2p)^4He$ cross-section factor, was to reduce the 1964 counting rate by a factor of 0.57. Other effects were, in order of decreasing importance, the decrease in Z/X (0.68), the increase in S_{11} (0.75), the decrease in $L_\odot$ (0.84) and the increase in ^7Be electron-capture rate due to bound electron-capture (Iben, Kalata & Schwartz 1967, Bahcall & Moeller 1969) (0.87). The increase in assumed age led to an increase in predicted counting rate by 1.06; the increase in S_{34}, by 1.23. Multiplying all the factors together gives a total reduction of 0.3 which is the ratio of $\Sigma(\varphi\sigma)$ (1971) to $\Sigma(\varphi\sigma)$ (1964) in Table 3. The calculation from Equation (1) does not include refinements in the equation of state, which are minor, or changes in opacity. The gross dependence of opacity on Z is of course included, but the specific increase in opacity (hence in $\varphi(^8B)$) due to the increased iron abundance (Watson 1970) is not; evidently this is nearly balanced by the decrease in interior opacities between the Keller-Meyerott and Cox-Watson tables and by the electron-correlation effects. Incidentally, we note that Bahcall & Ulrich (1971) have shown that solar-model calculations by different workers with the same parameters give the same results to an accuracy of $\sim \pm 10$ percent.

III. 2. 2 Most Recent Results

The most recent calculations by Bahcall, Huebner, Magee, Merts, and Ulrich (1972) have used revised Los Alamos opacities. The results (not yet in final form) with all other paremters having their standard values [Bahcall and Ulrich 1971] are: $\Sigma(\varphi\sigma) \simeq 6$ SNU, with 4.2 SNU's coming from ^8B neutrinos, 0.83 from ^7Be neutrinos, 0.26 from pep, and 0.26 from ^{13}N and ^{15}O. It appears that the earlier MIT opacity corrections were invalid. When meson exchange effects on the p-p reaction are included [M. Gari and A. H. Huffman (1972); D. O. Riska (1972)], S_{11} (meson exchange) $\cong (1.09)S_{11}$ (Bahcall-May 1969)), we find $\Sigma(\varphi\sigma) = 4.35$ with 3.15 coming from ^8B neutrinos, 0.74 from ^7Be neutrinos, 0.26 from pep, and 0.2 from ^{13}N and ^{15}O. These latter calculations take into account the revised Los Alamos opacities and estimated meson exchange effects for the p-p reaction.

We conclude that refinements in the input data for standard solar models are not likely to eliminate the present large discrepancy between calculated and observed neutrino fluxes. Most notable in the improvements of the past decade has been the remeasurement of the various pp-chain cross sections--by now each crucial reaction has been remeasured (or recalculated) at least twice by different groups with good agreement throughout [see the excellent reviews by Tombrello (1967) and Kavanagh (1972)].

III. 3 Non-Standard Solar Models

The preceding discussion has been concerned with standard solar models. Since the first experimental results of Davis et al. (1968) revealed a discrepancy with standard theory, many suggestions have been offered and debated. We list some of these suggestions below.

The possibility of mixing, which could maintain a high amount of hydrogen in the center and thus lower the central temperature compared to the standard models, was first suggested by Ezer & Cameron (1968), who found that complete mixing over the entire lifetime of the sun could reduce $\varphi(^8B)$ to one-fourth of the standard result (see also Shaviv & Beaudet 1968, Iben 1968). Various arguments have been given against the likelihood of maintaining such extreme mixing [see Bahcall, Bahcall & Ulrich (1968) and Shaviv & Salpeter (1968, 1971)]. Also Ulrich (1969) showed that a rapidly rotating solar core had little effect on the calculated neutrino fluxes. Iben (1969b, c) investigated the possibility that the sun had a convective core and concluded this was unlikely to significantly change the calculated fluxes.

The effect of a large-scale interior solar magnetic field has been explored by Bahcall & Ulrich (1971) [cf. Iben (1968), Abraham and Iben (1971)], who find (taking the primordial Z/X from observation) that the predicted neutrino flux increases.

Diffusion of heavier elements toward the solar center (Aller & Chapman 1960) would tend to increase the predicted neutrino flux. Schatzman (1969, 1970) has suggested turbulent diffusion of ^{3}He, but the sign of his effect has been questioned by Shaviv & Salpeter (1971) (who showed it was small) and by Bahcall & Ulrich (1971).

Kocharov & Starbunov (1970, a, b), pointed out that if the sun presently contains enough ^{3}He so that ^{3}He(^{3}He, 2p) ^{4}He is the major energy source, the central temperature would be too low to produce solar neutrinos detectable in the ^{37}Cl experiment. Abraham & Iben (1970) have constructed models to check this suggestion and find that the initial mass abundance required is $X_{^3He} = 0.09$, which they note is substantially larger than the values found in H II regions, in meteorites, and in cosmological models (see also Ulrich 1971).

The possibility of a secular (thermal) instability, in which the observed luminosity might not equal the energy production rate on a Kelvin-contraction or photon-diffusion time scale, has not been thoroughly explored. An investigation by Aizenman & Perdang (1971) of a slightly more massive evolved model did not reveal any growing modes of thermal instability. Most recently Harm and Schwarzschild (1972) have shown that there are no thermally unstable radial modes with periods greater than 10 years and Littleton et al. (1972) have shown that longitudinal waves are unimportant for energy transport.

The possibility that the gravitational constant G might decrease on a cosmic time scale has been incorporated into solar models by several workers; Ezer & Cameron (1966) found that a varying-G model (Brans & Dicke 1961) would have four times the $\varphi(^{8}B)$ of a standard model (see also Shaviv & Bahcall 1969). A model by Rouse (1969) implied that the CN cycle dominates in the sun; this was already ruled out on the basis of the result of Davis et al. (1968) by Bahcall, Bahcall & Shaviv (1968) (provided the neutrinos reach the earth). It has been suggested that quark catalysts might be important in the sun (Libby & Thomas 1969; cf. Salpeter 1970).

None of the above suggestions has been widely accepted.

IV. IMPLICATIONS OF THE OBSERVATIONS

The implications of the ^{37}Cl experiment are best appreciated when one distinguishes the various counting levels at which different things are being tested. If the entire energy source for the sun were the CN cycle, then the counting rate in the ^{37}Cl experiment would be 35 SNU independent of model parameters (Bahcall 1966). Thus at the level of tens of SNU's one is distinguishing between the principal energy source mechanisms (pp-chain or CNO cycle). Modern solar model calculations all suggest that the principal energy source for the sun is the pp-chain and that, with presently accepted parameters, the expected counting rate should be ~ 4-9 SNU (Bahcall & Ulrich 1971, Abraham & Iben 1971, Bahcall, Heubner, Magee, Merts, & Ulrich 1972). Most of the calculated counting rate with present parameters is due to the rare ^{8}B reactions (see Table 2) that have no effect on the structure of the sun. On the other hand, the basic terminations, pp I and pp II (see Table 2), contribute (along with a small amount ~ 0.3 SNU from the CN neutrinos) a counting rate ~ 1 to 1.5 SNU which is insensitive to most parameter variations (Bahcall, Bahcall & Ulrich 1969). Thus a counting rate below 1 SNU would be in conflict with the basic ideas of stellar evolution as described in the standard books of, for example, Schwarzschild (1958) and Clayton (1968). The basic idea that nuclear fusion among light elements is the energy source for main sequence stars guarantees (assuming ν_e flux conservation) a counting rate of 0.3 SNU (Bahcall, Bahcall

& Shaviv 1968). The various levels of meaning of the ^{37}Cl experiment are summarized in Table 5.

We recall that Davis et al. (1968) have set an upper limit to the solar-induced counting rate of 3 SNU (see equation (1)). Comparison of this experimental limit with Table 5 leads to the following conclusions: (1) the CNO cycle is responsible for less than three percent of the total luminosity of the sun; (2) there is a discrepancy between present model calculations and observations of the order of five and (3) the ^{8}B flux from the sun is $<10^{6}$ cm^{-2} sec^{-1}.

Three possible astrophysical explanations for the discrepancy have been most strongly advocated in recent years: (1) the ^{8}B production cross section is seriously in error; (2) neutrino oscillations occur; and (3) the solar interior value of the heavy element abundance, Z, is much less than the surface value of Z.

The experimental low-energy cross section factor for the critical reaction ^{7}Be(p, γ)^{8}B is plotted in Figure 3 as a function of incident proton energy. The data shown are from the experiment of Kavanagh, Tombrello, Mosher & Goosman (1969) and we are greatly indebted to these authors for permission to display their as-yet unpublished data. The data shown are in good agreement with the earlier results of Kavanagh (1960) and Parker (1966a, b, 1968) (cf. Tombrello 1965 and Aurdal 1970); the higher-energy data of Vaughn, Chalmers, Kohler & Chase (1971) are systematically lower than the results of Kavanagh et al. (1969) by about twenty-five percent. Nevertheless it seems from Fig. 2 highly unlikely that the extrapolated low-energy cross section factor for the $^7Be(p, \gamma)^8B$ reaction could be off by the factor of ten or more required to fit the observations.

The possibility that a tiny amount of lepton nonconservation would lead to oscillations on the scale of the solar system between electron neutrinos and muon neutrinos and a consequent decrease in the observable electron-neutrino flux at the earth has been proposed by Pontecorvo (1968) and Gribov & Pontecorvo (1969) [see also Pontecorvo 1958]; Pontecorvo & Zatsepin 1970 and Kocharov & Fergberg 1970]. However this hypothesis leads to a most probable decrease in counting rate due to the ^8B neutrinos of only a factor of two when averaged over neutrino energies (Bahcall & Frautschi 1969). For a special value for the dimensionless ratio of various parameters, one can (see Fig. 1 of Bahcall & Frautschi 1969) lower the calculated ^8B flux by a factor of five or more (and the flux of lower energy electron neutrinos by a factor of two) but there is no theoretical justification for this a priori unlikely value.

Bahcall & Ulrich (1971) have pointed out that if the primordial heavy element abundance in the solar interior satisfied $Z_{interior} \lesssim 10^{-3}$ (we have today $Z_{surface} \sim 2 \times 10^{-2}$), then the interior opacity is significantly reduced and the calculated counting rate is only ~ 1.5 SNU [meson-exchange corrections to the p-p rate may reduce this slightly].

The calculated primordial helium abundance is, however, uncomfortably small ($Y = 0.1$). Nevertheless it is important to investigate further whether or not it can be excluded that in the early stages of the formation of the sun the heavy elements were largely expelled from the interior (by radiation pressure or by collection in some dust cloud).

V. DO SOLAR NEUTRINOS REACH THE EARTH?

In order to reach the earth, solar neutrinos must travel $\sim 10^{13}$ cm ($\equiv$ 1 A.U. = earth-sun distance) and traverse $\simeq 10^{36}$ electrons (or protons) per cm^{-2} $\left[\int_0^{R_{sun}} n\, d\ell \sim 10^{36}\ \text{cm}^{-2}\right]$. It is at least conceivable that solar neutrinos do not reach the earth. If this were known to be true, the upper limits of the Davis experiment would have to be interpreted as constraints on the physics of the electron's neutrino, not the solar interior.

There are three processes by which solar neutrinos might be prevented from reaching the earth: absorption, scattering, or decay. As long as we restrict ourselves to particles that have already been discovered, none of these processes are possible with the required cross sections ($\sim 10^{-36}$ cm^2 at 1 MeV) or half life (< 500 sec.) An amusing possibility which I have considered is strongly forward-peaked elastic scattering, so forward-peaked that most of the interactions are not detected by experimentalists who work at appreciable momentum transfers. Even this is not nearly enough given the experimental constraints [Cundy et al. (1970); Reines et al. (1960)], as one can easily show using, for example, an assumed diffraction scattering law $f(q) \propto \exp - bq^2$.

We have recently suggested [Bahcall, Cabibbo, & Yahil 1972] that electron neutrinos of 10 MeV energy may be unstable with a half life $< c^{-1}$ (the distance between the earth and the sun) ~ 500 seconds. The various possibilities were discussed in our original paper, but a general requirement of all the proposals is that they require a previously unknown interaction or particle. One may also consider the possibility (Regge 1971) that there are particles present in the sun (but not yet observed on earth) which interact strongly with neutrinos but only very weakly with ordinary matter. The astrophysical consequences of this latter suggestion are somewhat model dependent so I will restrict myself to a discussion of neutrino decay.

There are three direct astrophysical consequences of the hypothesis that the instability of ν_e is responsible for the unexpectedly low counting rate in the ^{37}Cl experiment. These consequences are: (1) the ^{37}Cl experiment should yield a result consistent with $\Sigma(\varphi\sigma) \equiv 0$; (2) attempts to detect the astrophysically-guaranteed pep neutrinos [e.g., by the proposed ^{7}Li experiment--see Bahcall (1969)] should fail; and (3) all astronomy with ν_e will be impossible (i.e., electron neutrinos from black hole formation or supernova collapse will also not reach us).

The basis for these predictions is the remark that if neutrino-instability is the explanation of Davis' result then the terrestrially observed flux of solar neutrinos with energy E, φ(E_, satisfies

$$\varphi_{\text{earth}}(E) = \varphi_{\text{sun}}(E)\,[\exp{-\alpha(E)}]\left(\frac{R_\odot}{1\,\text{A.U.}}\right)^2 ,$$

where $\alpha(E) = [(\text{earth-sun distance})/c\,\tau_{\frac{1}{2}}(10\text{ MeV})]\,(\frac{10\text{ MeV}}{E})$ and $\alpha(10\text{ MeV}) \geq 2$ because of the absence of ^8B neutrinos. The energy dependence of $\alpha(E)$ is due to time-dilation; low-energy particles decay more rapidly than their high-energy counterparts. If the higher energy ^8B neutrinos decay, then the lower energy ^7Be, ^{13}N, ^{15}O, and pep neutrinos must certainly do so. The ^{37}Cl experiment of Davis probably has an ultimate sensitivity ~ 0.5 SNU compared to a predicted counting rate ~ 5-10 SNU. It is extremely unlikely that the parameters for ν_e decay are just such that most of the neutrinos decay in 1 A.U. but enough are left over to permit a detectable counting rate in the ^{37}Cl experiment. Of course if neutrinos from the sun do not reach us, then ν_e's from much more distant astronomical sources will also decay along the way (the anticipated energies are not very different from those of solar neutrinos but the distances are more than a million times larger).

In the absence of experimental information, I would like to summarize the arguments, pro and con, as to whether or not solar neutrinos reach the earth. There are two "arguments" supporting the suggestion that they do not reach the earth: (1) the suggestion "solves" the solar-neutrino problem; (2) there have been other surprises in weak and superweak interactions so why not this one. The arguments against decay are: (1) a new interaction or particle is required; and (2) the idea is suggested by astrophysical considerations (and no astrophysicist knows enough about the systems he investigates to plausibly suggest modifications in the conventional form of physical laws). This last argument, which is based on the complexity of astronomical systems, is frequently made by theoretical physicists and seems to me to have a lot of merit. We recall, however, that there are rare but important exceptions, e.g.: Newton's discovery of the law of gravitation from the motion of the planets; the inference by Roemer of the finite velocity of light from the motion of the moon's of Jupiter; the discovery of helium in the sun; and the precession of the perihelion of Mercury and its connection with the development of the theory of General Relativity.

VI. WHAT NEXT?

I am frequently asked where I believe the trouble lies, in the physics or the astrophysics. My answer is: I don't know and don't even have a strong opinion. It seems to me that more theoretical work on possible instabilities in the solar interior, on opacities, and on other fundamental assumptions and ingredients of the theory of stellar evolution is required before we can be "sure" that something is not seriously wrong with our understanding of the sun. The theory of stellar evolution does not have an impressive record of successful predictions; it seems to me we are entitled to question its foundations (see Bahcall 1971).

On the other hand, I think it is also conceivable that solar neutrinos do not reach the earth. A possibility worth keeping in mind is the idea that the solar-neutrino puzzle might be linked to one of the other current mysteries in particle physics: CP nonconservation, the absence of $K_L \rightarrow \mu^+ + \mu^-$, or the problem of constructing the correct Weinberg theory. I hope that high-energy physicists will at least consider this possibility. Further work on mesonic corrections to the p-p reaction are also needed.

The ultimate answer must of course come from experiment. The flux at the sun of pep neutrinos [from $p + e + p \rightarrow {}^2H + \nu_e$] is about as well known as the solar luminosity in photons. Experiments designed to detect pep neutrinos would distinguish between the two presently co-existing but antipodal hypotheses that the origin of the discrepancy is in our inadequate astrophysical theory or in our inadequate theory of leptons. If pep neutrinos are not observed at the predicted flux level, then we well know that the problem must lie in our lack of understanding of the neutrino. If pep neutrinos are detected, then we will have to find some fundamental correction to the theory of stellar evolution in its simplest application.

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TABLE 1. Total cross sections for neutrino capture[a]

Neutrino Emitter	Target ($10^{-46} cm^2$)				
	^{37}Cl	^{7}Li	^{87}Rb	^{71}Ga	^{2}H
p+p	0	0	7 E+1	3 E+1	0
$p+e^-+p$	1.7 E+1	5.9 E+2	5.5 E+2	6 E+1	0
^{7}Be	2.9[b]	1.6 E+1[b]	2.5 E+2	1.5 E+2	0
^{8}B	1.35 E+4	4.5 E+4	1 E+4	7 E+3	1.2 E+4
^{13}N	2.1	4.5 E+1	2 E+2	1 E+2	0
^{15}O	7.8	2.3 E+2	3.5 E+2	2 E+2	0.1

[a] The notation E+X means 10^{+X}. The ^{37}Cl cross sections are from Bahcall (1964a,b, 1966) and the ^{7}Li cross sections are from Bahcall (1969a) with the exception of the p-e-p source which was calculated for this paper following Domogatsky (1969). The ^{87}Rb cross sections are from Bahcall (1964c), the ^{71}Ga cross sections from Bahcall (1967), and the ^{2}H cross sections from Ellis & Bahcall (1968). See also Sunyar & Goldhaber (1960), Kuzmin & Zatsepin (1965, 1969), Domogatsky, Gavrin & Eramjan (1965), Kelly & Uberall (1966), Zumin (1966), and Kopysov & Kuzmin (1967, 1968). An asymmetry parameter, α, that gives the directional sensitivity possible in a given experiment if the direction of the produced electrons is measured can easily be calculated (see Bahcall 1964c, 1967), given specific assumptions regarding the incident neutrino energy spectrum.

[b] Does not include the fact that only ninety percent of the ^{7}Be reactions produce neutrinos above threshold for capture.

TABLE 2. The proton-proton chain[a]

	Reaction	Percentage	Maximum Neutrino Energy (MeV)
	$p+p \rightarrow {}^{2}H+e^{+}+\nu$	(99.75%)	0.420
	or		
	$p+e^{-}+p \rightarrow {}^{2}H+\nu$	(0.25%)	1.44 (monoenergetic)
	${}^{2}H+p \rightarrow {}^{3}He+\gamma$		
I.	${}^{3}He+{}^{3}He \rightarrow {}^{4}He+2p$	(86%)	
	or		
	${}^{3}He+{}^{4}He \rightarrow {}^{7}Be+\gamma$		
II.	${}^{7}Be+e^{-} \rightarrow {}^{7}Li+\nu$		0.861 (90%), 0.383 (10%) (both monoenergetic)
	${}^{7}Li+p \rightarrow 2\,{}^{4}He$	(14%)	
	or		
III.	${}^{7}Be+p \rightarrow {}^{8}B+\gamma$		
	${}^{8}B \rightarrow {}^{8}Be^{*}+e^{+}+\nu$		14.06
	${}^{8}Be^{*} \rightarrow 2\,{}^{4}He$	(0.02%)	

[a]The numbers in parentheses are the percentages at which various reactions are calculated to occur from the standard model of Bahcall & Ulrich (1971).

TABLE 3. Properties of some recent solar models[a]

	1964	1969	1971
INPUT PARAMETERS OF SOLAR MODELS			
$M_{\odot}$ (10^{33}g)	1.989	1.989	1.989
Age (years)	4.5×10^{9}	4.7×10^{9}	4.7×10^{9}
$L_{\odot}$ (10^{33} erg sec^{-1})	3.90	3.83	3.81
Z/X	0.028	0.019	0.019
S_{11} (10^{-22}keV-b)	3.36	3.78	3.78
S_{33} (keV-b)	1100	5000	5100
S_{34} (keV-b)	0.47	0.47	0.61
S_{17} (keV-b)	0.03	0.035	0.030
SOME OUTPUT RESULTS OF MODEL CALCULATIONS			
X	0.708	0.767	0.726
Y	0.272	0.218	0.260
Z	0.020	0.015	0.014
$\varphi(^{8}B)$ (cm^{-2}sec^{-1})	1.9×10^{7}	3.5×10^{6}	5.4×10^{6}
$\Sigma(\varphi\sigma)$ (SNU)	29	6.0	9.0
T_c (10^{6} °K)	15.7	14.9	15.3

[a]The 1964 results are from Sears (1964), the later results are from Bahcall, Bahcall & Ulrich (1969) and Bahcall & Ulrich (1971). The cross-section factors S_{ij}, where i and j are the mass numbers of the reacting particles, are directly related to the cross sections (see Fowler et al 1967, p. 539, or Clayton 1968, p. 297). The values of X, Y, and Z given are the inferred primordial values.

TABLE 4. Calculated neutrino fluxes[a]

Neutrino Source	Flux
^{8}B	5.4 E-4
p+p	6.0
$p+e^{-}+p$	1.5 E-2
$e^{-}+{}^{7}Be$	4.5 E-1
^{13}N	3.3 E-2
^{15}O	2.7 E-2

[a]All fluxes are given in units of 10^{10} per cm^2 per sec at the earth's surface ($10^{-x} \equiv$ E-X) and are taken from the standard model of Bahcall & Ulrich (1971). The calculated ^{8}B flux should be decreased by a factor of ten or more (see Section IV) when computing expected counting rates for future experiments.

TABLE 5. Levels of meaning of the ^{37}Cl experiment

Counting Rate (SNU)	Principle or Idea Tested
35	CNO cycle vs. pp-chain
9-4	best current predictions
1-1.5	basic ideas of stellar evolution
0.3	nuclear fusion among light elements as the energy source of main sequence stars

LEGENDS FOR FIGURES

FIGURE 1. Solar neutrino energy spectrum, from the model of Bahcall & Ulrich (1971). Solid lines: pp-chain neutrinos. Broken lines: CN cycle neutrinos. Fluxes are in units of no. cm^{-2} sec^{-1} MeV^{-1} for continuum sources and no. cm^{-2} sec^{-1} for line sources.

FIGURE 2. The low-energy cross section factor for the critical reaction $^7Be(p, \gamma)^8B$, largely from the data of Kavanagh, Tombrello, Mosher & Gossman (1969) (see also Kavanagh 1972), is exhibited. The smooth curve is a normalized replica of a theoretical curve given by Tombrello (1965). The boxes are the data of Parker (1966a, b, 1968). The data of Vaughn, Chalmers, Kohler & Chase (1970), in the region ≥ 1 MeV (not shown here), are systematically lower by about 25 %.

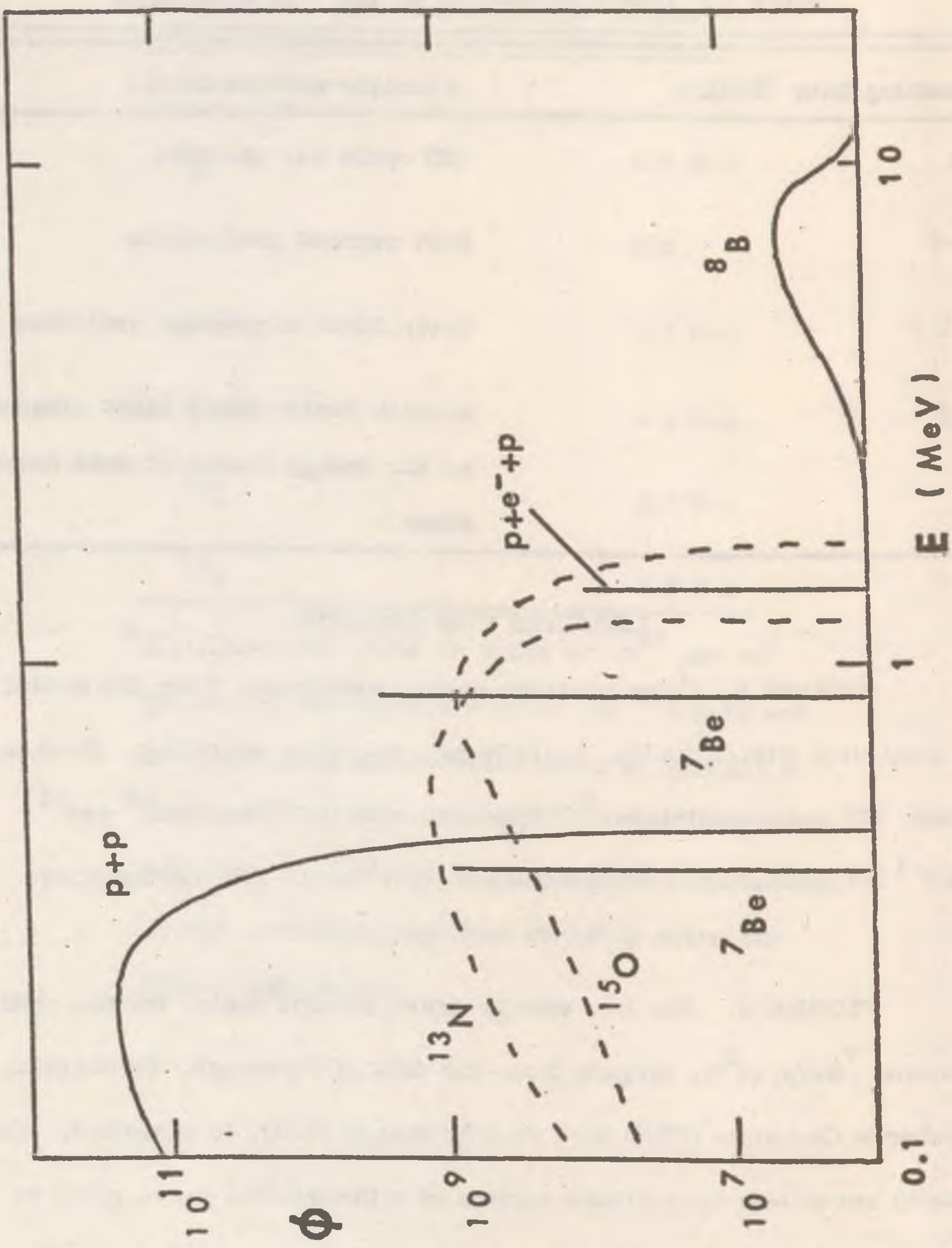

FIGURE 1

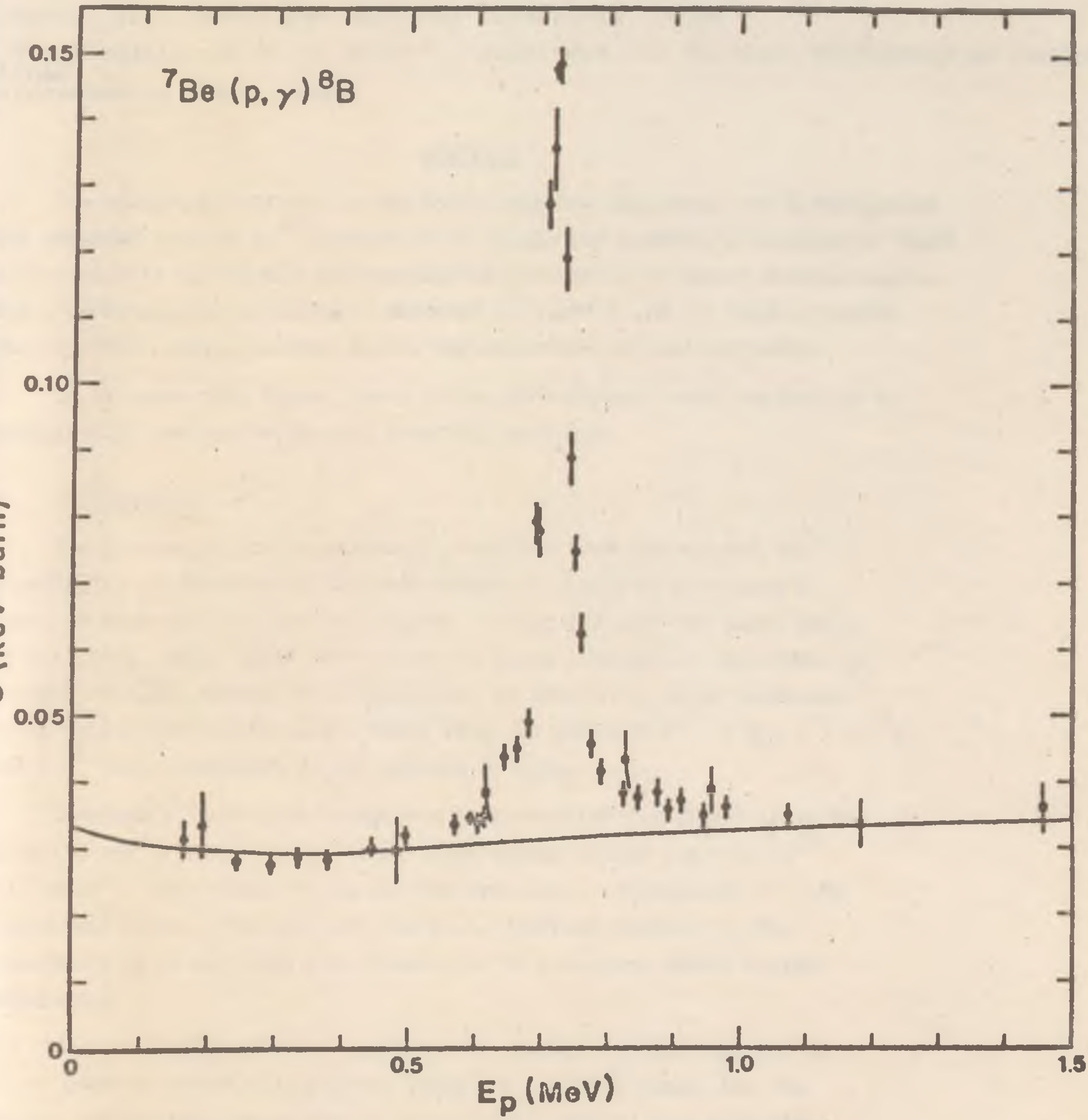

FIGURE 2

THE BACKGROUND EFFECTS IN THE SOLAR NEUTRINO EXPERIMENT

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/Presented by E.C.M.Young/

ABSTRACT

The background effects in the solar neutrino experiment are investigated. The expected rate of Ar^{37} production by cosmic ray muons as a function of depth underground is calculated and compared with measured values at shallow depths. The muon background at the main detector at 4,400 m.w.e. is then estimated. The expected background from cosmic ray neutrinos is also estimated.

It is shown that upper limits to the photonuclear cross section up to energies ~150 GeV can be derived from this estimate.

1. Introduction

The problem of the detection of neutrinos from the sun and the possibility of determining the temperature of the solar core from a measured neutrino flux are well known. During the last few years Davis et al. (1964, 1968, 1971) have mounted a large underground experiment at a depth of 4,400 metres water equivalent to search for solar neutrinos by detecting radioactive argon atoms from the reaction $Cl^{37}(\nu_e, e^-)Ar^{37}$, the Cl^{37} being contained in 10^5 gallons of C_2Cl_4.

Inevitably there is a background against which the signal is to be detected and in this case it is provided mainly by the reaction $Cl^{37}(p, n)Ar^{37}$, the protons being derived from local interactions of fast cosmic ray muons. The electron and muon neutrinos produced in the atmosphere by cosmic rays also contribute to a smaller extent to the background.

When the experiment was conceived the background was expected to be at least an order of magnitude below the neutrino signal but the signal has in fact turned out to be very much smaller than predicted (Davis, 1971, 1972). The reason for this low value is still the subject of much debate amongst astrophysicists and a number of explanations have been put forward (e.g. Bahcall et al., 1971.)

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*On leave from the University of Hong Kong

What concerns us in the present paper is the relevance of the measurements of Ar^{37} production to our understanding of underground cosmic ray phenomena. If the whole of the signal is attributed to muon and cosmic ray neutrino backgrounds, that is it is assumed that solar neutrinos give a negligible contribution, then an upper limit can be put on the magnitude of the muon cross section for interactions in which secondary protons and neutrinos are produced. Such a limit has interest for two reasons:

(i) Two experiments performed recently have suggested the existence of anomalies in the underground muon characteristics, namely, the Utah project (Bergeson et al., 1968, 1971) in which the angular distribution of muons above 1000 GeV is found to be not as expected, and the Turin work (Baschiera et al., 1970, 1971) in which an apparent excess of stopping muons is observed.

(ii) In the absence of anomalies, the source of the muon background is muon interactions by way of the virtual photon flux accompanying the muon. The yield of secondaries is related to the magnitude of the photonuclear cross section and there is thus the possibility of making an indirect determination of the magnitude of this quantity, averaged over quite a wide range of energy, at energies above those available in conventional accelerator studies.

In previous work (Ryajskaya and Zatsepin, 1965; Wolfendale, 1969) attention was directed towards predicting the expected muon background at 4,400 m.w.e.; in the present work predictions are made as a function of depth and comparison is made not only with the data from the main detector at 4,400 m.w.e. but also with measurements with subsidiary detectors at shallower depths. The cosmic ray neutrino background, which does not depend on depth, is also estimated.

2. Muon Background

As has been remarked already, the muon background arises from protons associated with muons. These protons come from a number of processes, viz: protons produced in μ^- capture, and fast muon interactions, and protons from neutron interactions: Cl^{35} (n,p) S^{35}, the neutrons coming from the same processes. The contribution from μ^- capture falls quite rapidly with depth, following the diminishing stopping/penetrating muon ratio (for example in a target of 100 g cm^{-2}

Grupen et al., 1972, estimate ratios of 6.8×10^{-2}, 6.5×10^{-3}, and 1.3×10^{-3} at 25, 300 and 4,400 m.w.e. respectively). The measurements of Short (1965) and the characteristics of the μ^- capture process and fast μ-star production have been used to estimate the contribution from the former as a function of depth. At 25 m.w.e. it is estimated that 20% of the measured background comes from μ^- capture and at the greater depths the μ^- capture contribution is found to be negligible.

In what follows we consider the problem of predicting the expected variation with depth of the background due to fast muon interactions. There can be no question of making absolute predictions since the details of the energy distribution, and cross sections for interaction, of the protons and neutrons in the C_2Cl_4 are not known with sufficient precision. It is illuminating to consider that the efficiency of producing Ar^{37} atoms with neutrons from Ra-Be source is only $\sim 7 \times 10^{-7}$ atoms per neutron (Davis et al. 1968) whereas, working back from the predicted count at the greatest depth and knowing the number of penetrating muons and details of their interaction cross section (Short, 1965) it appears that the efficiency here is $\sim 7 \times 10^{-4}$ atoms per secondary proton. There is thus a considerable difference in efficiency of protons and neutrons.

Protons will also be produced in (γ,p) reactions, the γ's arising from electromagnetic cascades coming from muon-electron collisions, direct pair production and bremsstrahlung.

Protons from all sources will be generated more efficiently the higher the muon energy and the variation of this efficiency with energy and thus with depth is required. Estimates of this variation are based on the following results:

(i) the predicted variation of neutron yield with muon energy (Ryajskaya and Zatsepin, 1965).

(ii) the measured neutron yeild from muon interactions in thick targets (Meyer et al. 1964; Allkofer and Andresen, 1968).

(iii) the predicted variation of slow pion secondaries with muon energy (e.g. Grupen et al. 1972).

All the variations can be represented as E_μ^α, where α is a slowly varying function of E_μ, and the different processes all give $\alpha \simeq 0.7$ for the range of muon energies in question: $\sim$5-300 GeV. In the present work we adopt the predictions referred to in (iii) which give α = 0.88, 0.77 and 0.67 for E_μ = 10, 50 and 300 GeV, respectively. Insofar as pions, like protons, are produced in the secondary cascades it seems reasonable to take this component although the results using the neutron component would not be significantly different. The relative Ar^{37} rate then follows as $R(d) = \int_0^{\pi/2} I(\theta,d)\, \bar{E}_\mu^{\alpha}(\theta,d) \cdot 2\pi \sin\theta \, d\theta$ where θ is the zenith angle. This integral has been evaluated using the well known relation $I(\theta) = I_v \cos^n \theta$, where n is an increasing function of depth, and the variation of $\bar{E}_\mu$ with depth given by Craig et al. (1968).

For the variation of vertical intensity, I_v, with depth the values given by Osborne et al. to 1000 m.w.e. and the recent survey of Groom (1971) at greater depths have been used. For example, the intensities are 1.90×10^{-3}, 2.90×10^{-5} and $4.44 \times 10^{-9} cm^{-2} s^{-1} sr^{-1}$ at 25, 300 and 4,400 m.w.e. respectively. After correcting for the 20% μ-capture at 25 m.w.e. R(d) has been normalised at that depth with the result shown in Figure 1. This curve then represents our estimate of the muon background as a function of depth.

3. Comparison of Prediction with Experiment

The theoretical predictions are compared with the experimental values which have been obtained with a smaller tank containing 1850 gallons of C_2Cl_4 exposed at several depths. At these shallower depths all the Ar^{37} production rate is due to muon interactions. Agreement at these depths will give greater confidence of the prediction at the main detector at 4,400 m.w.e.

As the ground surface above the detector tank locations is not flat, the vertical depths must be corrected for the change in muon intensities for the same zenith but different azimuthal angles. These corrections have been made and the experimental values of the Ar^{37} production rate at depths 275, 408, 620 and 1080 m.w.e. respectively are shown in Figure 1 where the predictions are compared. It can be seen that agreement is good. At the depth of 4,400 m.w.e. the predicted background from muon interactions is estimated to be 0.065 per day per 10^5 gallons of C_2Cl_4.

4. Cosmic Ray Neutrino Background

Cosmic ray neutrinos contribute to the background through the following processes: $Cl^{37}(\nu_\mu, \mu^-)Ar^{37}$ and $Cl^{37}(\nu_e, e^-)Ar^{37}$. Here the main contribution comes from neutrinos in the energy range 200-800 MeV. In this energy range, the neutrino intensity is affected by the local geomagnetic cut-off rigidity. In the present calculation the low energy neutrino spectra given by Tam and Young (1971) have been used with appropriate corrections for the geomagnetic effect. The cross sections for the above processes are different from the well-known cross sections (see, for example, Young, 1967; Budagov et al., 1969) for free nucleons, for two reasons; firstly, reactions can occur only when the recoil proton achieves a momentum allowed by the Exclusion Principle and, secondly, the recoil proton may be so energetic as to escape from the nucleus - the resulting nucleus will not contribute to the background. Allowance has been made for these effects using the results of Løvseth (1963) and Bell and Løvseth (1964). The result is a considerable reduction in the cross section when the nucleons are bound in the nucleus.

The background from the cosmic ray neutrino interactions has been estimated and the total contribution to the Ar^{37} production is found to be 0.024 per day per 10^5 gallons of C_2Cl_4.

5. Discussion

The prediction for the muon background is based on the assumption of a constant photonuclear cross section at photon energies above a few GeV. Any experimental excess above the predicted background can thus be attributed to a photonuclear cross section, σ_γ, higher than that pertaining to muons at 25 m.w.e., solar neutrinos or some anomalous behaviour of the muon (or a mixture of all three).

Adopting the first possibility, values of σ_γ can be derived, and are shown in Figure 2. The value of E_γ is taken to be a little less than $\bar{E}_\mu/2$, this being the median energy of the virtual photons responsible for the interactions which give the background. Also shown in Figure 2 are the recent values of σ_γ from the underground experiment of Bezrukov et al.

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Apart from the point at the greatest depth, the assumption that fast muons provide virtually the whole of the background is very reasonable and the agreement between the values of σ_γ derived by the two methods gives confidence in the method. There is thus not much room for uncertainty in the predicted background at 4,400 m.w.e: 0.09 Ar^{37} atoms in 10^5 gallons of C_2Cl_4 per day. The measured production rate is 0.18 $\pm$ 0.10 per day (Davis, 1972), i.e. less than one standard deviation above the background. Clearly, at the present time there is still no firm evidence that solar neutrinos have been detected.

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Fig 1 Ar^{37} production rate vs depth

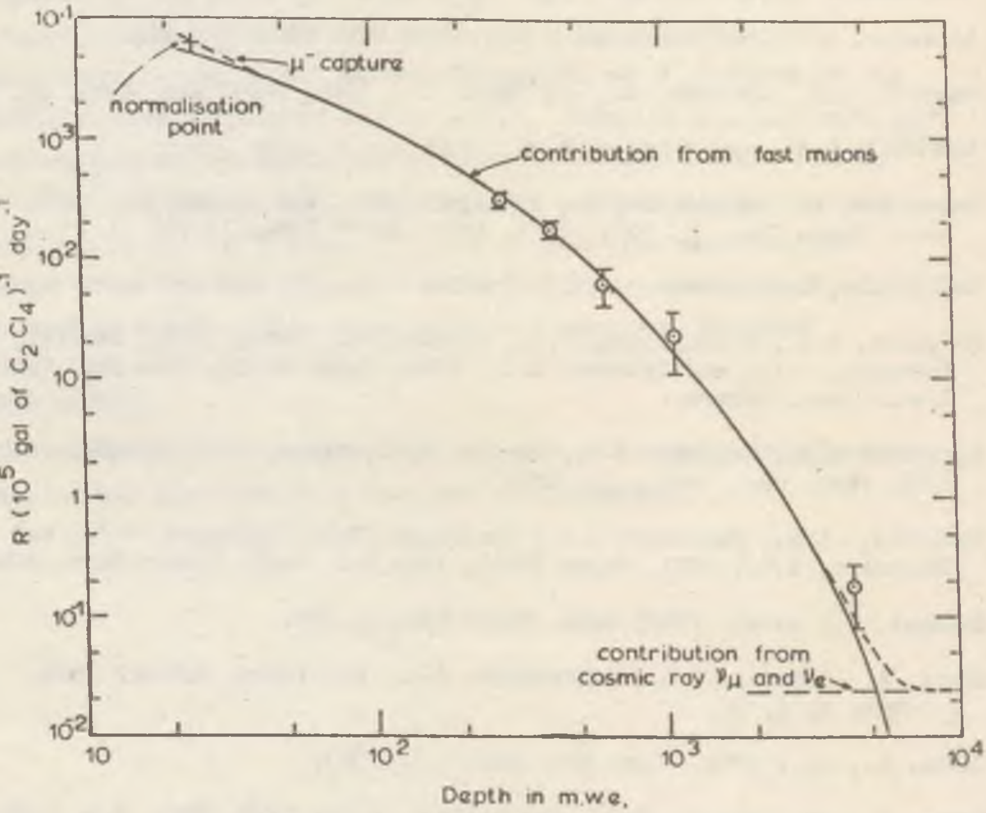
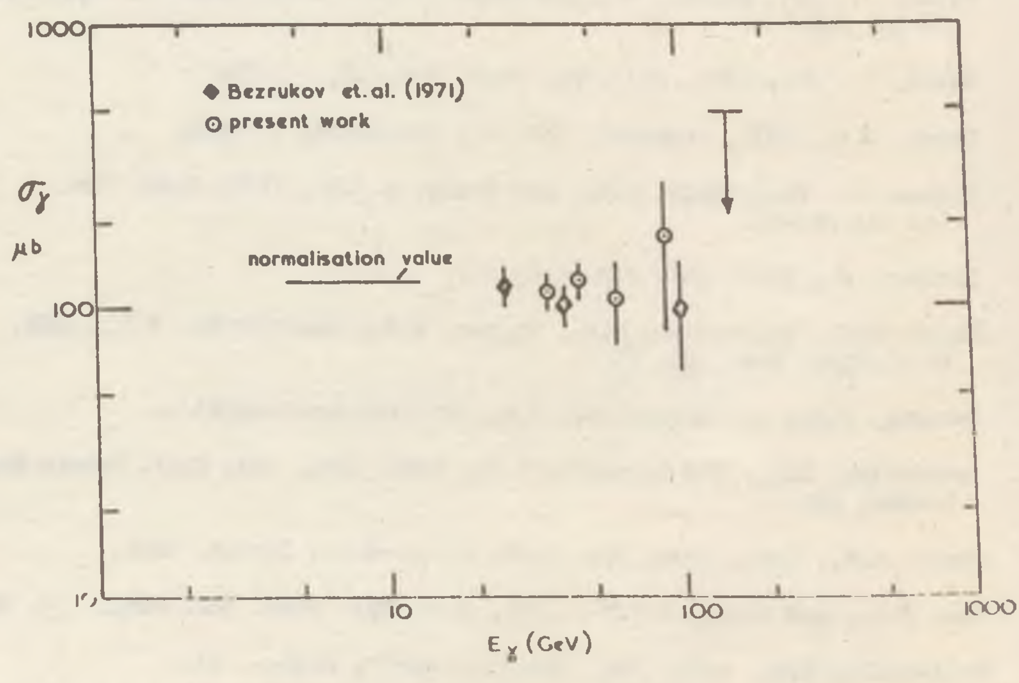


Fig. 2. Photonuclear cross section vs photon energy



THE INVESTIGATION OF THE BACKGROUND IN THE SOLAR NEUTRINO PROPORTIONAL COUNTERS (DISCUSSION REMARK)

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Solar neutrino experiments require counters with an extreme low intrinsic background /$\ll$I count per day/. The counter background, like that of any detector, depends on I/the radioactive contamination of the counter material; 2/the γ-radiation of natural radioactive isotopes being present in the laboratory environment; 3/cosmic rays.

All the calculations and experimental data presented in this paper are obtained for the counters with 30 mm length, 5 mm diameter and 0.3 mm wall thickness, because just the counter of this type has been used to count Ar^{37} in Davis' experiment. The obtained results could be applied to any detector.

The background of radioactive contaminations in counter materials depends mainly on β-activity because γ-activity is of the same order but the efficiency of γ-quantum registration in the counter is about I%. The main β-activity depends on the isotopes of the uranium and thorium radioactive families and potassium 40. Total count rate in all region of pulses from β-activity

$$N_\beta = I.5 x IO^9 \alpha_{U,Th} \text{counts/day}; \quad N_\beta = I.3 x IO^5 \alpha_K \text{counts/day}$$

$$N/Ar^{37}/ = \frac{I}{I5} N_{total}.$$

We shall have I count per month in Ar^{37} region if $\alpha_{U,Th} < IO^{-9}$ g/U,Th/ and $\alpha_K < IO^{-5}$ gK in one gramm of the counter material.

The background of γ-rays from ambient materials.
The materials of usual laboratory rooms /concrete, brick etc./ contain uranium and thorium in a quantity of IO^{-6}g/U+Th/ and potassium in a quantity of IO^{-2}gK in one gramm of ambient materials. Total count rate is

$$N_\gamma = 0.85 x IO^9 \alpha_{U,Th} \text{counts/day}; \quad N_\gamma = 4.3 x IO^4 \alpha_K \text{counts/day}.$$

We shall have that I count per month in Ar^{37} region requires the shield attenuating outer γ-background more than IO^3.

Cosmic ray background.

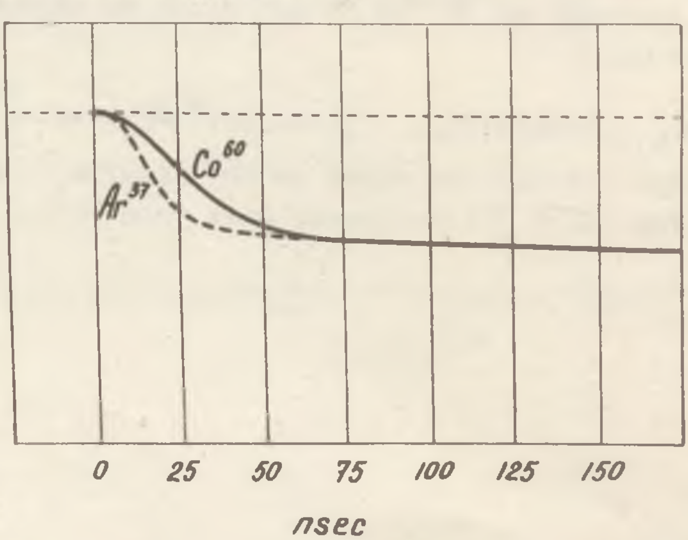
This kind of background is due to direct muon penetration through the counter as γ or corpuscular emission produced by interactions of muons. The cosmic ray background of the first type may be excluded by an anticoincidence system, therefore we shall take into consideration only the background of the second type. γ-emission results from: electromagnetic interactions of muons; capture of μ^-; inelasting scattering and radiation capture of neutrons produced from μ^{--}-capture and photonuclear interactions of muons; bremsstrahlung and following annihilation of positrons produced in μ^+-decay; production of γ-active isotopes taking place at μ^- capture and thermal neutron capture.

γ-ray background being in equilibrium with muons has been determined experimentally and was equal I/40 counts per hour. The calculation of the total background resulting from muon interactions gives $N_\mu = 2.4$ counts per day. To decrease this type of background the installation should be placed underground.

The experimental study of counter background.

Besides usual active and passive shields the background in Ar^{37} region may be decreased by using the rise-time discrimination. We have carried out the calculation of the pulse shape for the case of point and track ionizations. The calculation has shown that rise-time of an initial part of the pulses was equal to $\sim$ 20 nsec for point and $\sim$ 50 nsec for track ionizations at wire diameter 8μ. Experimental study of the pulse shape has confirmed the calculations.

It is reasonable to place the counter at great depth underground because in this case the necessity for an anticoincidence disappears due to small flux of muons. It must essentially simplify the construction of the shield from surrounding radioactivity to obtain $\alpha_{U,Th} \ll 10^{-9}$ and $\alpha_K \ll 10^{-5}$.



NEW APPROACHES TO THE SOLAR NEUTRINO PUZZLE*

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I. Introduction

The attempt to verify directly the nuclear fusion power source of the sun, the search for solar neutrinos, has reached a dramatic and unanticipated stage.

Recently Davis[1] has reported an upper limit for the interaction rate of solar ν_e with $C\ell^{37}$ of 1×10^{-36} interactions/$C\ell^{37}$ sec., compared with the predictions of the standard solar models[2] of 9×10^{-36} interactions/$C\ell^{37}$ sec. These results cannot be understood in terms of the usual solar models or in any reasonable variations of them[3].

Two classes of suggestions for explaining the puzzle of the missing solar neutrinos have been proposed. One of these suggests that the solar neutrino production rate is time dependent[3] and that the problem therefore lies in the nature of the solar model. This time dependence could either involve the nuclear energy generation rate and therefore may manifest itself in a time variation of the solar constant, or it could involve a time dependence of central solar temperature and thus mainly effect the neutrino emission rate from B^8 decay and to a lesser extent the formation of Be^7 and so its subsequent electron capture.

We do not propose to engage in the construction of solar models but rather will merely explore the terrestrial effects that might result from such solar behavior. A variation of the solar constant, of course, will cause a change in the earth's mean temperature and hence we are immediately led to an examination of the earth's temperature history.

II. Temperature History of the Earth

There appear two clear patterns; one involves a long term periodicity consisting of warm periods of about 200-250 x 10^6 years long interspersed with short cold periods of the order of 10 x 10^6 years[4,5,6)] (See Fig. 1). The other pattern consists of a fine structure within these cold periods with a 9 x 10^4 year period, the valleys of which correspond to the ice ages. At present we are in a cold period near a fine structure peak.

At least two previous cold periods have been established and the existence of a third and fourth one is suspected.

The evidence for the fine structure is well established only for the present cold period[(7,8)]. There is, as yet, no clear indication of its existence[(9)] during previous cold periods, nor, indeed is there any reason to believe that it is, or is not, present during the warm periods.

There are several geological models to account for the fine structure on the basis of various terrestrial effects[(10)]. We wish, however, to focus on the long term periodicity. Not only have no terrestrially-based theories been advanced to account for this phenomenon, but there is no terrestrial time constant that approaches 10^8 years[11]. It is difficult[12], if not impossible, to demonstrate any causal relationship between the earth's temperature variation and the nuclear reaction rate in the solar core. In order to establish such a link, it would be necessary to first establish the earth's neutrino bombardment history and then to correlate that history with the time dependence of the earth's temperature variation corrected for the $\sim 10^7$ years time lag between the thermal and neutrino transit times of the sun[3] - a formidable task.

III. Neutrino Archeology

The integrated neutrino flux incident on the earth during its history is recorded in the nuclear transitions induced by these neutrinos. Of particular interest is the simulated double beta decay[13] which involves

$$\nu_e + A^Z \rightarrow A^{Z+1} + e^-$$

followed by

$$A^{Z+1} \rightarrow A^{Z+2} + e^- + \bar{\nu}_e$$

for cases where the mass of A^{Z+1} > mass of A^Z. A measure of the fractional amount of A^{Z+2} in old ores of A^Z provides an integration of the incident ν_e flux over the age of the ore. The analysis techniques involved are similar to those employed by Kirsten et al.[14] and Takeoka and Ogata[15] to measure the Te^{128}, Te^{130} and Se^{82} double beta decay lifetimes. By examining various elements A^Z with different energy thresholds and ores of differing ages it is possible, in principle, to establish both the neutrino spectrum and the time dependence of the flux. Unfortunately, the sensitivities required, about 10^{-18} gms of A^{Z+2} per gm of A^{Z}, are several orders of magnitude greater than those involved in the tellurium and selenium experiments. Recently, a joint General Electric-Los Alamos group[(16)] has carried out a successful search for traces of Pu^{244} in an old precambrian Cesium ore. They were able to detect a concentration of 10^{-19} gms Pu^{244} per gm of ore, an order of magnitude greater sensitivity than required for the solar neutrino induced double beta decay.

One more point should be made with respect to the earth's temperature history. The average warm period temperature has been about 10°C. higher than the present average earth temperature[(17,18)]. If we assume that the earth is in radiative equilibrium with the sun and that the earth's albedo is the same now as it was during the warm periods, then the 3% increase in earth's temperature implies a 12% increase in the solar constant and so a 12% increase in the sun's typical central temperature. Because of the high power dependence of neutrino production, especially from B^{8}, on central solar temperature, during the warm periods there may have been effective solar neutrino induced transition rates an order of magnitude greater than those predicted by the present solar models. This thought at least harbors a little hope for attempting neutrino archeology.

IV. Neutrino Peculiarities

The second class of suggestions to elucidate the solar neutrino puzzle proposes changes in neutrino character in the interval between production in the sun and arrival at the earth so that the solar neutrino no longer behaves like a ν_e at the earth. One of these proposals, by Bahcall et al.,[19] points out that we cannot exclude the possibility that the ν_e decays with a mean free path of $\sim 10^7$ km, a distance large compared with terrestrial dimensions but small on the stellar scale. This decay process requires a new, extremely weak interaction that violates lepton conservation.

Another approach that also involves lepton nonconservation is a suggestion by Pontecorvo[20,21] that there occur neutrino oscillations between ν_e and ν_μ analogous to K°, $\bar{K}^\circ$ oscillations. These oscillations have the effect of reducing a ν_e beam to an average of 1/2 of its original intensity at distances large compared to the oscillation period. Recently, Pontecorvo[22] has extended this idea to a system of N leptons so that the ν_e beam is eventually reduced to 1/N of its original intensity.

The decay and oscillation hypotheses differ in their experimental detectability. If a generalized neutrino-electron scattering interaction exists with comparable strength for all neutrinos[24], then a detector based on ν-e scattering will not be effected by neutrino oscillation. The situation for such a detector in the case of neutrino decay may be quite different. The radio-chemical detector, however, is rendered insensitive by both proposals.

V. Neutrino Astronomy

The long characteristic lengths inherent in these suggestions eliminate the likelihood that the phenomena in question can be investigated with terrestrial neutrino sources. We must thus rely on stellar sources, either steady ones such as the sun, or pulsed ones such as collapsing stars. An absence of detected neutrinos from either of these classes of sources would leave uncertain whether this lack of neutrino signal should be attributed to the characteristics of the reactions in the source or to the behavior of the neutrinos between source and detector. The detection of neutrinos, on the other hand, would immediately eliminate the decay hypothesis and would limit the oscillation possibilities.

Collapsing stars represent the primary source of stellar neutrino pulses. The frequency of such events in our galaxy has a lower bound of 1/30 per year set by the frequency of supernova occurrence[24,25], a likely value of ~1/3 per year based on the stellar population and mass distributions[26], and an upper bound of ~1/day based on the possibility that the sources of Weber pulses[27] are also intense sources of neutrino pulses.

The feasibility of such searches has been discussed by

Zatsepin[28] and collaborators. They stress the necessity of using large volume detectors and of operating in a background free environment. Furthermore, the use of a technique, similar to that employed by Weber[29], involving correlations between widely separated detection stations will insure that detected pulses are of an extra-terrestrial, and not local, origin.

Since the required techniques are now in hand and adequate background reduction is feasible, only the uncertainties in the predictions of neutrino pulse intensity and their effect on detector volume stand in the way of this program of neutrino astronomy. It is exciting to look forward to such observations and to their contributions to the understanding of the solar neutrino puzzle.

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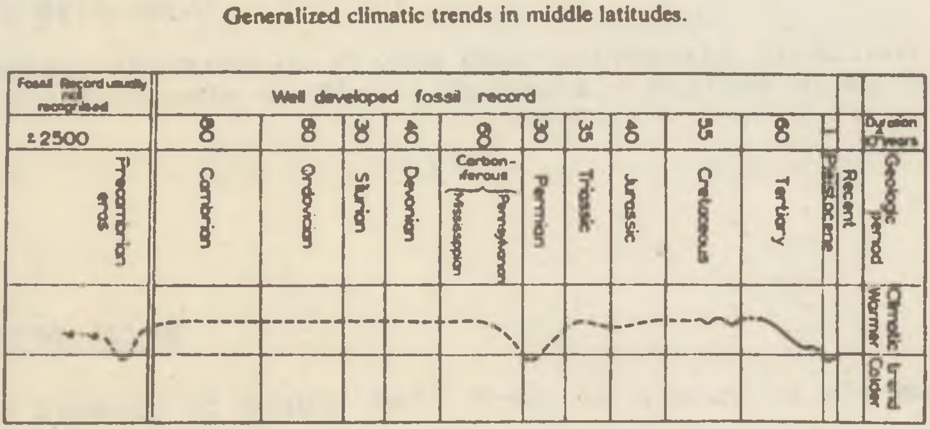

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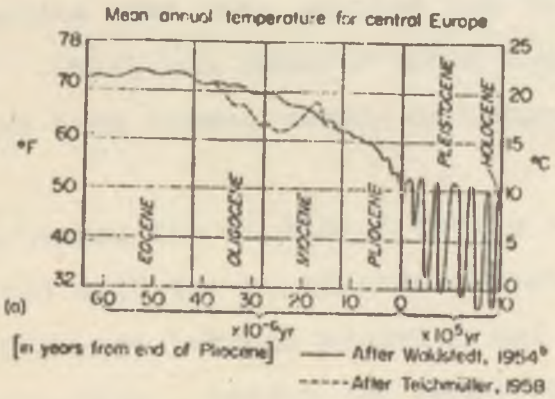

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Fig. 1

Temperature History of the Earth. /Fig. 1a from Ref. 4 and Fig. 1b from Ref 18/.

DOUBLE BETA DECAY

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1. Introduction

The history of double beta decay is almost as old as that of the neutrino . It was in fact only one year after the Fermi theory of beta decay that Maria Goeppert Mayer [1] suggested the existence of a process where " a metastable isobar can change in a more stable one by simultaneous emission of two electrons". In these 37 years a number of usually very sophisticated experiments have been carried out to search for this elusive process , which has however never been "seen" directly. Its existence has been proved rather unambigously by "indirect" methods.

Since the beginning the problem of double beta decay has been strictly connected with the highlights of weak interaction physics, like difference between neutrino and antineutrino, parity non conservation , lepton non conservation , neutrino mass and polarization and even , as suggested recently , CP non conservation.

Let us consider a nuclear triplet (A,Z) , (A,Z+1) and (A,Z+2), where the single beta decay of (A,Z) into (A,Z+1) is either energetically forbidden , or at lest strongly inhibited by large change in the spin-parity ground states of the two nuclei. Double beta decay is a transition where the charge of (A,Z) changes by two units with the emission of two electrons. This reaction could in principle occur in two channels :

$$(1) \qquad (A,Z) \rightarrow (A,Z+2) + 2\,e^-$$

$$(2) \qquad (A,Z) \rightarrow (A,Z+2) + 2\,e^- + 2\,\bar{\nu}_e$$

where the presence of channel (1) would imply , of itself, lepton non conservation . This neutrinoless $\beta\beta$ decay could occur according to the scheme

$$(3) \qquad n \rightarrow p + e^- + \bar{\nu}_e$$

$$(4) \qquad \nu_e + n \rightarrow p + e^-$$

where $\bar{\nu}_e \equiv \nu_e$ is a virtual neutrino which , interacting with the intermediate nucleus , induces its decay with the emission of a second electron.

Clearly other decays with double change in nuclear charge could occur like :

(5) $(A,Z) \rightarrow (A,Z-2) + 2\,e^+ \quad (+\,2\,\nu_e)$ (double positron emission)

(6) $2\,e^- + (A,Z) \rightarrow (A,Z-2) \quad (+\,2\,\nu_e)$ (double electron capture)

(7) $e^- + (A,Z) \rightarrow (A,Z-2) + e^+ \quad (+\,2\,\nu_e)$ (electron capture and positron emission)

The comparison with experiments is however much more difficult for these decays, where the predicted lifetimes are in general much larger (2). We will limit ourselves here to double beta decay.

The direct detection of double beta decay would already be a great success experimentally , since the present experimental limits on half lives are of $10^{20} - 10^{21}$ years. It would be however even more important to prove the existence of reaction (1) , whose rate should dominate on that of reaction (2) if the two matrix elements are comparable. In fact the phase space available to decay (1) would be much larger , since this decay is mediated by a virtual neutrino whose energy can be as high as 30-40 MeV , namely many times the energy available to the two neutrinos in two-neutrino $\beta\beta$ decay.

In reaction (1) the two electrons share the total transition energy , since the recoiling energy of the nucleus is negligible. A peak should therefore appear in the spectrum of the sum of the two electron energies.

2. Evaluation of the lifetimes for two neutrino and neutrinoless double beta decay

I will not enter into theoreticall details on double beta decay since lepton non conservation will be considered theoretically in the talk by prof.G.Marx. I will only remind here that two neutrino double beta decay is always possible ([2,3]), even if leptons are not conserved. This process has been evaluated as a second order effect of the $\Delta S=0$, $\Delta Q = \pm 1$ interaction which is responsable , in the first order , of single beta decay.

The half lifetimes for two-neutrino double beta decay are reported in Table 1 for all possible "nuclear triplets". In this calculation all nuclear matrix elements have been assumed to be the same and it is in fact the uncertainty on these matrix elements which is responsible for the large ($10^{\pm 2}$) error.

Neutrinoless double beta decay has been evaluated recently both as a "first order" ([4]) and as a "second order" ([5]) effect.

In both approaches the decay could involve the role of Δ, resonances , whose precence in the nucleaus has been recently indicated [6]. Reactions like :

$$\Delta^{-} \rightarrow p + e^{-} + e^{-} \quad \text{and} \tag{8}$$

$$n \rightarrow \Delta^{++} + e^{-} + e^{-} \tag{9}$$

should enhance the probability for neutrinoless double beta decay despite the low probability for the presence of these resonances inside the nucleus. Detailed calculations on neutrino_ less double beta decay on the basis of the "second order interactio approach have been carried out by H.Primakoff and S.P.Rosen[5]. I have computed from their formulae the lifetimes for neutrinoless $\beta\beta$ decay with and without resonance reported in Table I . Also reported in Table I are the values based on the assumption that the lifetime experimentally found for ^{130}Te (see later) be due only to neutrinoless $\beta\beta$decay , and that all nuclear matrix elements be the same. Both assumptions have very poor justifications. In fact V.A.Khodel [7] and A.N.Huffmann [8] have shown that matrix elements can differ among themselves by up to an order of magnitude.

As pointed out by H.Primakoff and D.S.Sharp [9] lepton non conservation could be correlated to CP non conservation . The lepton current in a single β decay could be written in the form

$$L_\lambda = \psi_e^{+} \gamma_4 \gamma_\lambda \left[(1+\gamma_5) + \eta(1-\gamma_5)\right] \psi_\nu \tag{10}$$

where ψ_e and ψ_ν describes the electron and the neutrino (a la Majorana) and the imaginary part of the complex parameter η gives CP non conservation , while its modulus gives lepton non conservation.

3. Experimental results

The many experiments carried out so far on $\beta\beta$ decay have been based on "indirect" methods , consisting in chemical and mass spectroscopic analysis of samples of known geological age and on "direct methods.

The most recent experiments with geological methods were on $\beta\beta$ decays of ^{82}Se , ^{128}Te and ^{130}Te; to ^{82}Kr , ^{128}Xe and ^{130}Xe , respectively. Kripton and Xenon have been extracted from rocks of known geological age and analysed with a mass spectroscope to detect abnormal aboundance of the daughter isotopes with respect to the atmospheric composition of Kripton and Xenon. The age of the rock , which in the earlier experiments was only determined by geological methods is presently obtained also from the aboundance in the rock of ^{4}He and ^{40}Ar from Uranium and Potassium decays.

In the experiment by E.K.Gerling et al (10) large excesses of Xenon isotopes 128,129,130,131,134 and 136 have been found, but the presence of isotopes 129,131,134 and 136 has been attributed to reactions other than double beta decay, like spontaneous and neutron induced fission of uranium , which is contained rather aboundantly in the rock . The presence of ^{128}Xe could be due to the reaction $^{127}I\ (n,\gamma)\ I^{128}$ followed by a single β decay of ^{128}I. The authors give therefore only a lower limit of 7.7×10^{20} years for $\beta\beta$ decay of ^{128}Te. The presence of ^{130}Xe has on the contrary been attributed to double beta decay, **but** the half lifetime is difficult to evaluate due to the possibility that some Xenon may have excaped from the rock as a result of hydrothermal alterations during its past history. A detailed calculation of the probability for Xe to have diffused from the rock gives an half lifetime of $(3 \pm 0.4) \times 10^{20}$ years for ^{130}Te. This value is lower than the limit of 3.3×10^{21} years previously obtained by R.J.Haiden and M.G.K.Inghram ([11]), and also lower than the value found by Kirsten et al ([12]). It should be taken with some care due to a possible overestimation of the Xenon loss.

Similar experiments have been carried out by N.Takaoka and K.Ogata ([13]) and by N.Takaoka ([14]) on various samples containing Tellurium. Excesses of Xe isotopes 128,129,130 and 131 have been found . The ^{130}Xe finds no explanation apart double beta decay of ^{130}Te with an half life of $(8.0 \begin{smallmatrix} +0.4 \\ -0.9 \end{smallmatrix}) \times 10^{20}$ years. The possibility of xenon loss due to hydrothermal alterations should be cancelled by the fact that the K-Ar method has been used to evaluate the age of the rock. In fact Argon should have excaped more than Xenon , but a comparison with the age of the rock determined geologically indicates that this effect should be negligible. The positive result on ^{128}Xe reqiires a more careful study of the background : for instance the mass spectrometer has some "memory effect" for Xenon and some ^{128}Xe was added to the atmospheric Xenon to test the apparatus. Moreover the excessive ^{128}Xe could be due to neutron capture by ^{127}I. There is therefore in my opinion no real evidence for double beta decay of ^{128}Te , even if a value of 6×10^{22} years was suggested . This has been taken as evidence for neutrinoless double beta decay.

This last result is not confirmed by the latest experiment by Kirsten et al ([12]) . These authors could use an excellent sample of almost pure - (99.4 $\pm$ 0.6) % - native Tellurium with a negligible contamination by radioactive elements . A very careful mass spectroscopic analysis indicates no abnormal distribution for Xe isotopes other than ^{130}Xe, whose aboundance is 50 times the natural aboundance in atmospheric Xenon. The "thermal history" of the rock indicates that some Helium , but no Argon, has escaped by diffusion ; a Xenon loss can therefore be excluded "a fortiori". The half life for ^{130}Te is $2.2 \times 10^{21 \pm 0.12}$ years.

A similar experiment ([15]) on the transition ^{82}Se - ^{82}Kr also gives a positive result , with an half life of $6 \times 10^{19 \pm 0.3}$ years. These are up to now the only positive evidences for $\beta\beta$ decay ; they should be confirmed by experiment based on direct methods for at least two reasons :

a) it is impossible with geological methods to distinguish clearly between neutrinoless and two neutrino double beta decay, since the theoretical values for half lives are , at least for the moment , uncertain by up to two orders of magnitude;

b) many other phenomena , which could produce the daughter isotopes and therefore an "apparent double beta decay"should be considered. Some of them , like for instance the reaction $^{8}Te(\nu ,e^{-})I^{128}$ and ^{130}Te (ν ,e^{-}) I^{130} by sun neutrinos, with subsequent decay of the I isotopes into Xe, are difficult to be completely ruled out. As pointed out by K.Lande at this conference "apparent " double negatron and positron decay can in principle be used to determine the fluxes of neutrinos and antineutrinos from the sun and from the radioactive materials on the Earth , respectively[16].

Almost all detectors , with the exception of bubble chamber, have been used to search double beta decay. Let me quote nuclear emulsions, cloud chambers[17], Geiger and proportional counters, scintillators , spark chambers and solid state detectors.

A review of the older experiments can be found in reference [2] ; let me limit myself here to three tipical recent experiment which gave limits comparable with those obtained with geological methods.

An underground experiment on the transition ^{48}Ca - ^{48}Ti has been carried out with the apparatus shown in Fig.1 by Bardin et al[18]. A thin source of 10.6 gr of 97% enriched ^{48}Ca , deposited as CaF_2, is sealed to form the central plate of a discharge chamber in a field of 400 Gauss .The chamber was triggered by means of two scintillation counters arrays, which were also used to measure the electron energies , while the pictures of the discharge chamber were mainly used to "visualize" the events and to determine the sign of the tracks. In 1150 h of running time , only four events were found with a sum of the two electron energies of more than 3 MeV and only one with a sum compatible with the transition energy (4.267 ± 0.009) MeV [19].Due to the presence of background this result can be used only to set a lower limit on the lifetime for neutrinoless $\beta\beta$ decay (2×10^{21} years with 80 % confidence level). The two electron events in this experiment give a continous sum spectrum of the two electron energies , as expected for two-neutrino double beta decay. However the comparison with the spectrum obtained with a dummy source of ^{40}Ca shows that a small contamination of Radium with an activity as low as one picocurie is present in the ^{48}Ca source. By limiting the analysis to events with an energy of more than 2.2 MeV, where the background due to this contamination is negligible, one can set a lower limit of 3.6×10^{19} years for the lifetime of two neutrino $\beta\beta$ decay.

The other two experiments are quite different between themselves in technique , but rather similar in philosophy: the use of a double beta active sample both as source and as detector of the decay. E.der Mateosian and M.Goldhaber (20) have used as scintillator a CaF_2 cristal with natural Calcium and later CaF_2 cristals enriched in ^{48}Ca and ^{40}Ca (the latter for background measurements). As shown in Fig.2 the ^{48}Ca and ^{40}Ca detectors were placed side by side within a plastic scintillator operated in anticoincidence with both detectors. The entire apparatus is placed within a section of a naval gun to provide additional shielding against local radioactivity. The spectra of pulses from the detectors are shown in Fig.2; no peak appears for the ^{48}Ca sample in the region corresponding to the transition energy . The limit on neutrinoless $\beta\beta$ decay was 2×10^{20} years. A new experiment with a carefully purified ^{48}Ca sample to reduce radioactive impurities and a better anticoincidence shielding is in progress (+) and the limit is presently increased to 10^{21} years.

A series of three experiments has been carried out by A.Pullia and myself in collaboration with the Nuclear Chemistry Division of Euratom (Ispra). In all experiments Ge (Li) cristals were used both as source and detector of the neutrinoless transition :

(11) $$^{76}Ge \longrightarrow {}^{76}Se + 2e^-$$

(†) Private communication by E.der Mateosian

The choice of these solid state detectors is based on the following reasons :

a) ^{76}Ge is contained in natural Germanium with an isotopic aboundance of 7.67 % : its transition energy to ^{76}Se is (2.045 ± 0.004) MeV;

b) the mass of Ge(Li) detectors can reach a few hundreds grams , thus providing a reasonable amount of $\beta\beta$ active isotope;

c) these detectors are remarkably free from impurities due to the way they are prepared : they do not require windows, which could contain radioactive elements etc.;

d) the energy resolution of these detectors is excellent, and this can be very useful to search neutrinoless $\beta\beta$ decay.

In these experiments detectors of different sizes were shielded from local activity with layers of mercury,copper, low activity and normal lead , and with layers of paraffin and cadmium to moderate and absorb the neutrons.A typical set-up is shown in Fig.3.Pulses from the detector were sent, after amplification , to a multichannel analyser . The detector and the electronic chain were routinely tested by extracting the detector from the shielding and by exposing it to radio_ active sources.A brief description of the three experiments follows :

a) in a first experiment at sea level a 17 cm^3 useful dimension detector was used for a total of 712 h of running time [21]. No peak was found in the pulse spectrum from the detector in the region corresponding to the transition energy; the corresponding lower limit for neutr noless $\beta\beta$ decay is 3×10^{20} years with a 68 % confidence level.

b) A second experiment was carried out with a similar set-up and a 24 cm^3 useful dimension detector in the underground laboratory "Monte dei Cappuccini" at a depth of 70m of water equivalent . No evidence fo neutrinoless $\beta\beta$ decay was found in the spectrum obtained in 2800 h of running time , with a limit of 1.2×10^{21} years at a 68 % confidence level.

c) A third experiment has been carried out with a 68 cm^3 useful dimension detector and an improved set-up in a laboratory situated in the Mont Blanc tunnel connecting Italy to France. The depth (about 4300 m w.e.) completely eliminates the background due to cosmic radiation . This fact and the use of a spectrum stabilizer connected to a new 4096 Laben analyser permits a detailed study of the spurious counting due to the presence of radioactive contaminants in the materials immediately surrounding the detector and perhaps in the detector itself. In the first series of measurements totalling about 2100 hours of running time , the support of the detector was made of 1100 Al and the vacuum cup of the cryostat of electrolitic copper . The results are shown by the upper curve of Figs.4, where many peaks appear, which clearly indicate the presence of the ^{235}U , ^{238}U and ^{232}Th series and of ^{40}K. This radioactivity is however very low , and can be seen only because the background due to cosmic rays is completely eliminated. As an example the clear peak at 2.615 MeV corresponds to only 0.25 counts per hour. No peak appears in the region of neutrinoless double beta decay , where the average counting rate is $(4.3 \pm 0.2) \times 10^{-3}$ counts $(keV)^{-1}h^{-1}$. The corresponding limit on neutrinoless $\beta\beta$ decay is 2.8×10^{21} years with 68 % confidence level.

In order to reduce the background we substituted both the aluminum of the support and the copper of the cup with oxigen-free high-conductivity copper. The results obtained in 2300 hours of running time are shown by the lower spectrum of Figs.4, where the contaminations from uranium and expecially from the thorium series appears to be considerably reduced. The average counting rate in the region of neutrinoless $\beta\beta$ decay is now (2.0 ± 0.2) counts $(\mathrm{keV})^{-1} h^{-1}$. No peak appears at 2.045 MeV ; the corrisponding limit on neutrinoless $\beta\beta$ decay is 4.6×10^{21} years with 68 % confidence level.

4. Limits on lepton non conserving amplitude based on double beta decay

Positive results on double beta decay have been obtained only for ^{130}Te ($\tau_{1/2} = (2.2 \pm 0.6)$ years) and for ^{82}Se ($\tau_{1/2} = (6^{+6}_{-3} \times 10^{19}$ years) with geological methods. They are both compatible with the theoretical predictions of Table I for two-neutrino decays if the errors in the evaluation of nuclear matrix elements are taken into account . The ratio between the two experimental lifetimes is $10^{1.56 \pm 0.33}$, while the expected ratios for two-neutrimo decay , and neutrinoless decay with and without resonance inside the nucleus are $10^{0.48}$, $10^{0.14}$ and $10^{0.29}$, respectively. In all cases these predictions are in good agreemnt with experiments if the large uncertainty in nuclear matrix elements are taken into account.

Since the evidence for abnormal aboundance of ^{128}Xe has later been disproved experimentally, no positive result exists on neutrinoless double beta decay. If the geological results on half lifetimes are taken as limits for neutrinoless $\beta\beta$ decay, one obtains the following limits on the lepton non conserving parameter η :

	^{48}Ca	^{76}Ge	^{82}Se	^{130}Te
No resonance	$10^{-3.6\pm1}$	$10^{-2.8\pm1}$	$10^{-2.4\pm1}$	$10^{-3\pm1}$
Resonance	$10^{-4.1\pm1}$	$10^{-3.5\pm1}$	$10^{-3 \pm 1}$	$10^{-3.7\pm1}$

We would like to note that the limit on ^{48}Ca and possibly others , are probably overestimated due to "unfavourable" nuclear structure . We can conclude from the preceding limits that the lepton non conserving amplitude should anyway be lower than one percent ; such a tiny violation cannot be detected , at present, in any experiment with elementary particles ([23]), and could easily be collocated , as CP non conservation, in a milliweak or a superweak theory. While experiments have been pushed to the limits of technical feasibility , theoretical calculations could be , in my opinion, much improved to take advantage of the impressive limits of many experiments.

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TABLE I.

Transition	A	Z	Isotopic abundance (%)	Transition energy (MeV)	Intermediate transition energy $(A, Z)-(A, Z+1)$	Two-neutrino half-life (years)	Neutrinoless half-lives (years): *a*) No resonance ($\eta = 1$)	*b*) No resonance [$\tau_{\frac{1}{2}}(^{130}Te) = 2.2 \cdot 10^{21}$ y]	*c*) Resonance $\eta = 1$	*d*) Resonance [$\tau_{\frac{1}{2}}(^{130}Te) = 2.2 \cdot 10^{21}$ y]
Ca-Ti	46	20	0.0033	0.985 ± 0.009	-1.382 ± 0.011	$7.4 \cdot 10^{26\pm2}$	$4.5 \cdot 10^{17\pm2}$	$4.7 \cdot 10^{23}$	$1.7 \cdot 10^{16\pm2}$	$5.4 \cdot 10^{23}$
Ca-Ti	48	20	0.185	4.267 ± 0.009	$+0.289 \pm 0.012$	$4.0 \cdot 10^{20\pm2}$	$1.6 \cdot 10^{14\pm3}$	$1.7 \cdot 10^{20}$	$1.2 \cdot 10^{13\pm2}$	$3.8 \cdot 10^{20}$
Zn-Ge	70	30	0.62	1.008 ± 0.006	-0.653 ± 0.009	$2.8 \cdot 10^{26\pm2}$	$3.7 \cdot 10^{17\pm2}$	$3.8 \cdot 10^{23}$	$1.1 \cdot 10^{16\pm2}$	$3.4 \cdot 10^{23}$
Ge-Se	76	32	7.67	2.045 ± 0.004	-0.923 ± 0.004	$3.3 \cdot 10^{23\pm2}$	$9.2 \cdot 10^{15\pm2}$	$1.0 \cdot 10^{22}$	$4.1 \cdot 10^{14\pm2}$	$1.3 \cdot 10^{22}$
Se-Kr	80	34	49.82	0.138 ± 0.007	-1.871 ± 0.005	$1.6 \cdot 10^{35\pm2}$	$3.0 \cdot 10^{21\pm2}$	$3.2 \cdot 10^{27}$	$5.2 \cdot 10^{18\pm2}$	$1.6 \cdot 10^{28}$
Se-Kr	82	34	9.19	3.003 ± 0.008	-0.089 ± 0.008	$5.6 \cdot 10^{21\pm2}$	$1.1 \cdot 10^{15\pm2}$	$1.1 \cdot 10^{21}$	$5.3 \cdot 10^{13\pm2}$	$1.6 \cdot 10^{21}$
Kr-Sr	86	36	17.37	1.240 ± 0.006	-0.0537 ± 0.008	$2.9 \cdot 10^{25\pm2}$	$1.2 \cdot 10^{17\pm2}$	$1.3 \cdot 10^{23}$	$3.7 \cdot 10^{15\pm2}$	$1.1 \cdot 10^{23}$
Zr-Mo	94	40	2.80	1.230 ± 0.005	-0.921 ± 0.014	$3.3 \cdot 10^{25\pm2}$	$1.2 \cdot 10^{17\pm2}$	$1.2 \cdot 10^{23}$	$3.4 \cdot 10^{15\pm2}$	$1.0 \cdot 10^{23}$
Zr-Mo	96	40	17.40	3.364 ± 0.005	$+0.215 \pm 0.025$	$1.7 \cdot 10^{21\pm2}$	$5.1 \cdot 10^{14\pm2}$	$5.4 \cdot 10^{20}$	$2.3 \cdot 10^{13\pm2}$	$7.1 \cdot 10^{20}$
Mo-Ru	100	42	9.62	3.034 ± 0.006	-0.335 ± 0.060	$4.3 \cdot 10^{21\pm2}$	$8.8 \cdot 10^{14\pm2}$	$0.9 \cdot 10^{21}$	$3.8 \cdot 10^{13\pm2}$	$1.1 \cdot 10^{21}$
Ru-Pd	104	44	18.5	1.321 ± 0.012	-1.145 ± 0.008	$1.6 \cdot 10^{25\pm2}$	$7.6 \cdot 10^{18\pm2}$	$8.0 \cdot 10^{22}$	$2.2 \cdot 10^{15\pm2}$	$6.6 \cdot 10^{22}$
Pd-Cd	110	46	12.7	2.004 ± 0.013	-0.868 ± 0.015	$2.5 \cdot 10^{23\pm2}$	$8.2 \cdot 10^{15\pm2}$	$8.7 \cdot 10^{21}$	$2.9 \cdot 10^{14\pm2}$	$8.9 \cdot 10^{21}$
Cd-Sn	114	48	28.86	0.547 ± 0.008	-1.439 ± 0.008	$9.4 \cdot 10^{28\pm2}$	$5.4 \cdot 10^{18\pm2}$	$5.6 \cdot 10^{24}$	$6.1 \cdot 10^{16\pm2}$	$1.9 \cdot 10^{24}$
Cd-Sn	116	48	7.58	2.811 ± 0.006	-0.517 ± 0.024	$8.0 \cdot 10^{21\pm2}$	$1.2 \cdot 10^{15\pm2}$	$1.3 \cdot 10^{21}$	$4.8 \cdot 10^{13\pm2}$	$1.5 \cdot 10^{21}$
Sn-Te	122	50	4.71	0.349 ± 0.008	-1.622 ± 0.008	$8.2 \cdot 10^{30\pm2}$	$4.2 \cdot 10^{19\pm2}$	$4.4 \cdot 10^{25}$	$2.5 \cdot 10^{17\pm2}$	$7.6 \cdot 10^{24}$
Sn-Te	124	50	5.98	2.263 ± 0.007	-0.653 ± 0.008	$6.8 \cdot 10^{22\pm2}$	$4.1 \cdot 10^{15\pm2}$	$4.3 \cdot 10^{21}$	$1.4 \cdot 10^{14\pm2}$	$4.3 \cdot 10^{21}$
Te-Xe	128	52	31.79	0.872 ± 0.007	-1.268 ± 0.010	$7.8 \cdot 10^{26\pm2}$	$5.5 \cdot 10^{17\pm2}$	$5.8 \cdot 10^{23}$	$9.7 \cdot 10^{15\pm2}$	$3.0 \cdot 10^{23}$
Te-Xe	130	52	34.49	2.543 ± 0.008	-0.407 ± 0.030	$1.7 \cdot 10^{22\pm2}$	$2.1 \cdot 10^{15\pm2}$	$2.2 \cdot 10^{21}$	$7.2 \cdot 10^{13\pm2}$	$2.2 \cdot 10^{21}$
Xe-Ba	134	54	10.44	0.731 ± 0.039	-1.328 ± 0.039	$4.3 \cdot 10^{27\pm2}$	$1.3 \cdot 10^{18\pm2}$	$1.3 \cdot 10^{24}$	$1.8 \cdot 10^{16\pm2}$	$5.5 \cdot 10^{23}$
Xe-Ba	136	54	8.87	2.718 ± 0.080	-0.112 ± 0.080	$7.6 \cdot 10^{21\pm2}$	$1.4 \cdot 10^{15\pm2}$	$1.5 \cdot 10^{21}$	$4.8 \cdot 10^{13\pm2}$	$1.5 \cdot 10^{21}$
Ce-Md	142	58	11.07	1.379 ± 0.049	-0.777 ± 0.050	$6.1 \cdot 10^{24\pm2}$	$5.0 \cdot 10^{16\pm2}$	$5.2 \cdot 10^{22}$	$1.2 \cdot 10^{15\pm2}$	$3.6 \cdot 10^{22}$
Nd-Sm	148	60	5.71	1.936 ± 0.021	-0.514 ± 0.030	$1.9 \cdot 10^{23\pm2}$	$8.1 \cdot 10^{15\pm2}$	$8.5 \cdot 10^{21}$	$2.3 \cdot 10^{14\pm2}$	$7.0 \cdot 10^{21}$
Nd-Sm	150	60	5.60	3.390 ± 0.020	-0.036 ± 0.062	$8.0 \cdot 10^{20\pm2}$	$3.6 \cdot 10^{14\pm2}$	$3.8 \cdot 10^{20}$	$1.2 \cdot 10^{13\pm2}$	$3.8 \cdot 10^{20}$
Sm-Gd	154	62	22.61	1.260 ± 0.022	-0.718 ± 0.024	$1.3 \cdot 10^{25\pm2}$	$7.5 \cdot 10^{16\pm2}$	$7.8 \cdot 10^{22}$	$1.6 \cdot 10^{15\pm2}$	$4.9 \cdot 10^{22}$
Gd-Dy	160	64	21.75	1.782 ± 0.030	-0.029 ± 0.030	$2.9 \cdot 10^{23\pm2}$	$1.2 \cdot 10^{16\pm2}$	$1.2 \cdot 10^{22}$	$3.1 \cdot 10^{14\pm2}$	$9.4 \cdot 10^{21}$
U-Pu	238	92	99.275	1.173 ± 0.034	$+0.117 \pm 0.031$	$8.9 \cdot 10^{24\pm2}$	$7.1 \cdot 10^{16\pm2}$	$7.5 \cdot 10^{22}$	$1.1 \cdot 10^{15\pm2}$	$3.3 \cdot 10^{22}$

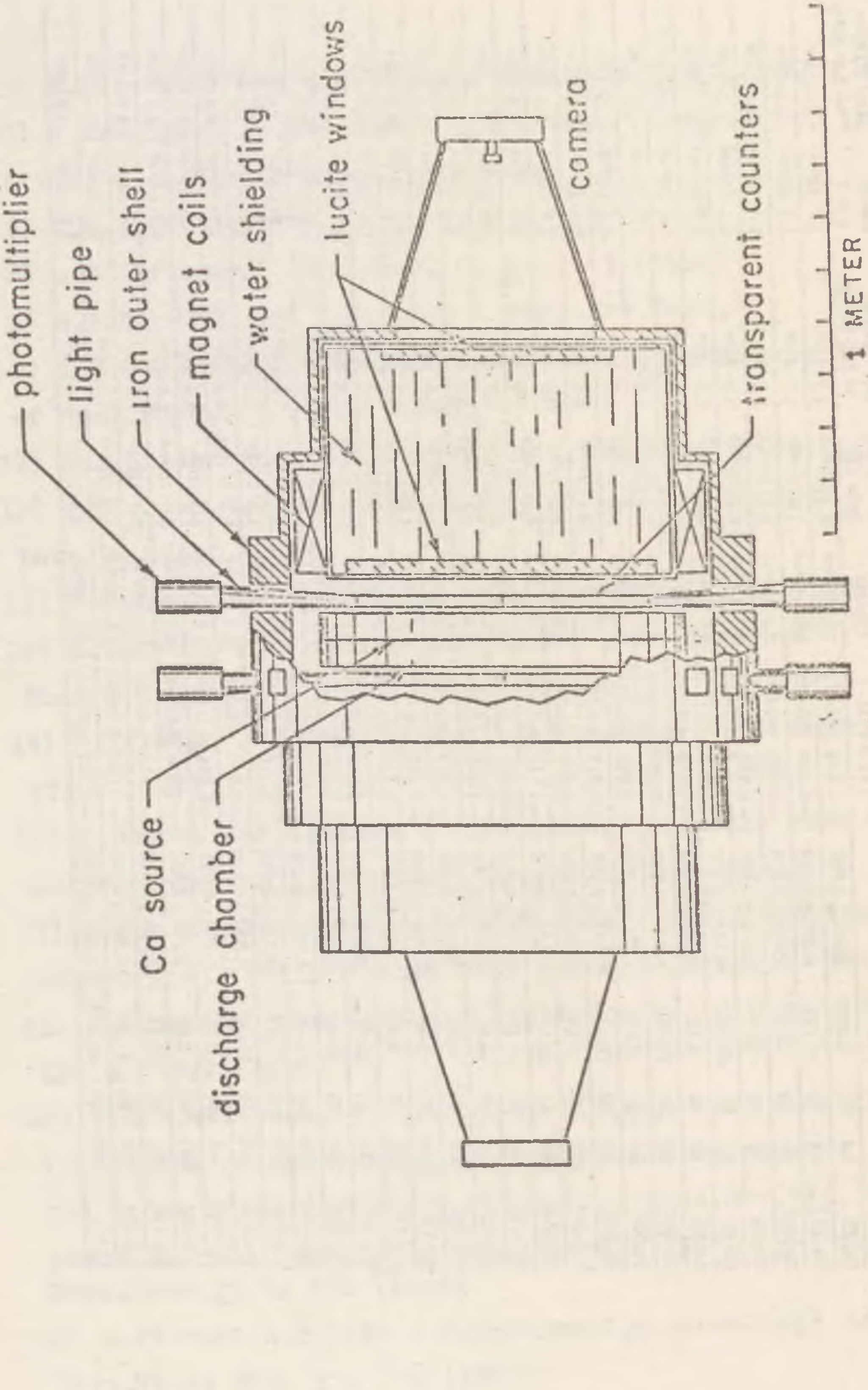

photomultiplier
light pipe
iron outer shell
magnet coils
water shielding
lucite windows
camera
transparent counters
1 METER
Ca source
discharge chamber

Fig. 1

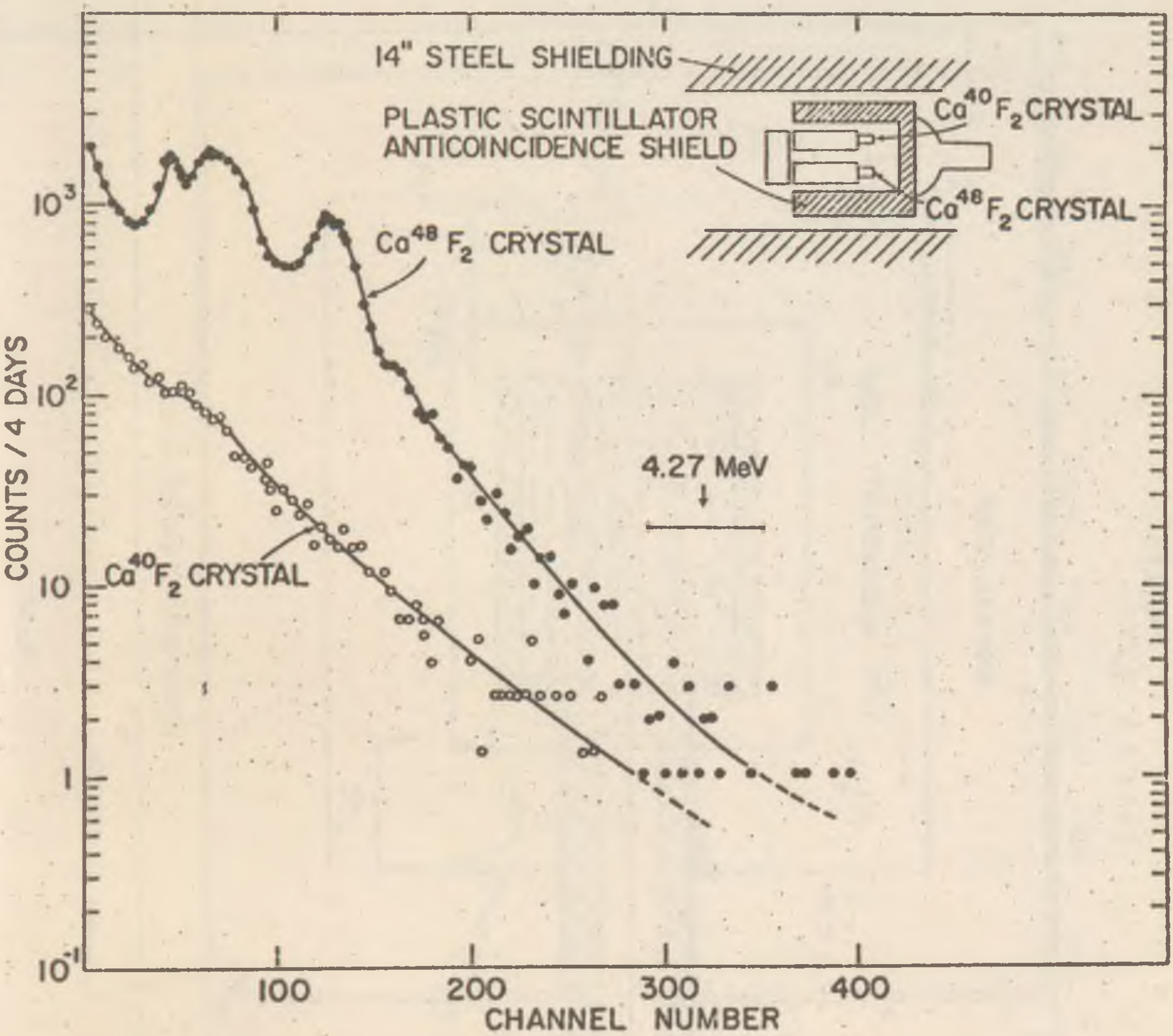


Fig. 2 EXPERIMENTAL RESULTS BY E.DER MATEOSIAN AND G.GOLDHABER

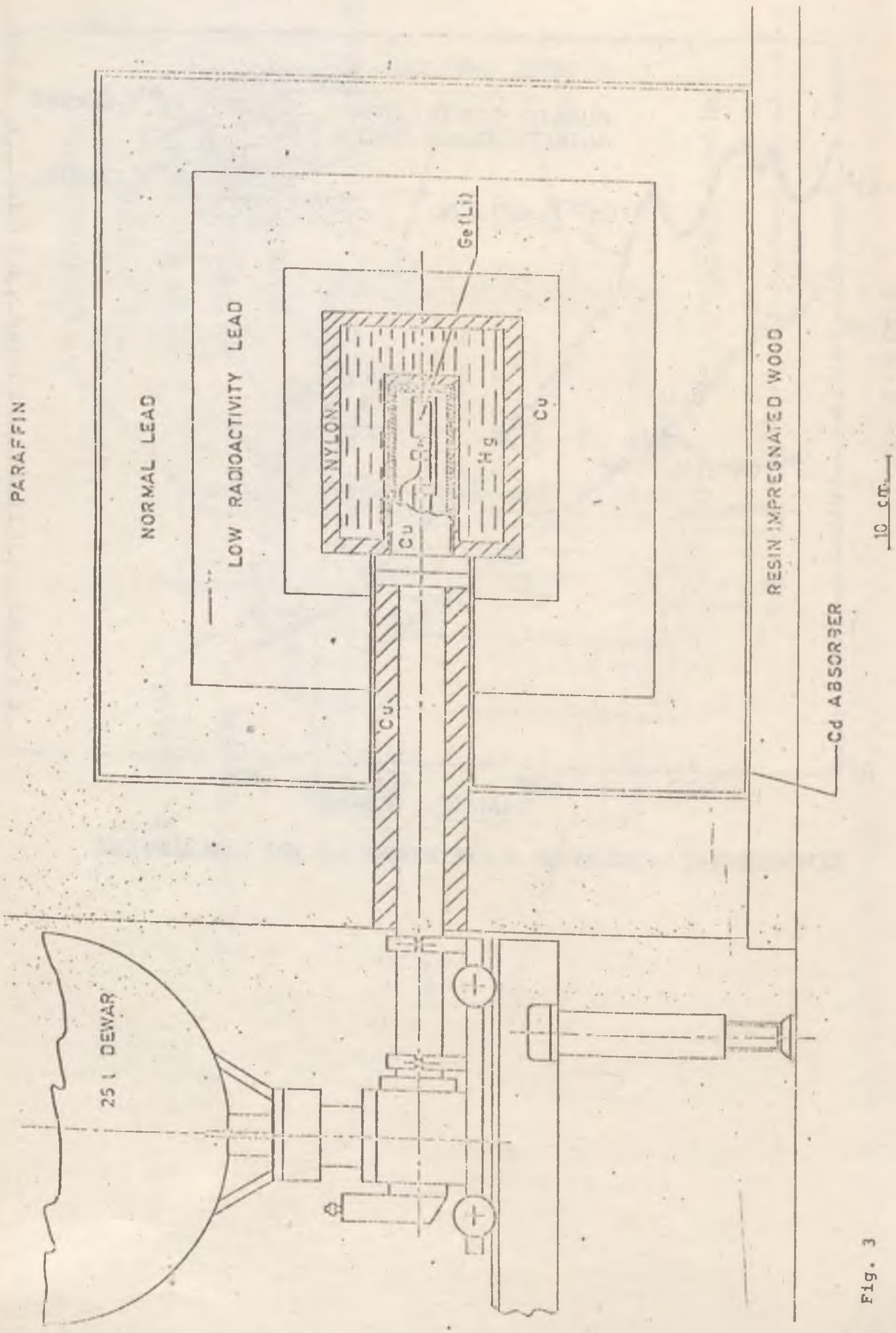
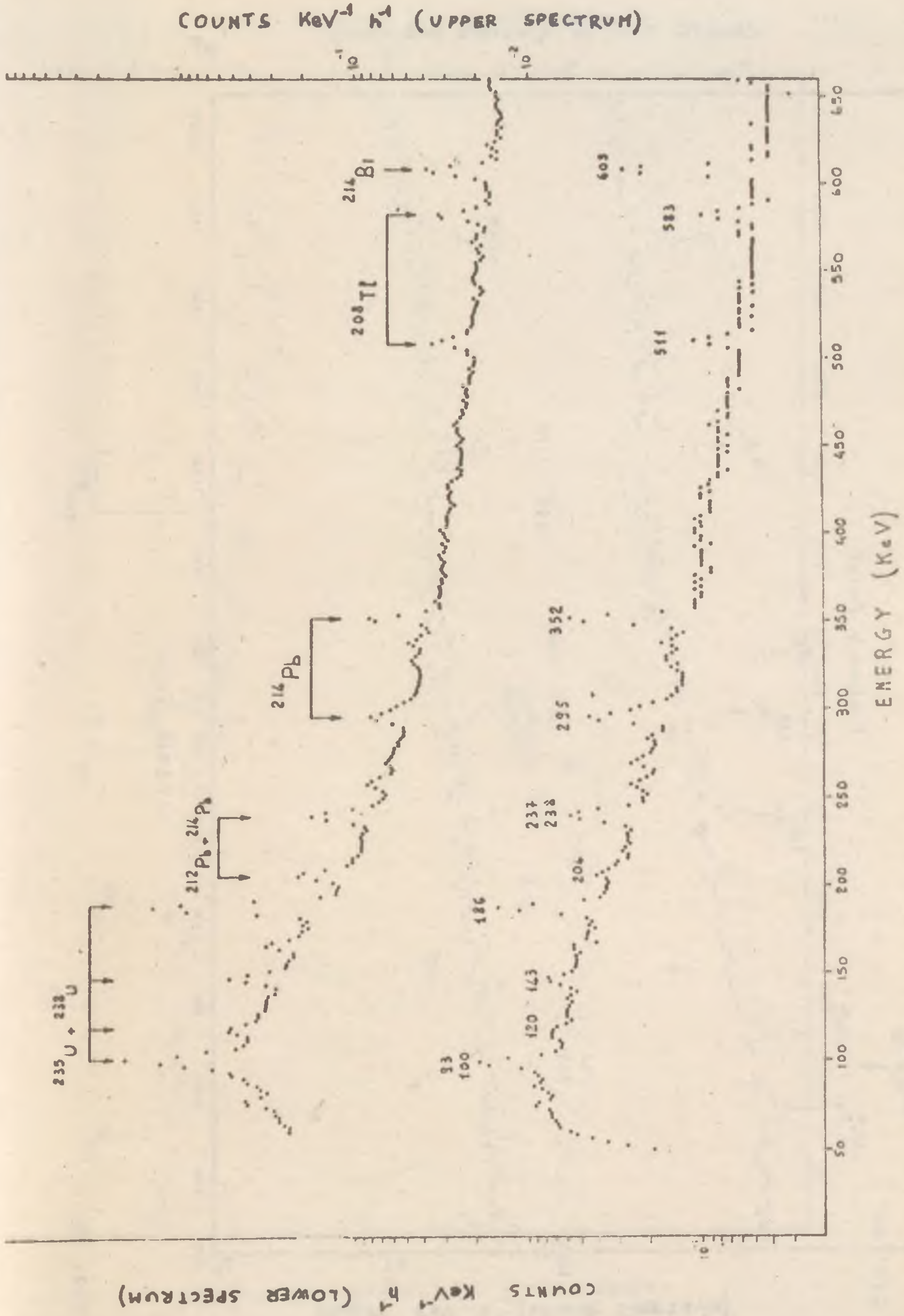

PARAFFIN
NORMAL LEAD
LOW RADIOACTIVITY LEAD
NYLON
Ge(Li)
Cu
Hg
Cu
Cu
RESIN IMPREGNATED WOOD
Cd ABSORBER
10 cm
25 l DEWAR

Fig. 3



COUNTS KeV^-1 h^-1 (UPPER SPECTRUM)
10^-1
10^-2
235U + 238U
212Pb + 214Pb
214Pb
208Tl
214Bi
93
100
120
143
186
204
237
238
295
352
511
583
609
50
100
150
200
250
300
350
400
450
500
550
600
650
ENERGY (KeV)
10^-1
COUNTS KeV^-1 h^-1 (LOWER SPECTRUM)

Fig. 4a

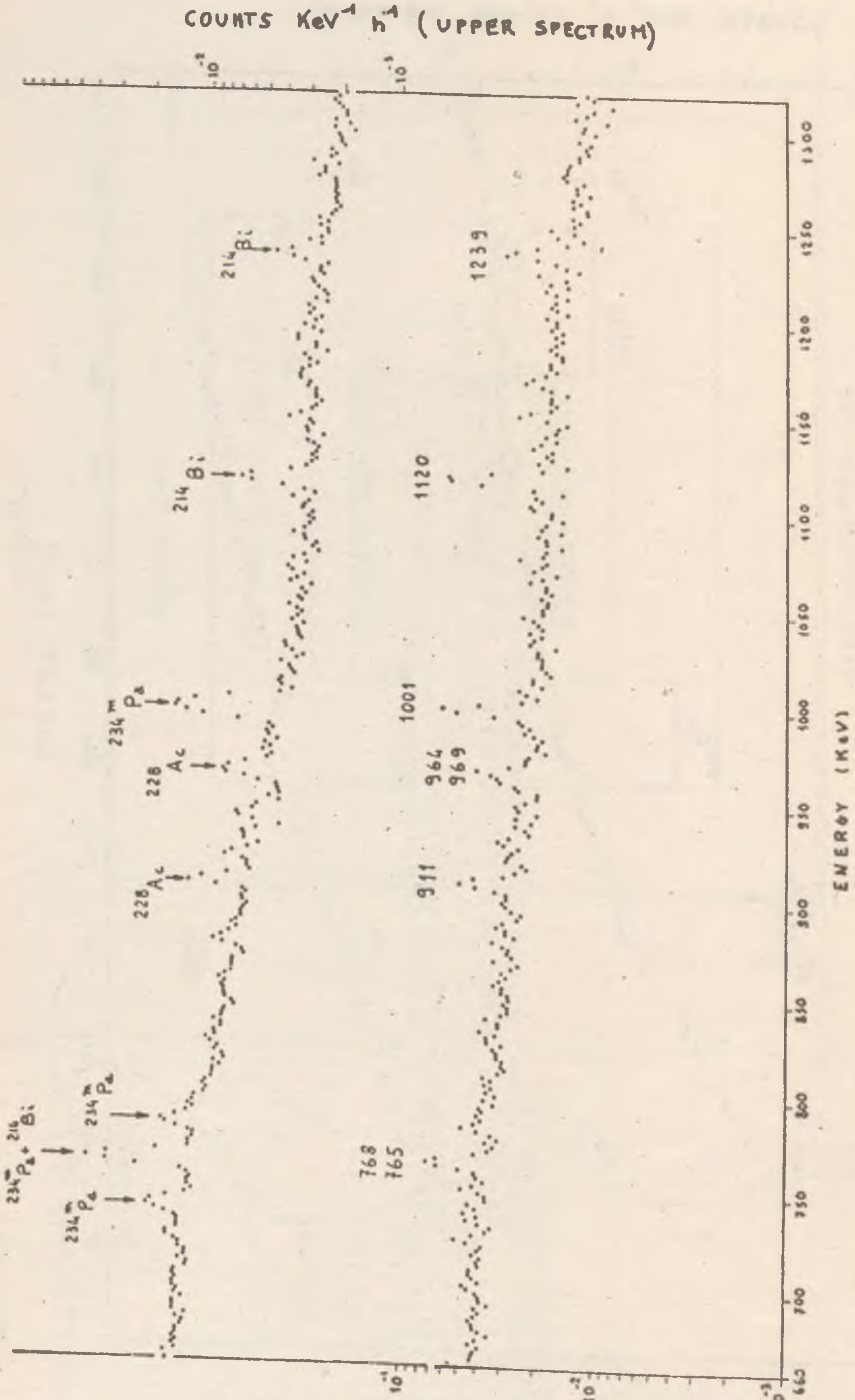

COUNTS KeV^{-1} h^{-1} (UPPER SPECTRUM)
COUNTS KeV^{-1} h^{-1} (LOWER SPECTRUM)
ENERGY (KeV)
^{234m}Pa
^{234m}Pa + ^{214}Bi
^{234m}Pa
^{228}Ac
^{228}Ac
^{234m}Pa
^{214}Bi
^{214}Bi
768
765
911
964
969
1001
1120
1239

Fig. 4b

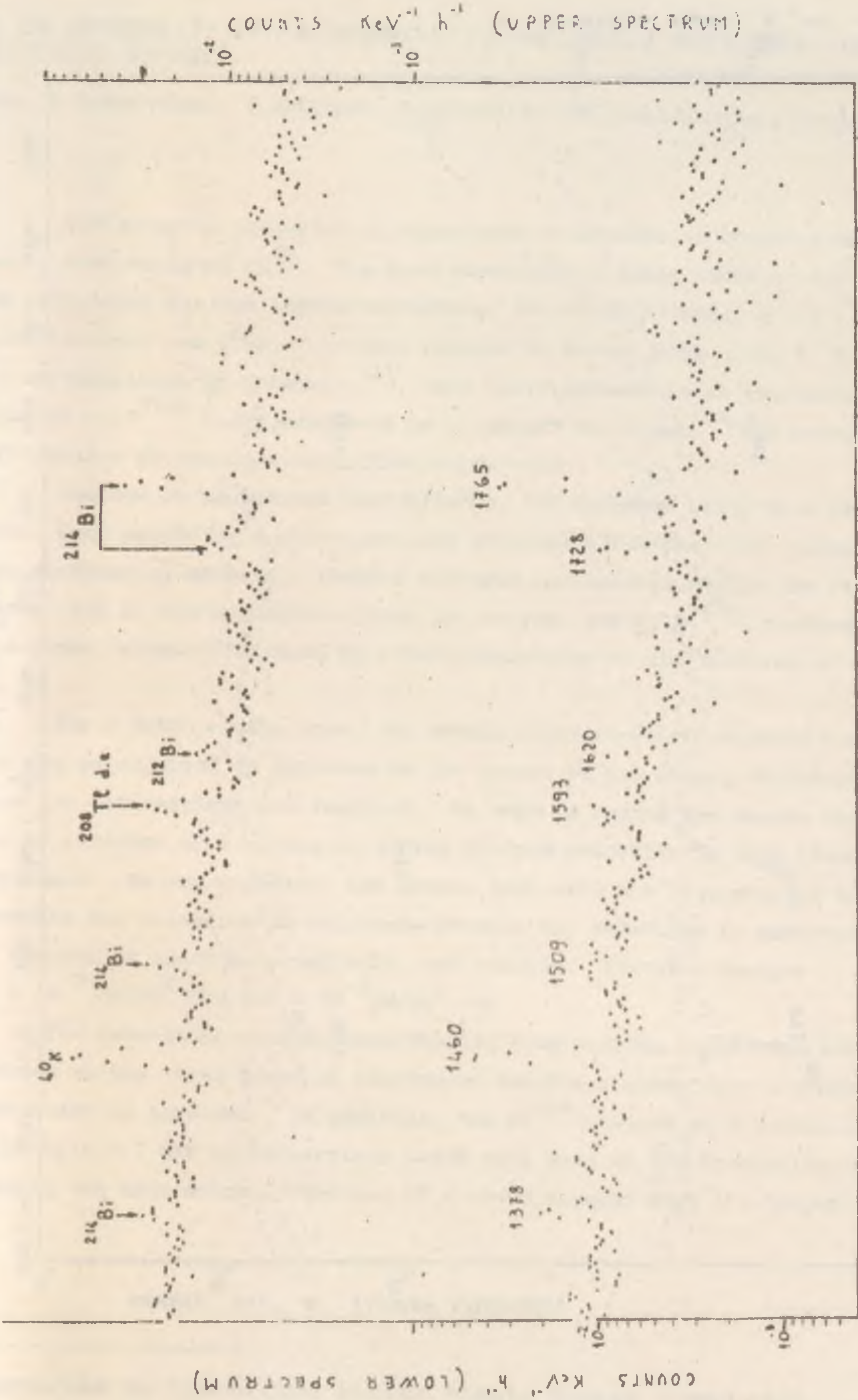
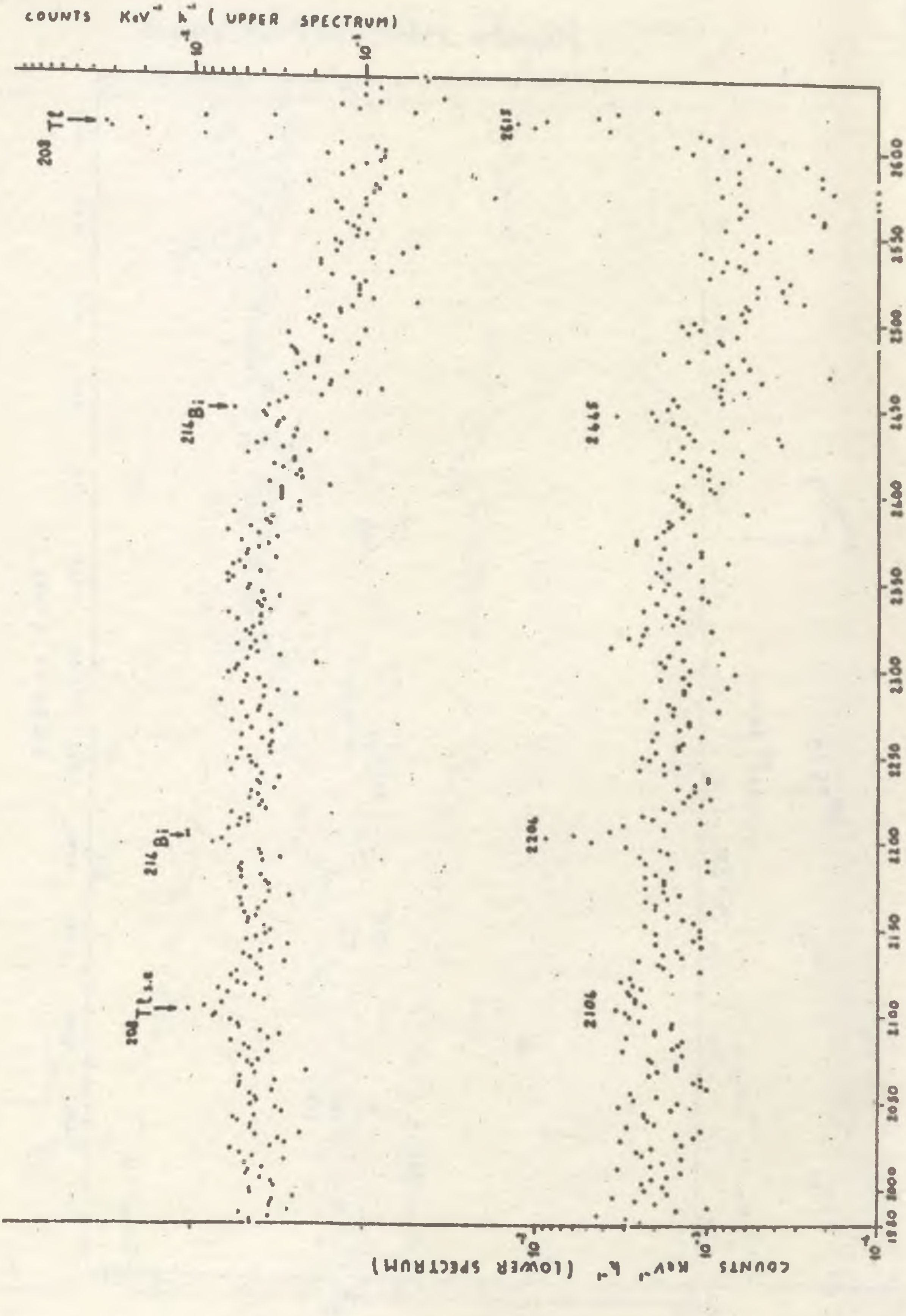

COUNTS KeV⁻¹ h⁻¹ (UPPER SPECTRUM)
COUNTS KeV⁻¹ h⁻¹ (LOWER SPECTRUM)
ENERGY (KeV)
214Bi
40K
214Bi
208Tl d.e
212Bi
214Bi
1378
1460
1509
1593
1620
1728
1765

Fig. 4c


COUNTS KeV⁻¹ h⁻¹ (UPPER SPECTRUM)
COUNTS KeV⁻¹ h⁻¹ (LOWER SPECTRUM)
ENERGY (KeV)
208 Tl
214 Bi
214 Bi
208 Tl s.e
2615
2448
2204
2106
1950
2000
2050
2100
2150
2200
2250
2300
2350
2400
2450
2500
2550
2600

NOTE ON UNIVERSITY OF CALIFORNIA, IRVINE DOUBLE-BETA-DECAY EXPERIMENT* (DISCUSSION REMARK)

M.Moe, D.Lowenthal, F.Reines, University of California, Irvine

Our group is preparing an experiment to attempt an observation of the double beta decay of Ca^{48}. The best experimental lower limit to date on the half life, for the lepton-conserving, two-neutrino mode, $> 3.6 \times 10^{19}$ years has been set with a streamer chamber in a salt mine by C. S. Wu and her collaborators at Columbia.[(1)] This limit encroaches on the theoretical value of $\sim 10^{21\pm2}$ years suggested by Primakoff and Rosen,[(2)] but background difficulties prevented a definitive experiment.

Sources of background recognized by the Columbia group were of two types, each producing double-beta-like pairs of electrons: 1) Gamma rays from surrounding materials induced multiple scatterings within the calcium source, and 2) contamination within the source, mostly Bi^{214}, produced beta-gamma cascades followed by single scattering of the electron or gamma ray.

For a given source area, the double scatterings of external gamma rays can be expected to decrease as the square of the source thickness, since two interactions are required. We hope to reduce the source thickness by a factor of 4 with a resulting 16-fold reduction in this cause of background. In our approach, the source area will not be increased to accommodate the reduction in thickness because the reduction is achieved by the removal of inert material only, not calcium. [Hopeful design: 4.8×10^{-3} gm/cm^2 Ca, 0.2×10^{-3}gm/cm^2 Al]

The beta-gamma cascade contribution from calcium impurities will decrease as the first power of the source thickness since only a single interaction is involved. In addition, the Bi^{214} cascade is followed in 164 μs by a 7.7 MeV alpha particle which will have an 80% probability of escaping the thin source. The use of a cloud chamber with its longer

* Supported by the United States Atomic Energy Commission

sensitive time will make the delayed alpha visible so allowing rejection of most bismuth-produced background.

Since the cloud chamber cannot be cycled much faster than once a minute, it was necessary to devise a triggering scheme to discriminate strongly against unwanted events. We are using a large number of internal proportional counters operating in the cloud chamber gas. Preliminary tests indicate that the trigger logic confines the rate to tolerable levels. Monte Carlo calculations indicate a reasonable efficiency for double-beta detection. Based on this efficiency if the half life is five times the limit set by Wu, we can expect about 150 recorded double beta events per year above a few hundred kilovolts.

The 20cm deep, by 51cm diameter cloud chamber will operate with a He-A filling in a 1500 gauss magnetic field. It appears that the background in our campus laboratory can be adequately controlled with the help of some very clean lead shielding, and salt-mine or underground operation may not be necessary.

The various components are now being assembled and tested. Ten grams of separated Ca^{48} have been obtained on loan from Oak Ridge. We hope to begin gathering data within the next year.

(1) Nucl. Phy. A 158, 337 (1970)

(2) Alpha-Beta- and Gamma-Ray Spectroscopy, K. Siegbahm, ed. (North-Holland, Amsterdam, 1965), 1499.

LEPTON CHARGE CONSERVATION

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After a brief review of the different definitions for the leptonic charge the experimental accuracy of the conservation laws is given. The hypothetical milliweak or superweak violation of the conservation law is discussed. The dependence of the strength of the superweak charge asymmetry on cosmological time is suggested, its astronomical consequences are considered.

The number of the exactly conserved quantities is very limited in Nature. According to our present knowledge one has the ten classical integrals of motion /energy E, momentum $\underline{\underline{P}}$, angular momentum $\underline{\underline{J}}$, speed of center of mass $\underline{\underline{V}}$/. We have also the CPT inversion, related to the former quantities by the celebrated CPT theorem. A completely independent conservation law is that of the electric charge Q. Our generation contributed with the conservation law of the baryonic charge B and - just 20 years ago - the lepton charges L were discovered. After a brief review of the past 20 years I am going to discuss the reopened question: can we be sure indeed that the lepton charge is exactly conserved?

The leptonic charge was introduced by several authors independently and it was distributed among leptons originally in two different ways.

Lepton charges

	γ	ν_e	$\bar{\nu}_e$	ν_μ	$\bar{\nu}_\mu$	e^-	e^+	μ^-	μ^+	hadrons
L	0	+	−	+	−	+	−	+	−	0
L'	0	+	−	−	+	+	−	−	+	0
L_e	0	+	−	0	0	+	−	0	0	0
L_μ	0	+	0	+	−	0	0	+	−	0
P_L		+	+	−	−	+	+	−	−	

$L=L_e+L_\mu$ /Marx [1] 1952/

$L'=L_e-L_\mu$ /Zeldovich [2], Konopinski-Mahmoud [3] 1953/

$P_L = /-1/^{L_\mu}$ /Cabibbo-Gatto [4] , Feinberg-Weinberg [5]/

After the discovery of parity violation in 1957 one learned, that in a certain sense L is the more natural definition: it associates left chirality with positive sign, right chirality with negative sign. On the other hand, L' is independent of chirality, so the L' and $\underline{\underline{J}}$ conservati together result in more restrictive selection rules.

The problem of the two different definitions was solved by proving experimentally in 1962, that essentially both leptonic charges are conserved. The lepton charge conservation splitted into two independent conservation theorems.

It was asked very early, how necessary is to introduce two absolute additive conservation principles, two superselection rules for leptons?

Are these firmly demanded by the experiments? Or can one weaken the stricktness of the lepton charge conservation principle? Three alternatives are worth of discussion:

I. Two additive conservation rules: L_e and L_μ are conserved independently.

II. One additive and one multiplicative conservation rule: L and P_L are conserved.

III. One additive conservation rule: L' is conserved.

Let us have a look at the experimental verifications. The best way of doing this is to compare the coupling constant G' of a lepton non-conserving four-fermion interaction with the G of the conventional lepton conserving weak interaction.

			L	L'	P_L
$\mathcal{H}\to e^-e^-$	in $\beta\beta$ decay	[6,7,8] $G'/G<10^{-10}$	no	no	yes
	in $K^+\to\pi^- e^+ e^+$ decay	[9] $<3.10^{-3}$			
$\mathcal{H}\to e^-\nu_e$	in $\beta\beta$ decay	[6,10] $<10^{-3}$			
	$\bar\nu_e Z\to e^-$... capture	[11] $<0,2$			
$\mathcal{H}\to\mu^-\nu_\mu$	in $\nu_\mu Z\to\mu^+$... capture	[12] $<7.10^{-2}$			
$\mathcal{H}\to e\bar\mu$	in $K_L\to e^-\mu^+$ decay	[13] $<4.10^{-5}$	yes	no	no
	$\mu^- Z\to e^-$... capture	[14] $<5.10^{-4}$			
$e\bar e\to e\bar\mu$	in $\mu\to 3e$ decay	[7,9] $<8.10^{-5}$			
$\mathcal{H}\to e\bar\nu_\mu$	in $\nu_\mu Z\to e^-$ ... capture	[15] $<10^{-1}$			
$\mathcal{H}\to e\mu$	in $\mu^- Z\to e^+$... capture	[16] $<2.10^{-2}$	no	yes	no
$\mathcal{H}\to e\nu_\mu$	in $\nu_\mu Z\to e^+$... capture	[15] $<10^{-1}$			
$\mu\nu_\mu\to e\nu_e$	in $\nu_\mu Z\to\mu^+e^-$ ...capture	[17] <4	yes	no	yes
$ee\to\mu\mu$	$\mu^-e^+\leftrightarrow\mu^+e^-$ (muonium)	[18] <5800			
	$e^-e^-\to\mu\mu$ (coll. beam)	[12] <6100			

/Here $\mathcal{H}$ means the hadrons, i.e. the hadronic currents./

One can conclude from these data, that not only Version I, but also Version II seems to be in accordance with all the experiments. /An L and P_L conserving, but L'violating coupling may be as strong as the conventional weak interaction itself./ The easiest experiment to prove proposal II is searching for an evidence of the $\nu_\mu Z \to Z \mu^+ e^- \bar{\nu}_e$ reaction.

The Version III is in a "centiweak" position: according to the present evidence, by having L' strictly conserved, L may be broken by a $G' < 0{,}01G$ coupling. To check this version, the best way is searching for positrons in μ^- and ν_μ capture reactions.

The conservation law of (at least one)* lepton charge seems to be in an excellent shape from phenomenological point of view. In spite of this brilliant look the possibility of a faint violation of the leptonic charge conservation turned up in the last years. The experimental hints were rather indirect:

a. The life time of ^{128}Te against double β decay was measured to be about 1000 times larger than expected from the theory of the L_e conserving weak interactions [19].

b. The neutrinos, produced at the center of the Sun, do not seem to reach South Dakota [20].

c. A strange new interaction of a very low intensity has been discovered, which violates the time reversal invariance T [21].

By seeing these, the existence of a new faint asmmetric force has been suggested, which violates the T symmetry and the lepton charge conservation at the same time.

1. Primakoff suggested the existence of a milliweak force ($G' = 10^{-3} G$), which violates T symmetry in the $\Delta S = 1$ channel and violates L conservation in the $\Delta Q_{lepton} = 1$ channel [10].

2. Pontecorvo suggested the existence of a superweak force $(G'=10^{-10}G)$, which violated T symmetry in the $\Delta S=2$ channel and violated L conservation in the $\Delta Q_{lepton}=2$ channel [8,22].

It is evident that a milliweak or superweak breaking of L conser-tion is not excluded by the present experimental evidence. It must not be overlooked, that the conservation of the electric charge Q is verified only up to the same accuracy. The life time of the electron has been found to be of the order or longer than 10^{21} years [23] and this is just the order of magnitude of the double β decay life times. If we assume a hypothetical $e\nu \leftrightarrow \nu\nu$ coupling, its strength is limited with similar accuracy: $G'' \lesssim 10^{-10}G$ (See also [7].) If one wants to prove or disprove the possibility of a simultaneous superweak T, L and Q breaking, one should increase the knowledge about the electron stability, too. (The conservation of the baryonic charge B has been verified up to a much higher accuracy [24].)

The experimental indications for a superweak violation of L_e conser-vation became fainter in the last monthes. Kristen remeasured the ^{128}Te life time and it turned out to be definitely longer ($t > 10^{23,3}$ instead of $t=10^{22,5\pm0,5}$ years) [6]. Davis found the neutrino rays of the Sun fainter [25] than predicted by the $\nu_e \rightleftarrows \nu_\mu$ oscillation theory [23,26]

This new development does not mean, that the solar neutrino puzzle has been solved. It became even more puzzling than earlier. The instability of ν_e still offers the simplest explanation:

α. ν_e decays into lighter particles: $\nu_e \rightarrow \nu_2$ +x [27]

β. ν_e oscillates between several ν states: $\nu_e \leftrightarrow \nu_\mu \leftrightarrow \nu_2 \leftrightarrow \ldots$ [28]

The solar neutrino puzzle cannot be solved by giving up only the lepton charge conservation. We would need new particles, even if L_e could change.

*

Instead of giving up the strict lepton charge conservation, an alternative astrophysical explanation of the faint solar neutrino flux could be the low He content of the original Sun [29]. y=10% He concentration would give a ν intensity, which were in accordance with the Davis experiment. (The assumption of the unstable neutrino predicts ~ 0 SNU, the assumption of the low He concentration predicts cca 1 SNU rate for the Davis experiment.*) This low He concentration is, however, in sharp contradiction with the conventional cosmology, which supposes a Hot Universe 10 billion years ago.

A low He concentration would result from the Big Bang if one attributed a high nonvanishing (positive) lepton charge to the universe [30].

$$L \simeq 10^{8} B > 0$$

could explain even y=1%. So we have two charge asymmetries in Nature:

i. A superweak charge asymmetry in the physical interactions, observed now in lab in the K_L decay.

ii. A giant charge asymmetry in the state of our world /$L \gg B > 0$/, which originated 10^{10} years ago.

Can these two phenomena have the same root? Let us suppose, that the strength of the charge asymmetric coupling is decreasing in cosmological time, e.g.

$$G' \approx \frac{1}{t},$$

as in the Dirac-Dicke theory of gravity or in certain theories of weak interactions [31]. The charge asymmetric force is superweak today, in our oldish world, but it might be wildly strong in the first moments of the fresh universe. An interference with the charge symmetric strong nuclear force could result in a faster concentration and evolution of antimatter, than that of matter. It might happen, that at the dawn of the universe much more black holes were formed from antimatter. In these black antiholes a tremendous amount of negative baryonic charge and lepton charge disappeared [33]. What survived this charge asymmetric era, was a universe with positive effective L and B . The present Sun is dark for our neutrino eyes, because it is poor in He, and this astronomical fact may be related to the C and T symmetry breaking (but L conserving) force, which is superweak now, but which might be strong long ago.

The change in the value of G' can be explained with a T-odd scalar field $\varphi(x) = -T^{-1}\varphi(x)\,T$ in the Hamiltonian, e.g. in

$$H' = i\,G_0\,\varphi(x)\left[\bar{\psi}_1(x)\,\psi_2(x) - \bar{\psi}_2(x)\,\psi_1(x)\right],$$

which possesses a nonvanishing expectation value [34]

$$G_0 \langle 0|\varphi(x)|0\rangle = G^* \neq 0.$$

G^* is a combination of masses and (dimensionless) coupling constants, having the dimension sec^{-1}. Once upon a time the universe was so hot, that all the particles were extrem relativistic, the world was scaling invariant [35]. This time the T-odd quantity G^* had to be of the form

$$G^* \simeq \frac{G_0}{t} + \dots$$

where t is the cosmological time.

It is hard to prove the - prehistoric - time dependence of the superweak coupling. One indication would be the observation of a scalar particle with charge parity $C = -1$ (consequently with a very complex quark structure). Its final states can be only complicated many - pion - - states, penetrating high centrifugal barriers, like

$$\bar{q}\gamma_\mu q \cdot \partial_\mu \bar{q}q \sim \omega_\mu \partial_\mu \varepsilon$$

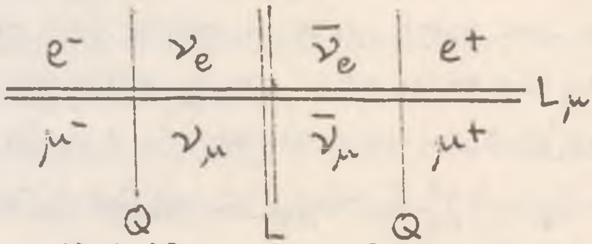
$$\bar{q}\gamma_5\gamma_\mu q \cdot \bar{q}\gamma_5\partial_\mu q \sim (\underline{A}_{1\mu} \cdot \underline{B}_\mu)$$

$$\bar{q}\gamma_\mu q \cdot \bar{q}\gamma_5 q \cdot \partial_\mu \bar{q}\gamma_5 q \sim (\underline{\pi} \cdot \underline{\rho}_\mu)\partial_\mu \eta.$$

It may be worth to search for such an exotic resonance.

*

The world of the hadrons is a very wild and alive jungle, in sharp contrast with the cool land of leptons. The eight leptons are separated by unpenetrable superselecting barriers.



It is a bit curious that if we were able to add two neutrinos with vanishing total E, $\underline{\underline{p}}$ and $\underline{\underline{J}}$ to our world, we would arrive at a completely strange new world, isolated from ours for eternity. Do the L_e and L_μ conservation laws generate superselection rules indeed? Or is there a very narrow path, connecting these nearby and still faraway state vector spaces? Should we long for a conserved or for a very slowly changing lepton charge? Or is the other question the right one: is the charge asymmetry of our world a consequence of an unchangeable initial condition, or is its original charge asymmetry hidden behind the lace of the dynamical evolution?

T.D.Lee told once [36], that in certain cases a broken symmetry produced a more estetic impression, (Fig. 1) than the boring exact symmetry. Let me dare to mention an other thesis: in certain cases a hidden symmetry may be more attractive, than an open one (Fig. 2).

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FOOTNOTE

* An other possible experiment to look for neutrinos with a shorter mean path would be the detection of terrestrial neutrinos [30].

Fig. 1

Fig. 2

NOTE ON ELECTRON STABILITY AND THE PAULI PRINCIPLE*
(DISCUSSION REMARK)

F.Reines and H.W.Sobel, University of California, Irvine

It was most interesting to learn from Dr. Marx that an extension of the old limit for electron stability obtained by Moe and Reines[1] would have important theoretical implications because as it happens one of us (H. W. Sobel) has been actively pursuing such an improvement - by perhaps a factor of ten.

An amusing aside is the possibility pointed out by M. Goldhaber that these same experiments can be interpreted in terms of a test of the Pauli principle for atomic electrons. The idea is simply that a vacancy associated with a violation of the Pauli principle would allow two electrons with identical quantum numbers to occupy the same orbit so permitting, say, the emission of K X-rays. A limit on the number of K X-rays observed, in fact, gives an exquisitely sensitive test of the Pauli principle in this case. The result is given in terms of an upper limit on the relative strength of the Pauli violating term to the strength of the electromagnetic coupling $\left(\frac{g_{pv}}{g_{em}}\right)^2 < 5 \times 10^{-44}$.

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* Work supported by the United States Atomic Energy Commission

1. M. K. Moe and F. Reines, Phys. Rev. 140, 4B, B992 (1965). We found a lower limit of $> 2 \times 10^{21}$ yr or $> 4 \times 10^{22}$ yrs depending on the mode of decay assumed.

REMARK ON STABILITY OF $\bar{\nu}_e$*
(DISCUSSION REMARK)

F.Reines, University of California, Irvine

If, as has been suggested by various authors [1] the ν_e has a non-zero rest mass, then we must allow the possibility that it might be unstable . There is, of course, a wide range of possible decay modes which might be postulated but work at Irvine by H. S. Gurr, H. W. Sobel and myself rules out at least one of the postulated decay modes, i.e.,

$$\bar{\nu}_e \to \bar{\nu}' + \gamma \qquad (1)$$

The limit for this decay mode to a lighter $\bar{\nu}'$ and a gamma ray has been set in a most conservative way by ascribing the entire reactor associated background in our extremely low background detector at the Savannah River Reactor to the process (1). We interpret our results subject to the assumption that $m_{\nu'}/m_{\nu_e} << 1$ and conclude from our data that the decay length is $> 1.5 \times 10^{18}$cm, or $> 10^{-5}$ astronomical units.

For this reason we conclude that (1) can be ruled out as the cause of the unexpectedly low observation of solar neutrinos reported by Davis. We recognize of course that we can make little comment as to other decay modes which may give rise to particles more difficult to detect than gamma rays.

In the most general case, a comparison in reactor experiments between the predicted and observed inverse beta rate shows that, whatever the cause, there is an agreement which indicates a $\bar{\nu}_e$ decay length $>$ 100 meters but this more general result falls far short of testing the stability on the astronomical scale required for a coagent comment on the dilemma posed by the solar neutrino experiment of Davis.

(A more detailed description of the experiment which leads to this result for (1) is in preparation for submission to Phys. Rev. Letters).

* Sponsored by the United States Atomic Energy Commission

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NEUTRINO EXCITATION OF NUCLEAR LEVELS IN ^{12}C*

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(Presented by Clyde L. Cowan, Catholic University)

Neutrino induced reactions will be carried out both at low energies ($E_\nu \leq$ 53 MeV, using neutrinos from the decay of stopped muons at high-intensity meson facilities such as LAMPF, Los Alamos, N.M.) and at high energies (several GeV, obtained at high-energy accelerators such as at NAL, Batavia, Ill.). While elementary particle reactions will be of primary interest, neutrino-induced reactions on complex nuclear targets will be considered also. Besides being of interest per se, such reactions occur as a background in all neutrino processes where complex nuclei are present in the counters and/or the target material, and consequently their cross sections have to be known. We present here the results of a calculation predicting the neutrino excitation of the T = 1 levels of ^{12}C, via the reaction

$$\nu + {}^{12}C \rightarrow {}^{12}N^{(*)} + \ell^- \qquad (1)$$

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*Talk Presented at the "Neutrino '72" conference, 11-17 June 1972, Balatonfüred, Hungary.

+Supported in part by a grant of the National Science Foundation.

where ℓ = electron or muon. The threshold of the reaction is

$$E_\nu^{thresh} = \omega + m_\ell - m_e \tag{2}$$

with ω the level energy of ^{12}N measured from the ^{12}C ground state. The energy of the emitted lepton is

$$E_\ell = E_\nu - \omega + m_e \tag{3}$$

and one has a momentum transfer to the nucleus $\vec{q} = \vec{p}_\nu - \vec{p}_\ell$. The non-relativistic Hamiltonian

$$H = \langle \phi_p \psi_\ell | (G_F + G_{GT}\,\vec{\sigma}\cdot\vec{\sigma}^N) | \psi_\nu \phi_n \rangle \tag{4}$$

is taken to contain just the large terms (Fermi and Gamow-Teller), with $G_F = 10^{-5} m_p^{-2} f_p(q_\mu^2)$, $G_{GT} = -1.20\, G_F$, with f_p the proton form factor. We obtain a differential cross section

$$\begin{aligned} d\sigma/d\Omega_\ell = (p_\ell E_\ell/4\pi^2)\Big\{ & G_F^2 \left(1 + \vec{p}_\nu\cdot\vec{p}_\ell/E_\nu E_\ell\right) |\mathcal{M}|^2 \\ & + G_{GT}^2 \left[2\,\mathrm{Re}\,\vec{\mathcal{M}}^*\cdot\vec{p}_\nu\,\vec{\mathcal{M}}\cdot\vec{p}_\ell/E_\nu E_\ell + |\vec{\mathcal{M}}|^2\left(1 - \vec{p}_\nu\cdot\vec{p}_\ell/E_\nu E_\ell\right)\right]\Big\} \end{aligned} \tag{5}$$

in terms of scalar ($\mathcal{M}$) and vector ($\vec{\mathcal{M}}$) nuclearmatrix elements that are functions of $\vec{q}$. They can be expressed by "transition densities" $\rho(\vec{r})$, $\vec{\rho}(\vec{r})$, e.g.

$$\mathcal{M}(\vec{q}) = \int d^3r\, e^{i\vec{q}\cdot\vec{r}} \rho(\vec{r}), \tag{6}$$

for which we perform a multipole expansion,

$$\rho(\vec{r}) = \sum_{LM} (J_i M_i, LM | J_f M_f)\, \rho_L(r)\, Y_{LM}^*(\hat{r}) \tag{7}$$

corresponding to nuclear states $J_iM_i \to J_fM_f$, so that

$$\mathcal{M}(\vec{q}) = 2\pi\sqrt{2}\sum_{LM}(J_iM_i, LM|J_fM_f)\, I_L(q)\, Y^*_{LM}(\hat{q}) \tag{8}$$

$$I_L(q) = i^L \int r^2 j_L(qr)\, \rho_L(r)\, dr \tag{9}$$

For the vector matrix element, one obtains similarly

$$\vec{\mathcal{M}}(\vec{q}) = 2\pi\sqrt{2}\sum_{LL'M}(J_iM_i, LM|J_fM_f)\, I_{LL'}(q)\, \vec{Y}^{M*}_{LL'}(\hat{q}), \tag{10}$$

$$I_{LL'}(q) = i^{L'} \int r^2 j_{L'}(qr)\, \rho_{LL'}(r)\, dr \tag{11}$$

where L' = L, L ± 1. In terms of the "form factors" $I_L(q)$ and $I_{LL'}(q)$, one obtains e.g. for electric multipole transitions EL:

$$\begin{aligned}(d\sigma/d\Omega_\ell)_{EL} = {} & (2p_\ell E_\ell/\pi)\Big\{G_F^2\,(1+\vec{p}_\nu\cdot\vec{p}_\ell/E_\nu E_\ell)\,|I_L(q)|^2 \\ & + G_{GT}^2\Big[(1-\tfrac{1}{3}\vec{p}_\nu\cdot\vec{p}_\ell/E_\nu E_\ell) + (-1)^L\sqrt{6}\,(\vec{p}_\nu\cdot\vec{q}\;\vec{p}_\ell\cdot\vec{q}/q^2 - \tfrac{1}{3}\vec{p}_\nu\cdot\vec{p}_\ell)/E_\nu E_\ell \\ & \times (2L+1)(L0,L0|20)\,W(LL11;2L)\Big]\,|I_{LL}(q)|^2\end{aligned} \tag{12}$$

It can be shown that the multipole transition density $\rho_L(r)$ is the same quantity that describes the Coulomb excitation of nuclear levels by electrons, while $\rho_{LL'}(r)$ is related by

$$\rho_{LL'}(r) = \frac{2m_p}{\mu_p-\mu_n}\,\mu_{LL'}(r) \tag{13}$$

to the magnetization density $\mu_{LL'}(r)$ in the theory of electroexcitation, with $\mu_p - \mu_n = 4.70$ the difference between p,n magnetic moments. The electroexcitation densities ρ_L and $\mu_{LL'}$ can be approximated by delta-functions at a transition radius R, convoluted with a Gaussian surface smearing ("generalized Helm model"[1]). We have used this model to fit the measured electroexcitation cross sections of the ^{12}C levels, and have thus determined the strength parameters in the Helm model for each level. Subsequently, the same model has been used to predict the differential and total (integrated over lepton emission angles) ^{12}C-neutrino cross sections with excitations of nuclear levels.

It was found that at low neutrino energies, the largest contribution to the neutrino cross section was the M1 transition to the ^{12}N ground state (analog of the 15.1 MeV level in ^{12}C) while at high energies, the E1 giant resonance predominated. Our fits to the M1 and E1 electroexcitation form factors of the corresponding levels in ^{12}C are presented in Figs. 1 and 2, respectively. The data are taken from the experimental paper of Yamaguchi[2], and Fig. 2 presents the summed dipole strength over the regions of 21-26 MeV and 21-37 MeV excitation of ^{12}C.

The predicted low-energy neutrino total cross sections for $\ell = e$ are shown in Fig. 3. Level No. 1 is the 1^+ ground state of ^{12}N ($\omega = 17.34$ MeV), seen to dominate below $E_\nu = 60$ MeV. "Levels" No. 18 and 19 correspond to the summed dipole strengths shown in Fig. 2 (No. 19 is the 21-37 MeV sum) and dominate above 60 MeV, while all other ^{12}N levels below the giant resonance are small. The high energy neutrino cross sections are shown in Fig. 4 ($\ell = e$) and Fig. 5 ($\ell = \mu$),

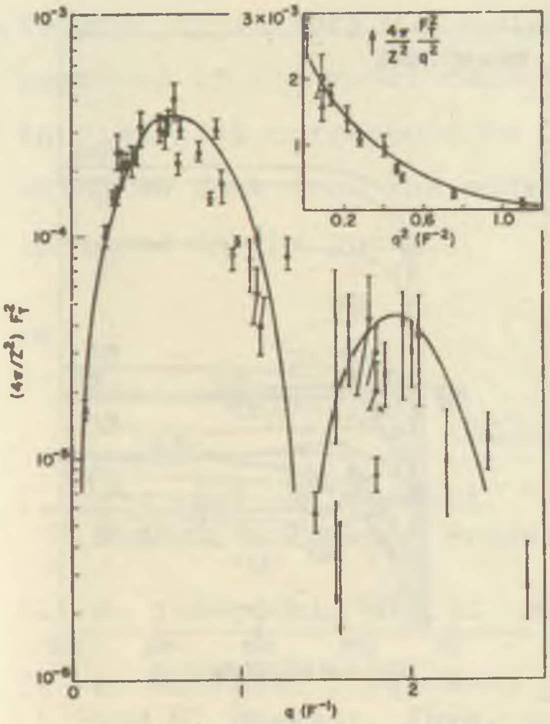

Fig. 1 M1 electroexcitation form factor of the 15.1 MeV level in ^{12}C, fit by the Helm model.

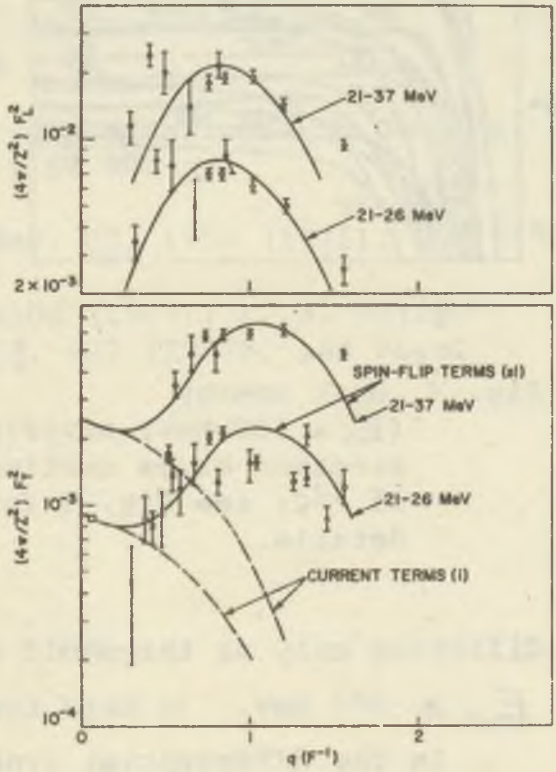

Fig. 2 Longitudinal (top) and transverse (bottom) E1 electroexcitation form factor of ^{12}C giant resonance, summed over 21-26 MeV or 21-37 MeV region as indicated and fit by the Helm model.

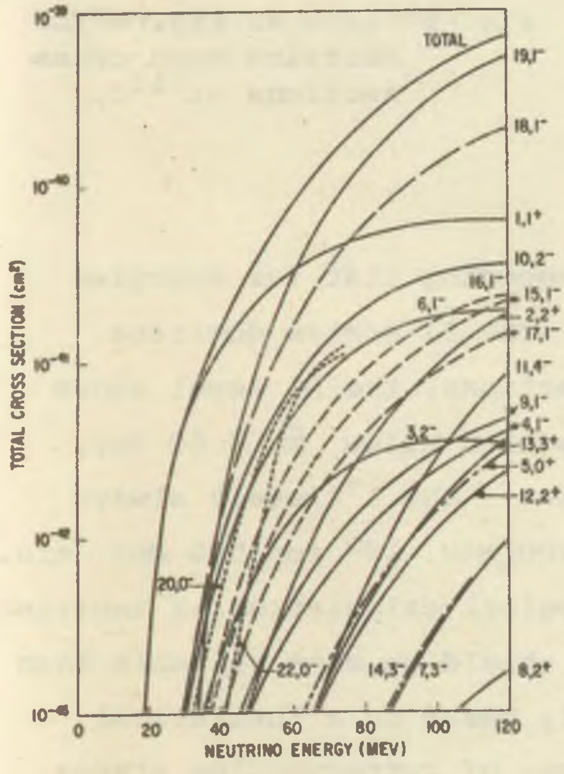

Fig. 3 Low energy ($E_\nu \leq 120$ MeV) neutrino excitation cross sections of ^{12}C leading to ^{12}N levels. Dashed curves are individual states included in the summed 1^- level No. 19. Dotted curves are 0^- levels, given by theory only and not included in the total.

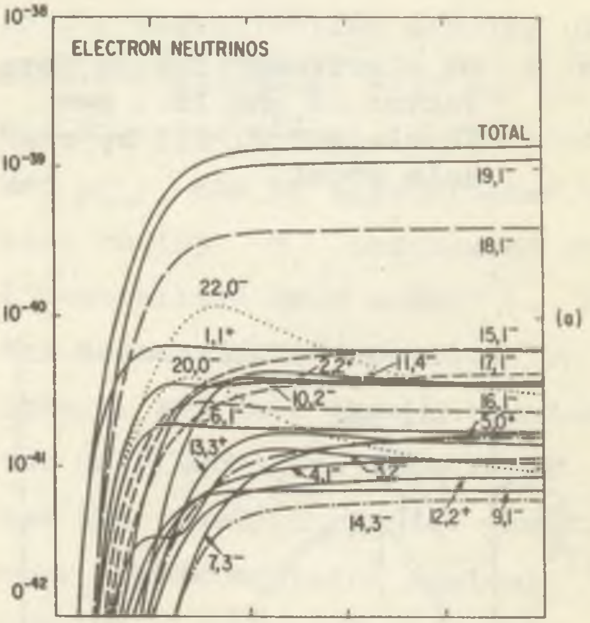

Fig. 4 High energy ($E_\nu \leq$ 500 MeV) neutrino-electron cross sections of ^{12}C; see Fig. 3 for details.

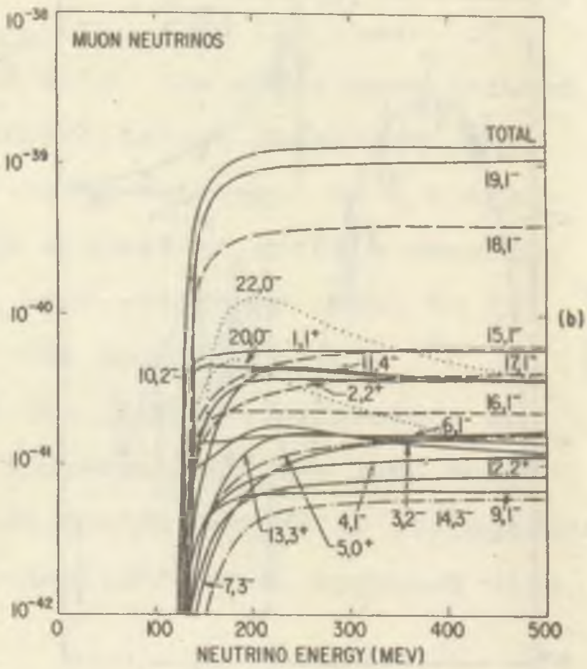

Fig. 5 Same as Fig. 4 for neutrino-muon cross sections of ^{12}C.

differing only at threshold and becoming flat for energies $E_\nu \geq$ 300 MeV. In this region, the E1 states dominate.

In the differential cross sections, the M1 level shows some backward emission of the electron below E_ν = 60 MeV, but peaks at 0° at high energies. The 1^- levels always peak forward (at 30° for E_ν = 300 MeV, 17° for 500 MeV, etc.)

The present semi-phenomenological calculation of neutrino excitation of ^{12}C nuclear levels should be more reliable than previous work[3] which was entirely based on a theoretical nuclear model. Indeed, comparisons of corresponding states show that the realistic calculation predicts cross sections reduced by factors 0.4 - 0.6 (or even 0.1 for a 2^- level) as compared to the model calculation. (The dotted cross sections in Figs. 3-5 correspond to 0^- states not excited by electrons, which we took from the model calculation[3]. They are not included in the total.)

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ANTINEUTRINO-ELECTRON SCATTERING [x]

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INTRODUCTION AND HISTORICAL BACKGROUND

The search for the antineutrino-electron elastic scattering process

$$\bar{\nu}_e + e^- \longrightarrow \bar{\nu}_e + e^- \qquad /1/$$

has proceeded since Fermi's /1/ pioneering paper on beta decay appeared in 1934. It quickly became evident that the process was not going to be an easy one to observe. In 1935 Nahmias, /2/ using a natural Radium source looked for the process using a geiger counter as both target and detector. His lower limit oh the path of neutrinos in air /3×10^{10} cm/ was interpreted by Bethe /3/ as presenting an upper limit on the magnetic moment of the neutrino at less than 2×10^{-4} Bohr magnetons /Bm/. Unlike the inverse beta process there was no theoretical basis on which to calculate the magnitude of a magnetic moment. However in 1935, such an additional feature was not a priori unreasonable.

in 1950, Barrett /4/ used tritium in an experiment similar to that of Nahmias and obtained an upper limit of 10^{-5} Bm. Also about this time Wollan /5/ took advantage of the much larger neutrino flux emitted from

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[x] Supported by the United States Atomic Energy Commission

a uranium pile and was able to give an upper limit for the interaction cross-section of fission neutrinos as $< 10^{-31}$ cm^2. This implied /6/ that the experimental lower limit on the path of neutrinos in air was now 2×10^{12} cm or an upper limit on the magnetic moment of 10^{-6} Bm.

Earlier, Crane /7/ had used geophysical arguments involving solar neutrinos to arrive at an upper limit of 2×10^{-7} Bm.

In 1954, Cowan, Reines and Harrison /8/ working at a large fission reactor with a liquid scintillator were able to set an experimental upper limit of 10^{-7} Bm. In 1957, the same experimental location but with a larger and better-shielded detector allowed Cowan and Reines /9/ to set a new upper limit of 10^{-9} Bm.

About this time, Salam, /10/ Landau /11/ and Lee and Yang /12/ published the two-component neutrino theory in which a theoretical basis was provided for distinguishing neutrinos and anti-neutrinos. It was pointed out by Salam/10/ that according to this theory the neutrino magnetic moment must be zero.

It wasn't until 1958 with the appearance of the V-A theory Marshak and Sudarshan /13/ and Feynman and Gell-Mann /14/ that experimentalists once again had a theoretical picture to describe the elastic scattering process. The success of this theory in describing the inverse beta decay process or "non-diagonal" terms of the weak interaction Hamiltonian made it natural to assume the theory's validity for the "diagonal," or elastic scattering, terms also.

Analyzed in this new light, the data of Cowan and Reines /9/ implied that the square of the ratio of the coupling constant for elastic scattering $[g^2]$ to the beta-decay constant $[g_\beta^2]$ was $< 10^2$.

No new data were to be forthcoming until 1970. Before that, however, in 1969, Gell-Mann, Goldberger, Kroll and Low /15/ suggested that the "diagonal" and "non-diagonal" terms of the weak interaction Hamiltonian may be of a quite different character. If this is true no valid prediction as to the cross-section of the elastic scattering process can be made on the basis of V-A and we would have to abandon universality.

In 1970, Stothers /16/ using astrophysical data concluded that the cross-section was $< 10^{0\pm 2}$ times V-A predictions and Steiner /17/ using CERN neutrino experiment data gave an upper limit of 40 times V-A, for the related reaction

$$\nu_e + e^- \longrightarrow \nu_e + e^- . \qquad /2/$$

Also in 1970, Reines and Gurr, /18/ working at a reactor, set a new upper limit of < 4 times the V-A prediction.

Meanwhile, evidence has been mounting that a unified theory of weak and electromagnetic lepton interactions proposed by Steven Weinberg in 1967 /19/ may be renormalizable. /20/ H.H. Chen /21/ and H.H. Chen with B.W. Lee /22/ have used the data of Reines and Gurr to restrict the parameter $x = \sin^2\theta$ which appears in the Weinberg theory. For the limit $g^2/g_p^2 < 4$, x is restricted to be ≤ 0.45.

Recently, the prospects for observing reaction /2/ have been emphasized at the Los Alamos Meson Physics Facility [LAMPF] /23a/ where an intense neutrino flux is about to become available. The neutrinos which originate in the LAMPF beam stop are predominantly from the decay of μ and π at rest. The muon-neutrinos from these decays are below the energy necessary for muon production. The observation of /2/ is therefore not overwhelmed by inverse beta reactions which are so prevalent at higher energy accelerators. However the existence of neutrinos from both μ and π decay does present other possibilities not available at the reactor or at higher energy accelerators, especially if the multiplicative lepton conservation law is valid and / or neutral leptonic currents [e.g., as in the Weinberg theory] exist /23b/.

The results to be considered today are the most recent in the continuing search for the process [eq. 1] at a fission reactor, /24/ using an approach which has been under development for this explicit purpose for the last dozen years.

EXPERIMENTAL CONFIGURATION

The geometry of this experiment is dictated in large part by the overwhelming gamma ray background. According to V-A theory, the 7.84 Kgm plastic scintillator target can be expected to experience less than 5 elastic scattering events per week above our data threshold of 3 MeV. The target region was therefore segmented into 16 optically isolated bars in an effort to label a gamma ray by its characteristic sequence of scatterings in the low z medium. We thus accepted as valid events only those in which the recoil electron was confined to a single bar [about 60 %].

The bars were viewed through NaI scintillators, and were surrounded by a NaI annulus. The result was a target entirely surrounded by at least 10.7 cm of NaI. Advantage was taken of the widely different pulse shapes from the plastic and NaI scintillators to determine whether the event was located in the plastic or in the NaI light pipe. The NaI was in turn completely enclosed in a lead jacket 6.3 cm thick and then wrapped in a thin cadmium sheet. This unit was entirely immersed in a cylindrical 2,2oo liter liquid anti-coincidence detector.

The cylindrical anti-coincidence tank was itself immediately below a large rectangular liquid detector, and surrounded by water tanks and another lead shield 20 cm thick. More water tanks outside this lead completed the shielding. Fig. 1 is a schematic of the detector unit inside the cylindrical tank.

The photomultiplier signals representing the target portion of the apparatus as well as the many live regions surrounding the target were displayed on an oscilloscope and photographed. Trigger constraints required a pulse of the shape appropriate to the plastic scintillator with an energy deposition $\geq$ 2.7 MeV and unaccompanied by a pulse in either the rectangular or cylindrical liquid tanks which would be characteristic of a cosmic ray muon. The anti-coincidence gate was 4oo μ sec wide to exclude triggers from neutrons produced by an earlier muon. Further screening of events was accomplished by analyzing the film record. Such an inspection revealed the location of the event in the target, the energy of the event, and the presence, if any, of NaI or liquid scintillator pulses. The NaI annulus was capable of detecting an energy deposition of $\sim$ 40 KeV, while the NaI light pipes could reveal the presence of as little as $\sim$ 200 KeV energy deposition. Signal lights also displayed on each frame were used to denote the occurrence of various characteristic events.

ENERGY CALIBRATION AND STABILITY

The rapidly falling spectrum characteristic of the background makes it essential to maintain a continuous check of the system gain. Such a check is provided by the beta decay [3.2. MeV end point] of Bi^{214} contaminant. The procedure used was to measure the rate of Bi^{214} decay as identified by a delayed daughter α particle [7.7. MeV] $\tau_{1/2}$ = 164 μ se The energy associated with a preselected Bi^{214} rate was taken as the reference value and all runs normalized to it. The absolute energy scale was set to within ± 2 % by means of a Th^{206} gamma source [2.62 MeV]. Both sources were viewed by the entire plastic in anticoincidence with the surrounding NaI. The correction to the overall gain of the system as measured by the Bi^{214} was within ± 2 % . Fig. 2 shows these calibration spectra.

Reactor Associated Backgrounds

A. $\bar{\nu}_e + p \rightarrow n + e^+$

Inverse beta decay produces in our detector a positron which will masquerade as an elastic scattering event if we fail to detect both of the accompanying positron annihilation gammas and the delayed neutron capture gammas. Fortunately, we can estimate this source of background by measurements of the annihilation gamma detection efficiency for positrons in the plastic detector and the neutron detection efficiency for neutrons from inverse beta decays in the plastic detector.

If we denote the probability of observing a single annihilation gamma as η_γ and the probability of seeing the neutron, hydrogen capture gamma as η_n then, the probability of an inverse beta decay event masquerading as an elastic scattering event is given by $P = [1 - \eta_n] [1 - \eta_\gamma]^2$. From observation of positron annihilation in our detector, for those events exhibiting neutron capture, we conclude that $\eta_\gamma = 0.85 \begin{smallmatrix} +0.02 \\ -0.03 \end{smallmatrix}$. The neutron detection efficiency is estimated to be ~ 0.2 by comparing the rate of observed inverse beta decays with the predicted rate. Hence

$$P \approx [1 - 0.2] [1 - 0.85]^2 = 0.018$$

. The number of $\bar{\nu}_e + p$ reactions per day in our plastic target capable of depositing between 3.6 and 5.0 MeV in a single plastic element is calculated to be 6.2/day. This calculation took into account the fraction of single element events which we measured with our detector and allowed for the 1 MeV annihilation energy which added to the positron recoil energy in these cases. The number mistaken as elastic scattering is, therefore,

$$\left[0.018 \begin{smallmatrix} +.008 \\ -.005 \end{smallmatrix}\right] \times [6.2/\text{day}] = 0.11 \begin{smallmatrix} +.05 \\ -.03 \end{smallmatrix} \; / \text{d}$$

It is amusing to observe that in searching for the elastic scattering $\bar{\nu}_e$ interaction we must contend with this relatively large neutrino produced background.

B. Neutrons

It can be imagined that neutrons produce background events by direct recoil from protons in the plastic. However, this process is discriminated against by the poor scintillation efficiency of the plastic for recoil protons in the relevant energy range. [It is $\sim$ 0.5 that for electrons at $\sim$ 6 MeV.] Neutrons of the requisite energy, $>$ 6 MeV cannot arise from natural photo processes.

A test of maximum possible neutron associated background was made using a Pu-Be source. It produced a count rate in the NaI detector of 92.9/min, [4 $\rightarrow$ 10 MeV] and a single element rate of 0.002/min in the plastic detector [3.6 $\rightarrow$ 5.0 MeV]. The reactor associated background in the NaI was measured to be $\lesssim$ 0.15/min. Assuming the neutron spectrum from the reactor to be as energetic as that associated with the source <u>and</u> further assuming that all the annulus up-down difference is due to neutrons, we find that the maximum reactor associated background in the plastic arising from neutrons is

$$\frac{0.15}{92.9} \times [0.002/\text{min}] \quad [1440 \text{ min/day}] = 0.005/\text{day}, \quad [3.6 \rightarrow 5.0 \text{ MeV}]$$

C. Gammas

Although reactor associated gammas probably arise in large measure from neutron capture, we make a separate limit for completeness. Deductions as to the maximum background contribution from these gammas can be made by comparing reactor associated rates in the NaI anticoincidence detector before the addition of 6.3 cm Pb outside the NaI [12/min] and scaling the pre-Pb plastic rate [6/day]

$$\frac{0.15}{12} [6/\text{day}] = 0.075/\text{day}$$

Since 1/3 of these will be single element events, we obtain a maximum background due to reactor associated gammas of

$$0.025/\text{day} \quad [3.6 \rightarrow 5.0 \ \text{MeV}]$$

RESULTS

Fig. 3 shows the observed plastic detector events in the elastic scattering mode in the energy range $3.0 \rightarrow 5$ MeV for reactor "on" [95 days] and "off" [52 days]. The rates and run times have been corrected for dead time [$\sim$ 22 %]. Also shown are detector rates predicted by the V-A theory /13,14,25/. These predictions take into account the increase in rate due to detector energy resolution [12 % half width at half maximum], and a decrease in rate [$\sim$ 40 %] due to electrons escaping a single plastic element. Also included in the predicted rate is a statement of errors based on the uncertainties in the fission $\bar{\nu}_e$ spectrum and in the β decay coupling constants. /26/

The sum of the above reactor associated backgrounds [$\sim$0.14/day] is approximately the same as our observed reactor on-off difference [for E_{min} = 3.6 MeV] and the data suffer from a large statistical uncertainty. We chose the results for E_{min} = 3.6 MeV as a compromise between minimizing $\frac{\delta}{\text{V-A}}$ and maximizing $\frac{\text{V-A}}{R_{off}}$

Based on the data in Table I, we find a one standard deviation upper limit on this elastic scattering signal to be

$$< .15 \pm .17 - .11 \begin{smallmatrix} \pm .05 \\ .03 \end{smallmatrix} \begin{smallmatrix} < \\ < \end{smallmatrix} \begin{smallmatrix} 0.4 + \sqrt{[.17]^2 + [.05]^2} \\ 0.22/\text{day} \end{smallmatrix}$$

In obtaining this limit, we have conservatively not subtracted off our small reactor associated neutron and gamma background. This corresponds to an elastic scattering cross-section limit σ_{expt} $< 6.6 \times 10^{-47}$ cm^2 fission $\bar{\nu}_e$ for the production of recoil electrons in the energy range $3.6 \rightarrow 5.0$ MeV.

It is apparent from Fig. 3 that we are searching for a small elastic scattering signal superimposed on a very rapidly falling background. Thus gain shifts which may be $\leq \pm 1\ \%$ could produce counting rate shifts $< {}^{+.14}_{-.10}$ day^{-1}. Combining the statistical and gain uncertainties gives an uncertainty in the measured rate of $\pm .23$. In this way we arrive at an upper limit, relative to the V-A prediction, of

$$\frac{\sigma_{expt}}{\sigma_{V-A}} < \frac{[0.04 + 0.23]}{0.16}$$
$$< 1.7$$

If we, instead of taking the central value and the maximum deviation we allowed for the asymmetry in the $\pm$, the limit would be the same i.e.

$$< \frac{.05 + .21}{.16} < 1.7$$

The most recent model presented by Weinberg [20] takes the parameter X = 0.25 $\theta = 30^\circ$. This would imply a rate in our detector in the energy range $3.6 \rightarrow 5.0$ MeV of about 1.3 times [22] the rate predicted by V-A theory. Our limits in this case are therefore

$$\frac{\sigma_{expt}}{\sigma_{\substack{Weinberg \\ \theta = 30^\circ}}} < 1.3$$

In terms of a magnetic moment this implies $< 2 \times 10^{-10}$ Bm.

Fig.4 indicates the limits which these data place on various theoretical possibilities. /19,21,22/ We note that the cross-section in the case for $x = 0$ [$\Theta = 0^{\circ}$] is $\frac{\sigma_{V-A}}{4}$ This is at the limit of accessibility for the improved experiment which we now describe.

Improved Experiment

Consideration of the data in Table I indicates that the system used in these measurements was seriously limited by the small target mass and a high background which placed a premium on instrumental stability. An estimate shows that the primary source of background was from unlabeled Bi^{214} decay. Accordingly, we have rebuilt the detector, doubling the size of the plastic target and thinning the aluminized mylar optical dividers to 5.3×10^{-4} g/cm^{2} so as to make visible the otherwise absorbed alpha particle which follows the Bi^{214} decay. In addition, we have lowered the unlabled inverse beta background by removing the MgO packing material, which separated the NaI and plastic scintillators. Preliminary runs with the modified system indicate a marked reduction in the background from unlabeled B^{214} . The overall increase in the predicted V-A signal to background appears to be greater than a factor of 3 [Fig. 5].

Although the elastic scattering data in this improved system are still too sketchy to be discussed here, we can report that the inverse beta decay process has already been observed in the modified detector. This process, seen at a rate of about 15 day^{-1} will enable us to make a direct determination of the fission antineutrino spectrum. As mentioned earlier, it will in addition, provide valuable evidence that the detector is functioning properly; the reaction can also be used to experimentally determine our efficiency for observing the annihilation gamma rays and the eventual neutron capture. Whereas the neutron detection efficiency remains at about 20 %, the .51 MeV annihilation gamma detection efficiency has been reised to $.96 \pm .04$. The result is that the inverse beta background between 3.6 and 5.0 MeV which was previously $.11/.16 \approx 70$ of the "expected" signal, is now < 15 % of the V-A predicted rate in the same energy range.

It should also be noted, that the ± 1 % gain uncertainty we chose increased our upper limit from $\frac{.22}{.16}$ to $\frac{.27}{.16}$ or by .3 /V-A/. That same 1 % gain shift with our improved signal to background corresponds to 0.1/V-A/ at the same energy. Accordingly, the improved detector should yield spectral information if the signal is at the V-A level.

For all of the above reasons, we feel confident that our modified detector will allow us to improve significantly the results we report to you here today, and will make possible an experimental statement as to the validity of V-A for the elastic scattering reaction and in general provide some further constraints on theoretical conjecture.

Acknowledgements

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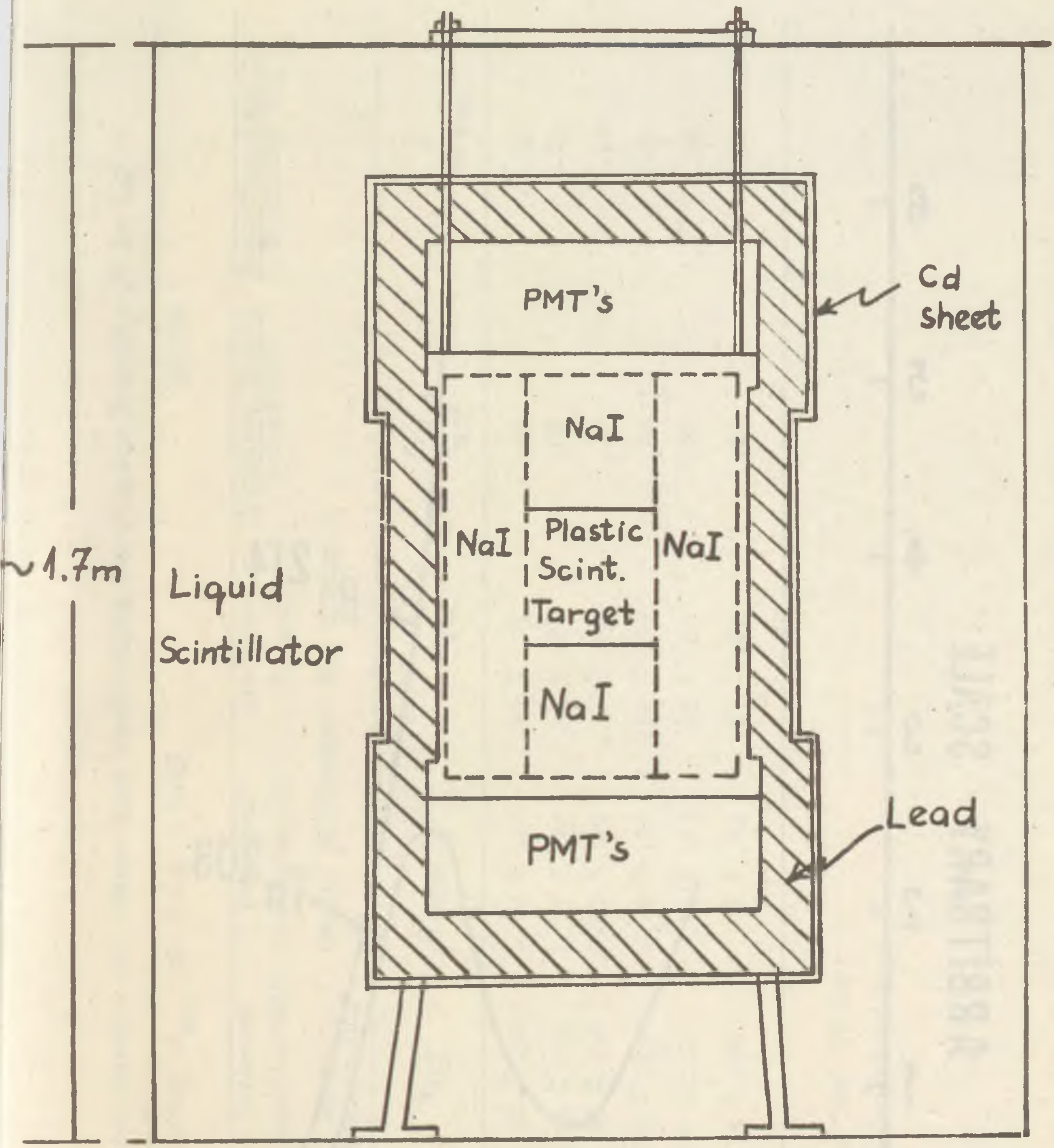

FIG. 1 SCHEMATIC OF ELASTIC SCATTERING DETECTOR CONFIGURATION.

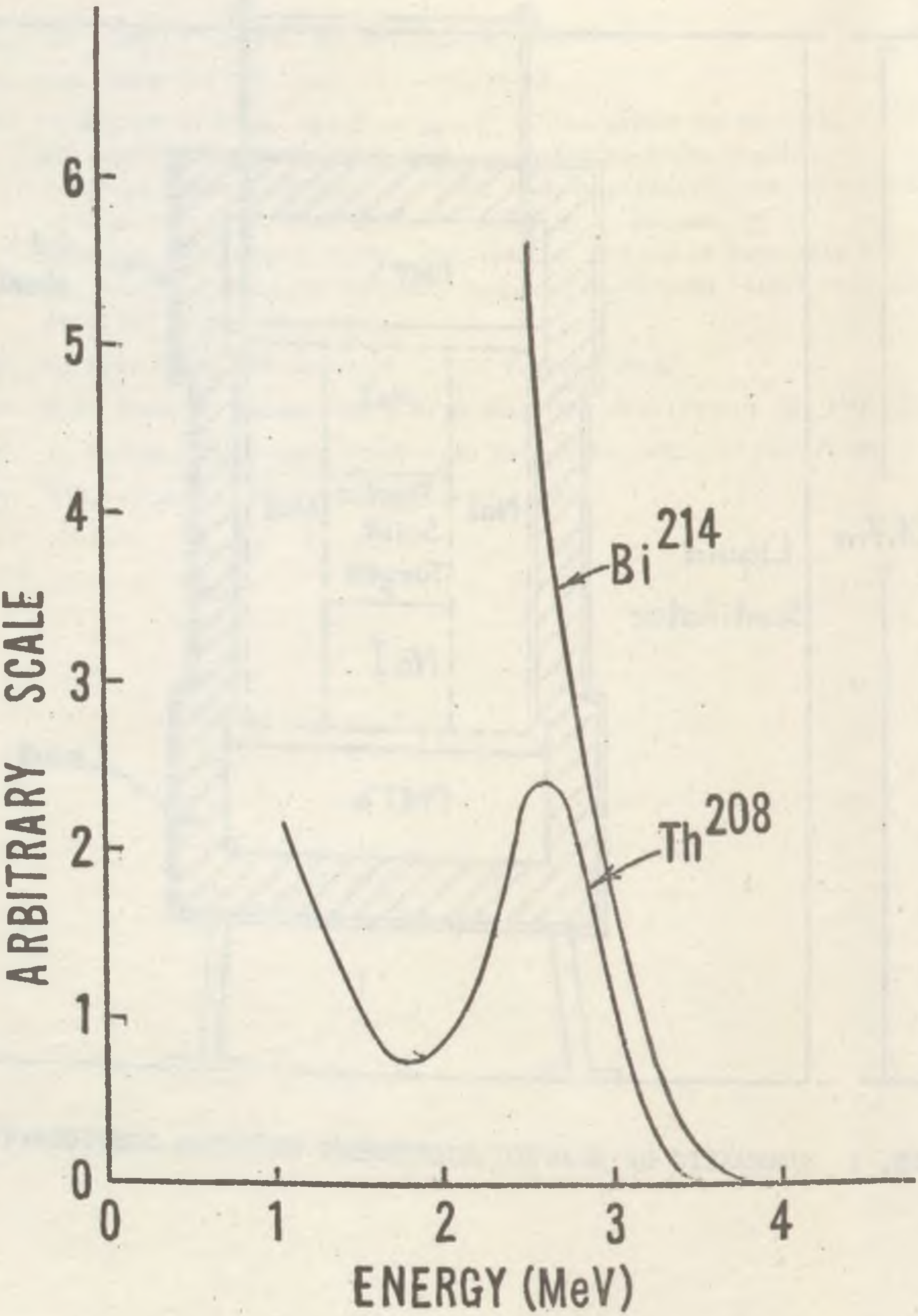

FIG. 2 SPECTRA OF CALIBRATION SOURCES.

Experimental Results and Theoretical Predictions for Plastic Detector [7.84 Kg]

	RATES [counts/day, $E_{min} \leq E \leq 5$ MeV]				RATIOS	
	Observed Reactor "On" Rate	Observed Reactor "Off" Rate [Background]	Observed Gross Reactor On-Off Difference	Predicted Signal V-A Theory	$\frac{\text{Predicted Signal}}{\text{Background}}$	$\frac{\text{Expt.Uncertainity}}{\text{Predicted Sig}}$
E_{min} MeV	R_{on}	R_{off}	$\Delta_{expt} \pm \delta$	Δ_{V-A}	$\frac{\Delta_{V-A}}{R_{off}}$	$\frac{\delta}{\Delta_{V-A}}$
3.0	6.43 ± .26	6.49 ± .35	−.06 ± .44	.40 ± .03	.06	1.10
3.2	3.28 ± .18	3.49 ± .35	−.21 ± .31	.29 ± .02	.08	1.06
3.4	1.82 ± .18	1.81 ± .18	.01 ± .22	.21 ± .02	.13	1.05
3.6	1.05 ± .11	.90 ± .13	.15 ± .17	.16 ± .01	.18	1.06
3.8	.68 ± .08	.54 ± .10	.14 ± .13	.12 ± .01	.22	1.08
4.0	.47 ± .07	.38 ± .09	.09 ± .11	.07 ± .01	.18	1.57

Run Time: Reactor "On" 95 days

Reactor "Off" 52 days

Fig. 3

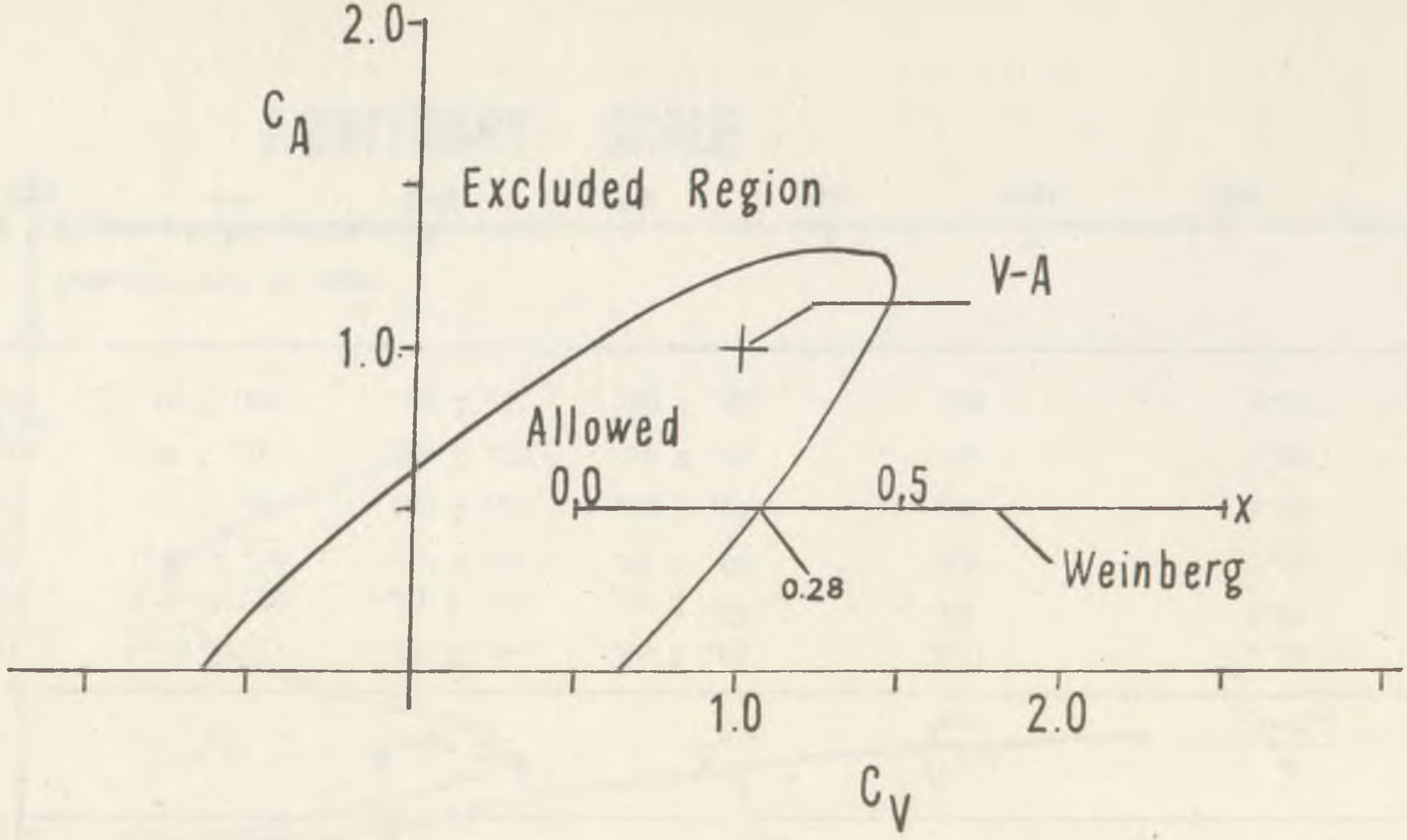

FIG. 4 EXPERIMENTAL LIMITS FOR THE VECTOR, C_V, AND AXIAL VECTOR, C_A, COUPLING CONSTANTS APPROPRIATE TO THE ELASTIC SCATTERING PROCESS /EQ.1/.

If C_V, C_A are the only couplings involved

$$\frac{1}{m_e}\frac{d\sigma}{dT} = \frac{G_w^2}{2\pi}\left\{(c_V-c_A)^2+(c_V+c_A)^2\left(1-\frac{T}{\omega}\right)^2-(c_V^2-c_A^2)\frac{m_e T}{\omega^2}\right\} \text{(ref. 22)}$$

$$\left.\begin{array}{l} C_V = 1/2 + 2x \\ C_A = 1/2 \end{array}\right\} \text{Weinberg conditions relating } C_V \text{ and } C_A$$

$$G_w = 1.026 \times 10^{-5}\ m_p^{-2}$$

E_{min} [MeV]	Observed Rate [day^{-1}] [Signal plus Background] in Energy Range $E_{min} \leq E \leq 5$ MeV		Predicted Rate [day^{-1}] V-A Theory		$\dfrac{\text{Total rate}/\text{Prediction}_{new}}{\text{Total rate}/\text{Prediction}_{old}}$
	[New]	[Old]	[New]	[Old]	
3.0	7.0 ± .7	6.5 ± .4	.80	.40	1.8
3.2	2.4 ± .4	3.5 ± .4	.58	.29	3.0
3.4	.99 ± .25	1.8 ± .2	.42	.21	3.5
3.6	.56 ± .19	1.1 ± .1	.32	.16	3.8
3.8	.37 ± .15	.68 ± .08	.24	.12	3.6
4.0	.25 ± .12	.47 ± .07	.14	.07	3.7

"New" Reactor "On" 15 days

"Old" " " 95 days

Fig. 5

Signal plus Background and Theoretical Predictions for "Old" [7.84 Kgm plastic detector] and Improved detector [15.7 Kgm, plastic]

THE LEPTON ERA OF THE BIG BANG

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1. The Hot Big Bang model of the Universe

In 1929 the observation of red shifts in spectra of remote galaxies indicated that the Universe is expanding (Hubble, 1929) and that it can be described by the models of Friedmann (1922,1924). These models were given as the following non-static solutions of the Einstein field equations, applied to the Robertson-Walker metric for a homogeneous and isotropic distribution of matter in the Universe:

(1) $$\left(\frac{\dot{R}}{R}\right)^2 + \frac{kc^2}{R^2} = \frac{8\pi G}{3}\rho$$

(2) $$2\frac{\ddot{R}}{R} + \left(\frac{\dot{R}}{R}\right)^2 + \frac{kc^2}{R^2} = -\frac{8\pi G}{c^2}p$$

In these formulae R(t) is the scale factor which characterizes the "size" of the Universe; $\dot{R}$ and $\ddot{R}$ are its first and second order time derivatives. The curvature index k determines whether the model is closed and spherical (k = +1), open and euclidean (k = 0) or open and pseudo-spherical (k = -1). The density ρ and pressure p are also functions of time, for which the behaviour depends on the equation of state. In two extreme cases this equation is given by

(3) $p = 0$ for a matter-dominated model and

(4) $p = \frac{1}{3}\rho c^2$ for a radiation-dominated model.

In the Big Bang model it is assumed that the evolution of the Universe starts with a point-singularity at expansion time $t = 0$. Gamow (1946,1949) suggested further that in the initial stages the temperature was extremely high. This Hot Big Bang model describes a Universe which is filled with a mixture of photons and relativistic particles for which $0 \leq mc^2 < kT$ and where equation (4) holds. Non-relativistic particles form a very small background; therefore we can use the radiation-dominated model. The conservation of entropy for a specific comoving volume during the expansion (a consequence of equations (1) and (2)) is given by

$$(5) \qquad \frac{d}{dt}(\rho R^3) + \frac{p}{c^2} \cdot \frac{dR^3}{dt} = 0$$

This becomes with $p = \frac{1}{3}\rho c^2$:

$$(6) \qquad \frac{d}{dt}(\rho R^4) = 0.$$

Thus ρR^4 is a constant during the expansion, which we can call Q. From equation (1) we get after substitution of $\rho R^4 = Q$:

$$(7) \qquad \left(\frac{dR}{dt}\right)^2 = \frac{8\pi GQ}{3} \cdot \frac{1}{R^2} - kc^2.$$

For small values of R (in the initial stages of expansion) we can neglect the term $-kc^2$ and find after integration the following solution for the expansion time:

$$(8) \qquad t = \left(\frac{3}{32\pi GQ}\right)^{\frac{1}{2}} R^2$$

This relation is valid for all models in early stages whereas it holds for the Einstein- De Sitter model (with k = 0) during the whole expansion. For the density we find from equation (8):

$$(9) \qquad \rho = \frac{Q}{R^4} = \frac{3}{32\pi G t^2} \simeq 4.5\text{x}10^5 \, t^{-2} \text{ g-cm}^{-3}$$

From equation (9) we see that for expansion time $t \approx 10^{-5}$ sec the density was $4.5\text{x}10^{15}$ g-cm^{-3} (which is larger than the density of nuclear matter).

If we assume that the known laws of physics can be extrapolated to these extremely high values of the density, we can relate the processes which played a role in the Hot Big Bang to the properties of the elementary particles in the cosmic plasma.

2. Physical processes during the Hot Big Bang

With the initial high temperatures and densities in the Universe the interaction processes between all kinds of elementary particles are sufficiently frequent in order to establish thermodynamic equilibrium. This implies that particles of a certain kind are described by a specific thermal distribution function with for all particles the same value of the temperature T. For fermions (baryons, leptons) one has the Fermi-Dirac distribution (+) and for bosons (π-mesons, photons) the Bose-Einstein distribution (-) with particle density (Landau and Lifschitz, 1958; Chiu, 1968)

$$(10) \qquad n = \frac{g}{2\pi^2\hbar^3} \int_0^\infty \frac{p^2 dp}{\exp(E-\mu)/kT \pm 1}$$

and energy density

$$(11) \qquad \varepsilon = \frac{g}{2\pi^2\hbar^3} \int_0^\infty \frac{Ep^2 dp}{\exp(E-\mu)/kT \pm 1}$$

In these formulae E is the relativistic energy of the particle $E = (p^2c^2 + m^2c^4)^{\frac{1}{2}}$, k the Boltzmann constant, g the spin statistical weight and μ the chemical potential (which is related to the degeneracy parameter $\xi = \mu/kT$).

As examples we give the following cases:

a. Photons: bosons with $g = 2$, $E = pc$, $\mu = 0$

$$(12) \qquad n_\gamma = \frac{1}{\pi^2 \hbar^3 c^3} \int_0^\infty \frac{E^2 dE}{\exp E/kT - 1} = \frac{1}{\pi^2} \left(\frac{kT}{\hbar c}\right)^3 \times 2\zeta(3) \simeq 20T^3$$

$$(13) \qquad \varepsilon_\gamma = \frac{1}{\pi^2 \hbar^3 c^3} \int_0^\infty \frac{E^3 dE}{\exp E/kT - 1} = \frac{kT}{\pi^2} \left(\frac{kT}{\hbar c}\right)^3 \times \frac{\pi^4}{15} = aT^4$$

with radiation density constant

$$(14) \qquad a = 7.6\times10^{-15} \text{ erg-cm}^{-3}\text{-degrees}^{-4}$$

b.1. Non-degenerate neutrinos: fermions with $g = 1$, $E = pc$, $\mu = 0$ ($n_\nu = n_{\bar\nu}$)

$$(15) \qquad n_\nu^{ND} = \frac{1}{2\pi^2 \hbar^3 c^3} \int_0^\infty \frac{E^2 dE}{\exp E/kT + 1} = \frac{1}{2\pi^2} \left(\frac{kT}{\hbar c}\right)^3 \times \frac{3}{2}\zeta(3) \simeq 7.5T^3$$

$$(16) \qquad \varepsilon_\nu^{ND} = \frac{1}{2\pi^2 \hbar^3 c^3} \int_0^\infty \frac{E^3 dE}{\exp E/kT + 1} = \frac{kT}{2\pi^2} \left(\frac{kT}{\hbar c}\right)^3 \times \frac{7\pi^4}{120} = \frac{7}{16}\, aT^4$$

These relations hold for all neutrinos ν_e, $\bar\nu_e$, ν_μ and $\bar\nu_\mu$.

b.2. Degenerate neutrinos: in this case we have $\mu \neq 0$, and get densities $n_\nu \neq n_{\bar\nu}$ which depend on the degeneracy parameter $\xi = \mu/kT$:

$$(17) \qquad n_\nu^{D} = \frac{1}{2\pi^2 \hbar^3 c^3} \int_0^\infty \frac{E^2 dE}{\exp[(E-\mu)/kT] + 1} = \frac{1}{2\pi^2} \left(\frac{kT}{\hbar c}\right)^3 \mathcal{F}_2(\xi)$$

$$(18) \qquad \varepsilon_\nu^{D} = \frac{1}{2\pi^2 \hbar^3 c^3} \int_0^\infty \frac{E^3 dE}{\exp[(E-\mu)/kT] + 1} = \frac{kT}{2\pi^2} \left(\frac{kT}{\hbar c}\right)^3 \mathcal{F}_3(\xi)$$

where the Fermi functions $\mathcal{F}_n(\xi)$ are defined as (Chiu, 1968)

$$(19) \qquad \mathcal{F}_n(\xi) = \int_0^\infty \frac{x^n dx}{\exp(x-\xi) + 1}$$

For antineutrinos we get similar expressions with $\mathcal{F}_n(\xi)$ replaced by $\mathcal{F}_n(-\xi)$. For non-degeneracy $\xi = 0$ we get back relations (15) and (16), and for extreme degeneracy with $\xi \to \infty$, we get $\mathcal{F}_2(\xi) \to \frac{1}{3}\xi^3$ and $\mathcal{F}_3(\xi) \to \frac{1}{4}\xi^4$ and thus:

(20) $$n_\nu^{ED} = \frac{1}{2\pi^2}\left(\frac{kT}{\hbar c}\right)^3 \times \frac{1}{3}\xi^3 = \frac{1}{6\pi^2(\hbar c)^3}\mu^3$$

(21) $$\varepsilon_\nu^{ED} = \frac{kT}{2\pi^2}\left(\frac{kT}{\hbar c}\right)^3 \times \frac{1}{4}\xi^4 = \frac{1}{8\pi^2(\hbar c)^3}\mu^4$$

c.1. <u>Relativistic non-degenerate electrons</u>: fermions with $g = 2$, $\mu = 0$, $kT \gg mc^2$:

(22) $$n_e^R = \frac{1}{\pi^2\hbar^3c^3}\int_0^\infty \frac{E^2dE}{\exp E/kT + 1} \simeq 15T^3$$

(23) $$\varepsilon_e^R = \frac{1}{\pi^2\hbar^3c^3}\int_0^\infty \frac{E^3dE}{\exp E/kT + 1} = \frac{7}{8}aT^4$$

Because of the different value for g these quantities are 2x as large as for the neutrino gas.

c.2. <u>Non-relativistic non-degenerate electrons</u>: here we have $kT \ll mc^2$ and get for the particle density

(24) $$n_e = \frac{1}{\pi^2\hbar^3}\int_0^\infty \exp\left(-\frac{mc^2 + p^2/2m}{kT}\right)p^2dp = \frac{2}{h^3}(2\pi mkT)^{\frac{3}{2}}\exp(-mc^2/kT)$$

or numerically:

(25) $$n_e = 1.5\times10^{29}\, T_9^{3/2}\exp(-5.9/T_9)$$

where T_9 is the temperature in units 10^9K.

Relation (24) shows that the particle density of electrons decreases in a very fast way if the temperature becomes lower than $T = T_e^{ann} = m_ec^2/k \simeq 5.9\times10^9$K. During the expansion the temperature decreases and for values around this annihilation temperature T_e^{ann} the electrons and positrons disappear from the cosmic plasma by pair annihilation processes. The same holds in general for all elementary particles with finite rest mass m, like muons, mesons, and baryons: for $T > mc^2/k$ their density is comparable to the photon density, during further expansion they annihilate at temperatures $T < mc^2/k$.

The total energy density can be given as the sum of contributions from all available particles; in general one has the expression

$$(26) \qquad \varepsilon = \kappa a T^4$$

where κ is determined by the separate values for particles (e.g. 1 for photons, 7/16 for neutrinos, 7/8 for electrons) and decreases at lower temperatures. If there are only photons, neutrinos ν_e, $\bar{\nu}_e$, ν_μ, $\bar{\nu}_\mu$ and electrons e^-, e^+ one has, for instance

$$(27) \qquad \kappa = 1 + 4\times\frac{7}{16} + 2\times\frac{7}{8} = \frac{9}{2}$$

By comparing equation (26) with equation (9) we find the following relation between temperature and expansion time(Zel'dovich and Novikov, 1967):

$$(28) \qquad T = \left(\frac{3c^2}{32\pi G a}\right)^{\frac{1}{4}} \kappa^{-\frac{1}{4}} t^{-\frac{1}{2}} \approx 1.5\times10^{10} \kappa^{-\frac{1}{4}} t^{-\frac{1}{2}} \text{ K}$$

or

$$(29) \qquad t = 2.3\times10^2 \kappa^{-\frac{1}{2}} T_9^{-2} \text{ sec}$$

Particles with zero mass remain relativistic during the expansion of the Universe and form a thermal background which does not annihilate. The only component of this background which has been detected could be the background microwave radiation which may be interpreted as a cosmic black body radiation of photons resulting from the Hot Big Bang (Penzias and Wilson, 1965; Dautcourt and Wallis, 1968). For its present temperature the value T ≈ 2.7 K has been established. This corresponds to an energy density

$$(30) \qquad \varepsilon_\gamma = aT_\gamma^4 \approx 0.25 \text{ eV-cm}^{-3}$$

We indicate the present temperature of the photon gas by T_γ, which can be different from the present temperature of the neutrino gas, T_ν, as we shall see in Section 4.

Another important physical process in the early Universe is the decoupling of particles during certain stages of cosmic evolution. This decoupling can be defined for those particles whose interaction with the cosmic plasma becomes too weak for the further establishment of thermal equilibrium with the rest of the cosmic plasma. This occurs when the characteristic collision time for these particles becomes longer than the expansion time of the Universe, which is given in equation (29). For this characteristic collision time one has the relation

$$(31) \qquad \tau = \left[< \sigma v > n \right]^{-1}$$

where n is the density of target particles and $< \sigma v >$ is the averaged product of cross-section and velocity. Both quantities decrease steeply for lower temperatures (e.g. relativistic particle densities proportional to T^3) which implies that τ increases much faster than the expansion time t_{exp}. At a certain stage of expansion both times can become equal and the corresponding temperature is defined as the decoupling temperature. Obviously this decoupling happens much earlier for neutrinos than for photons, as the weak interaction processes have much smaller cross-sections than the electromagnetic processes: the electron-neutrinos decouple for $T \approx 10^{10}$ K, the photons for $T \approx 3000$ K.

3. Stages in the Early Universe

During the evolution of the Early Universe one can distinguish the following stages in the Hot Big Bang in relation to the physical properties of specific elementary particles.

a. The Hadron Era

In this stage the temperature is so high that hadrons can be relativistic. The end of this hadron era is determined by the temperature at which the hadrons with the lowest mass (pions with $M_\pi \approx 135$ MeV) annihilate. This happens for $T^{ann} = M_\pi c^2/k \approx 1.7\times10^{12}$ K, which corresponds to an expansion time of 4×10^{-5} sec and a density of 3×10^{14} g-cm^{-3}. The annihilation of baryon-antibaryon pairs takes place in such a way that we are left with a finite number of baryons. The fact that the baryon number is a conserved quantity can be explained by assuming that initially the particle density of baryons n_B was larger than the particle density of antibaryons $n_{\bar{B}}$. Therefore the Hot Big Bang can be characterized by the baryon number

$$(32) \qquad B = \frac{n_B - n_{\bar{B}}}{n_\gamma}$$

Like other parameters which can be introduced the value of this number is arbitrary and not determined by the model. In a "symmetric" model of the Universe (see e.g. Omnès, 1971) it is assumed that originally the Universe was symmetric between particles and antiparticles with B = 0 and that during the expansion a phase transition took place which separated matter and antimatter into different regions. Because of the theoretical difficulties with this separation mechanism and the lack of evidence for the present existence of antimatter in the Universe (Steigman, 1972) such a model is rejected by many astrophysicists who prefer to introduce an arbitrary free parameter in the initial stages of the Universe. The value of this parameter is determined by the present matter density for which one can take (Oort, 1958)

(33) $$\rho_B = 3\times10^{-31} \text{ g-cm}^{-3}$$

which corresponds to a particle density $n_B = 6\times10^{23}\ \rho_B = 1.8\times10^{-7}\ \text{cm}^{-3}$ and to a baryon number

(34) $$B = n_B/n_\gamma \approx 5\times10^{-10}$$

This number could be much higher if the total density in the Universe would have a larger value. The density could, e.g., be equal to the critical density

(35) $$\rho_c = \frac{3H_o^2}{8\pi G} \approx 1.1\times10^{-29} \text{ g-cm}^{-3}$$

which depends on the Hubble constant $H_o = (\dot{R}/R)_{t=t_o}$ and determines whether the Universe is open ($\rho \leq \rho_c$) or closed ($\rho > \rho_c$). The relation between baryon particle density and photon temperature can be given by

(36) $$\rho_B = mn_B = mBn_\gamma = hT_9^3$$

where h is another parameter which can be used instead of B (Wagoner, Fowler and Hoyle, 1967). The relation between both quantities is

(37) $$B = 3.0\times10^{-5}\ h$$

The ratio of the energy densities in the form of baryons and photons which decrease, respectively, as T^3 and T^4, is for $B = 5\times10^{-10}$ given by

(38) $$\varepsilon_B/\varepsilon_\gamma = 2\times10^3\ T^{-1}$$

Both energy densities become equal for $T \approx 2000$ K, which can be taken as the final temperature of the "primeval fireball". In much earlier stages of the expansion the baryons contribute a very small fraction to the total energy density which is mainly determined by the background of photons and relativistic particles.

b. The Lepton Era

During this stage in the Early Universe the temperature decreases from 1.7×10^{12} K to the value $T = m_e c^2/k \approx 5.9 \times 10^9$ K where the electrons annihilate. Because of relations (9) and (29) this corresponds to expansion times between 4×10^{-5} sec and 3 sec, and to densities between 3×10^{14} g-cm^{-3} and 5×10^4 g-cm^{-3}.

The cosmic plasma consists of photons, neutrinos ν_e, $\bar{\nu}_e$, ν_μ, $\bar{\nu}_\mu$, electrons and muons; at a temperature of $\approx 1.2 \times 10^{12}$ K (corresponding to $kT = m_\mu c^2$) the muons annihilate and the value of κ decreases from 25/4 to 9/2. As we saw before, the baryon energy density is negligible, but there could be other components of the plasma which increase the value of the density and change the relation (29) and the behaviour of physical processes.

In the lepton era the neutrinos form a large part of the cosmic background which implies the important role of weak interaction physics in this stage of the Early Universe.

c. The Radiation Era

This stage begins after electron annihilation at $T \approx 5.9 \times 10^9$ K and it ends when the energy density of the photons becomes equal to the energy density of the baryons. We saw that this happens at a temperature of about 2000 K or higher, which corresponds to a density of $\approx 10^{-21}$ g-cm^{-3} and to an expansion time of $\approx 10^6$ years. After this period the Hot Big Bang comes to its end and the stellar era starts, in which galaxy and star formation

and all kinds of other phenomena take place.
An important event during the transition between the lepton and radiation era is cosmological nucleosynthesis: the creation of chemical elements in thermonuclear reactions which can occur at the given temperatures and densities. In fact Gamow (1946,1949) has proposed the Hot Big Bang theory for the early Universe in order to explain the observed abundance of helium. His mechanism of cosmological nucleosynthesis has been modified and extended by Hayashi (1950), Peebles (1966), Alpher et al. (1953,1967) and Wagonèr et al. (1967).

4. Neutrino decoupling during the lepton era

First we want to consider the processes between elementary particles in the lepton era and to compare the characteristic times for these processes with the expansion time (29). Strong interactions do not play an important role, for the hadrons are non-relativistic in this stage and as we saw their density is very small compared to the lepton and photon density. Electromagnetic interactions are responsible for scattering and annihilation processes like

$$(39) \qquad \gamma + e^{\pm} \rightarrow \gamma + e^{\pm} \quad ; \quad e^{+} + e^{-} \rightleftarrows \gamma + \gamma$$

which determine the thermodynamic equilibrium between photons and charged leptons. The cross-sections of these processes are high enough to maintain the equilibrium long after the lepton era.

Weak interaction processes are usually described by the Universal Fermi Interaction (UFI), which has been proposed by Feynman and Gell-Mann (1958). If this theory gives an appropriate description of the physical phenomena, the following processes are possible and their cross-sections can be calculated:

Processes with electrons and electron-neutrinos:

(40) $\nu_e + e^- \rightarrow \nu_e + e^-$; $e^+ + e^- \rightleftarrows \nu_e + \bar{\nu}_e$;

with muons and muon-neutrinos:

(41) $\nu_\mu + \mu^- \rightarrow \nu_\mu + \mu^-$; $\mu^+ + \mu^- \rightarrow \nu_\mu + \bar{\nu}_\mu$;

with both kinds of leptons:

(42) $\mu^- \rightarrow e^- + \bar{\nu}_e + \nu_\mu$; $\nu_e + \mu^- \rightleftarrows \nu_\mu + e^-$.

The processes (40) and (41) can be described by the "diagonal terms" in the Universal Fermi Interaction Hamiltonian, which can be written symbolically as

(43) $H_I = \frac{g}{\sqrt{2}} (\bar{p}n + \bar{\nu}_e e + \bar{\nu}_\mu \mu)(\bar{n}p + \bar{e}\nu_e + \bar{\mu}\nu_\mu)$

With the UFI theory one can calculate the cross-sections for the given neutrino scattering processes and find, e.g.

(44) $\sigma(\nu_e + e^-) = 3\sigma(\nu_e + e^+) = \frac{2g^2 m_e}{\pi\hbar^4} E_\nu^R$

in the electron rest frame with $E_\nu^R >> m_e c^2$. In the cosmic frame with non--degenerate neutrinos we can take for the average value of E_ν^R

(45) $<E_\nu^R> = <E_\nu> <E_e> /m_e c^2 = (3.2kT)^2/m_e c^2$

The characteristic time for scattering of neutrinos on the electron-neutrino gas is given by the relation (31), where we take for n the relativistic electron density (22). We find for this characteristic time

$$(46) \qquad \tau_{\nu e} \simeq 6.7\times10^{5}\ T_9^{-5}\ \text{sec}$$

Decoupling of the electron-neutrinos takes place when this time becomes equal to the expansion time (29). This happens when the temperature becomes equal to the decoupling temperature (De Graaf, 1970)

$$(47) \qquad T_{dec}^{\nu_e} \simeq 1.8\times10^{10}\ \text{K}$$

We see from its calculation that the decoupling temperature is related to the other quantities in the following way

$$(48) \qquad g^2 \langle E_\nu \rangle\ \kappa^{-\frac{1}{2}}\ T_{dec}^3 = A$$

where A is a numerical constant. The value (47) is larger than the electron annihilation temperature $T_{ann}^{e} = 5.9\times10^{9}$ K, which implies that the electron pairs started to annihilate in a stage where the electron-neutrinos were already decoupled. Therefore the temperature T_γ of the photon gas which is interacting with the annihilating electrons, is increased relative to the temperature T_ν of the electron-neutrino gas. From the conservation of entropy one finds (Alpher, Follin and Herman, 1953):

$$(49) \qquad T_\nu = \left(\frac{4}{11}\right)^{\frac{1}{3}} T_\gamma$$

which gives a present neutrino temperature of about 2 K.

For muon-neutrinos analogous considerations can be given (De Graaf, 1970). In this case the muons annihilate before muon-neutrino decoupling and one has to consider the non-relativistic expressions for the cross-section and the muon density. Therefore the decoupling temperature becomes higher:

$$(50) \qquad T_{dec}^{\nu_\mu} \simeq 1.2\times10^{11}\ \text{K}$$

Other values for the decoupling temperature and maybe for the present neutrino temperature would occur if the parameters in the characteristic time (g^2, $\langle E_\nu \rangle$) or in the expansion time (κ) would be different. The following possibilities can be distinguished:

a. Different theories of weak interaction

In a foregoing review (De Graaf, 1971a) we have considered the consequences of theories with neutral lepton currents, non-conservation of lepton number and other values for the coupling constant of diagonal interactions. The Weinberg theory (Reines, 1972) predicts a somewhat higher value for the cross-section of neutrino-electron scattering, whereas the present experimental upper limit is given by (Reines, 1972)

$$(51) \qquad \sigma_{exp}\,(\nu_e + e) < 1.7\,\sigma_{V-A}\,(\nu_e + e)$$

From relation (48) we see that the decoupling temperature depends on the value of the coupling constant according to $T_{dec} \propto g^{-2/3}$.

b. Degeneracy of the electron-neutrinos

Until now we have assumed that the electron-neutrinos are non-degenerate, which implies $n_{\nu_e} = n_{\bar{\nu}_e}$ and $\xi = 0$ according to relations (15) and (16). In a more general theory we can consider a new parameter in the Early Universe, the electron lepton number (Fowler, 1971)

$$(52) \qquad L_e = \frac{n_{\nu_e} - n_{\bar{\nu}_e} + n_{e^-} - n_{e^+}}{n_\gamma}$$

which at present reduces to

$$(53) \qquad L_e = \frac{n_{\nu_e} - n_{\bar{\nu}_e}}{n_\gamma} + \frac{n_p}{n_\gamma}$$

According to relations (12) and (17) we get for the first term

$$(54) \qquad \frac{n_{\nu_e} - n_{\bar{\nu}_e}}{n_\gamma} \simeq 0.21\,(T_\nu/T_\gamma)^3 \{\mathcal{F}_2(\xi) - \mathcal{F}_2(-\xi)\}$$

If we consider a neutrino degeneracy with degeneracy parameter $\xi > 0.1$, the tables for $\mathcal{F}_2(\xi)$ (Chiu, 1968) show that

$$(55) \qquad \frac{n_{\nu_e} - n_{\bar{\nu}_e}}{n_\gamma} > 0.02$$

This is much larger than values for the baryon number B like the one given in (34); therefore $n_p/n_\gamma < B$ can be neglected in (53) for a neutrino degeneracy with $\xi > 0.1$.

The average energy of the neutrinos in the degenerate Fermi gas is, according to relations (17) and (18) equal to

$$(56) \qquad \varepsilon_\nu/n_\nu = \{\mathcal{F}_3(\xi)/\mathcal{F}_2(\xi)\} \times kT$$

instead of 3.2 kT in the non-degenerate case. For large values of ξ we have $\mathcal{F}_3(\xi) \to \frac{1}{4}\xi^4$, $\mathcal{F}_2(\xi) \to \frac{1}{3}\xi^3$ and

$$(57) \qquad \varepsilon_\nu/n_\nu \simeq \frac{3}{4}\xi kT$$

The average neutrino energy is thus proportional to ξ, whereas the characteristic collision time becomes proportional to ξ^{-1}. On the other hand, the increase in density also changes relation (29) between expansion time and temperature. A large degeneracy $\xi >> 1$ would correspond to a neutrino energy density (cf. equation (21))

$$(58) \qquad \varepsilon_\nu = \frac{kT_\nu}{2\pi^2}\left(\frac{kT_\nu}{hc}\right)^3 \times \frac{1}{4}\xi^4 = \left(\frac{T_\nu}{T_\gamma}\right)^4 \times \frac{15\,\xi^4}{8\pi^4}\,\varepsilon_\gamma$$

Before electron annihilation the temperatures T_ν and T_γ were equal and we would have

(59) $\quad \varepsilon_\nu \approx 0.02\ \xi^4\ \varepsilon_\gamma$

At large values of ξ, e.g. $\xi > 10$, the value of κ in $\varepsilon_{tot} = \kappa\varepsilon_\gamma$ is mainly determined by the contribution (59). In relation to (48) we have $<E_\nu>$ proportional to ξ and κ proportional to ξ^4; the decoupling temperature increases thus as $\xi^{1/3}$ for higher values of the degeneracy parameter. In a previous review (De Graaf, 1971a) we compared the possible energy density of neutrinos with several limit values. The limit which can be established from nucleochronology (Zel'dovich and Novikov, 1967) is given by

(60) $\quad \varepsilon_\nu < \varepsilon_{nucl} \approx 4\ \varepsilon_{crit} \approx 2.5\text{x}10^4\ \text{eV-cm}^{-3}$

From relation (58) we find that ε_ν becomes larger than this limit value for $\xi > 100$. In the Hot Big Bang model extreme degeneracy with $\xi > 100$ would therefore not be compatible with condition (60).

c. Other particles in the cosmic background

One cannot exclude the possible existence of a large number of zero-mass particles which are difficult to observe and which might give an important contribution to ε_{tot} of equation (26). Possible particles are muon-neutrinos with a large degeneracy, which can increase the value of κ as in the case of degenerate electron-neutrinos. Further one can think of the possibility that there might exist other kinds of neutrinos ν_ℓ which differ from ν_e and ν_μ and are the partners of other charged leptons ℓ. Recently some attention has been given to the possible existence of such heavy leptons and to their consequences for elementary particle physics, astrophysics and cosmology (De Graaf, 1971b).

We see from equation (48) that an increase in κ due to an unknown background would increase the value of T_{dec} according to

$$(61) \qquad T_{dec} = \text{const} \times \kappa^{1/6}$$

The same holds for the equilibrium temperature which plays a role in the process of cosmological nucleosynthesis. This starts at the end of the lepton era and will be considered in the following Section.

Cosmological Nucleosynthesis and parameters of the Hot Big Bang

In the theory of Gamow (1946,1949) it was assumed that the original matter of the Universe consisted solely of neutrons. Through nuclear reactions such as β-decay and radiative capture element formation would be possible. This picture has been modified by Hayashi (1950), who observed that for high temperatures with $kT > m_e c^2$ (thus in the lepton era) the following weak interaction processes are fast enough to establish statistical equilibrium between neutrons and protons:

$$(62) \qquad n + e^+ \rightleftarrows p + \bar{\nu}_e \quad ; \quad p + e^- \rightleftarrows n + \nu_e$$

For these processes one can also define a characteristic time like the relation (46) for neutrino electron scattering. In this case the characteristic time becomes (Fowler, 1971)

$$(63) \qquad \tau_{char} \simeq 10^5 \, T_9^{-5} \text{ sec}$$

Neutrons and protons remain in equilibrium as long as this time is shorter than the expansion time $t_{exp} \simeq 10^2 \, T_9^{-2}$ of the Universe. As in the case of neutrino decoupling the equilibrium "freezes" at the equilibrium temperature for which $\tau_{char} = t_{exp}$. This happens for

(64) $$T_{eq} \approx 10^{10}\ K$$

In equilibrium the ratio between neutron and proton density n and p is determined by their mass difference Δmc^2:

(65) $$n/p = \exp(-\Delta mc^2/kT) \approx \exp(-15/T_9)$$

At very high temperatures $n/p \approx 1$; later it decreases for lower T_9. For temperatures below $T_9 \approx 10$ the ratio is fixed to

(66) $$(n/p)_{eq} \approx \exp(-15/10) \approx 0.22$$

Later on chains of strong interaction processes become possible at temperatures $T_9 < 2$ which produce isotopes like D, He^3, He^4, Li^7. A possible reaction channel is given by

(67) $$n\ (p,\gamma)\ D\ (p,\gamma)\ He^3\ (n,\gamma)\ He^4$$

In this way one can solve the so-called helium problem: the abundance of about 25% for helium in the Universe, which cannot easily be explained with stellar nucleosynthesis or with other processes (Searle and Sargent, 1972). The maximum helium abundance Y_{max} which can be obtained is found by combining all neutron pairs with proton pairs into helium nuclei:

(68) $$Y_{max} = 2\ n_{eq} = 2\ \frac{(n/p)_{eq}}{1 + (n/p)_{eq}} \approx 0.36$$

The real helium abundance will also depend on the free parameters of the Hot Big Bang, like the baryon number B and the lepton number L_e, which were introduced in (32) and (52).

For $L_e = 0$ (no electron-neutrino degeneracy) the resulting element abundances have been calculated by Wagoner et al. (1967); they find for instance that the value of the helium abundance Y varies as $0.21<Y<0.26$ for $5\times10^{-10}<B<5\times10^{-8}$. The results of these calculations are illustrated by Figure 1 (adapted from Fowler, 1971).

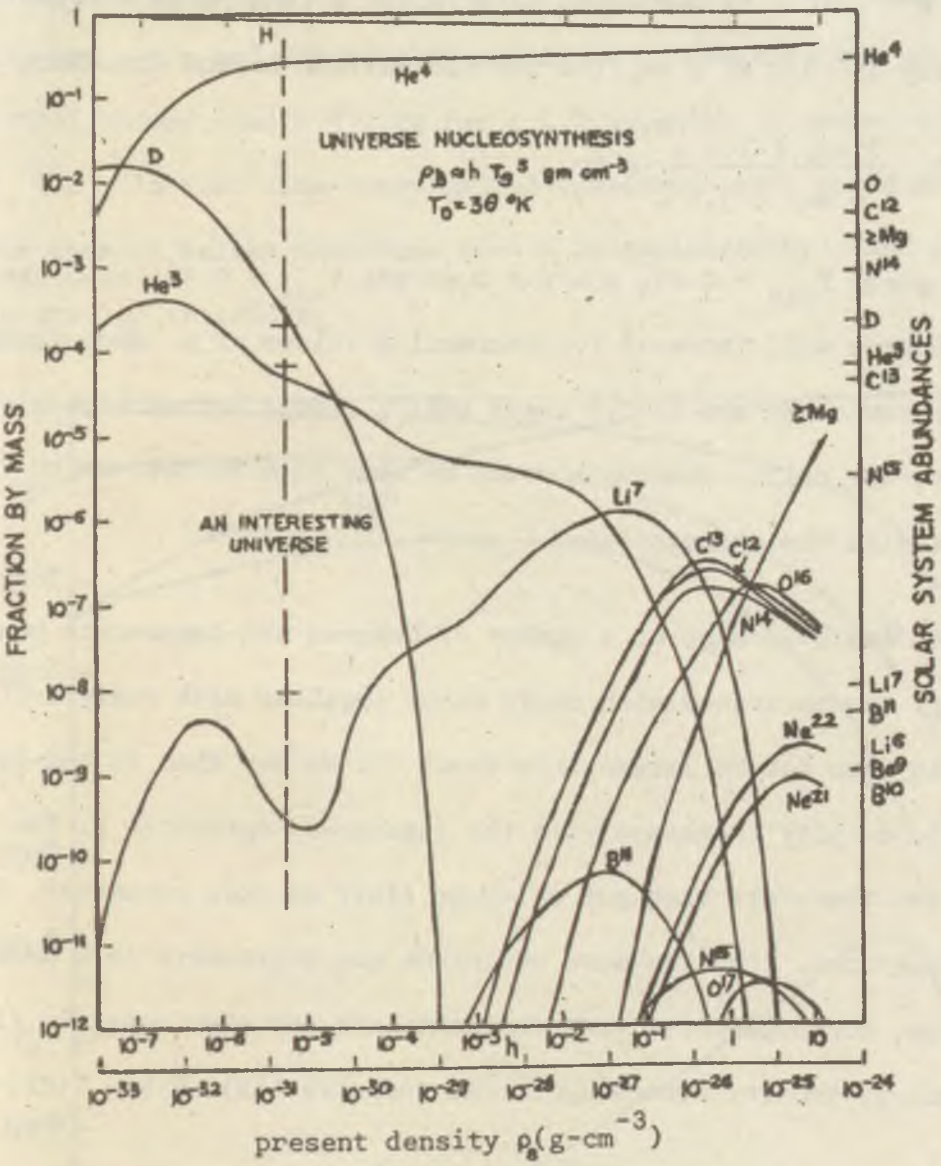

Fig. 1 Element production in the Hot Big Bang for different values of the present density without neutrino degeneracy (Fowler, 1971)

The calculations are done for the canonical Hot Big Bang with $L_e = 0$ in the usual background of particles and for $\varepsilon_{tot} = \kappa a T^4$ with $\kappa = 9/2$ in the lepton era. If the value of κ would be higher, e.g. because of the degeneracy of muon-neutrinos or the existence of other, unknown, zero mass particles, this would have the same effect on the equilibrium temperature as on the decoupling temperature in the foregoing Section. We find in a similar way that T_{eq} is proportional to $\kappa^{1/6}$. If κ is increased by a factor μ (the ratio between total energy density and $9/2\ aT^4$) we find for the maximum helium abundance

$$(69) \qquad Y_{max} = \frac{2 \exp(-1.5\ \mu^{-1/6})}{1 + \exp(-1.5\ \mu^{-1/6})}$$

The value $\mu = 2$ gives $Y_{max} \approx 0.42$, for $\mu = 3$ we get $Y_{max} \approx 0.45$; also the real helium abundance will increase for increasing values of μ. Shvartsman (1969) and Doroshkevich et al. (1971) argue that μ cannot become much higher than 3; otherwise the helium abundance would be more than 40-50%, which is in contradiction with the observations.

If the background would consist of a number of unknown non-degenerate bosons or fermions (e.g. new neutrinos wich could occur together with heavy leptons), this number could thus not be larger than about 10. We saw that in the case of degeneracy the density increases with the degeneracy parameter ξ. The given argument can therefore also put an upper limit on this parameter. If we assume, for instance, that the muon neutrinos are degenerate in a background of photons, non-degenerate electron-neutrinos and electrons, we find for the total energy density according to the formulae (13), (16), (18), and (23):

$$(70) \qquad \varepsilon_{tot} = \frac{15}{2\pi^4} \{ \mathcal{F}_3(\xi) + \mathcal{F}_3(-\xi) \}\ aT^4 + \frac{29}{8}\ aT^4$$

For $\xi = 4$ this becomes $\varepsilon_{tot} \approx 15\ aT^4$, which is more than three times larger than the canonical value 9/2 aT^4. The upper limit on the muon-neutrino degeneracy is thus given by $\xi \lesssim 4$.

Degeneracy of the electron-neutrinos does not only change the density and thus the equilibrium temperature T_{eq}, but also the way in which through reaction (62) the equilibrium between neutrons and protons is established. Neutrino degeneracy ($\xi > 0$) implies $n_\nu > n_{\bar{\nu}}$ and for a larger neutrino density the ratio n/p and the resulting helium abundance become smaller. Fowler (1971) showed that $Y < 0.02$ for $\xi \approx 2.5$, which occurs when $n_\nu \approx n_\gamma$ and also $\varepsilon_\nu \approx \varepsilon_\gamma$. In that case cosmological nucleosynthesis would be impossible and the present helium abundance should be explained by other processes, e.g. in stellar interiors.

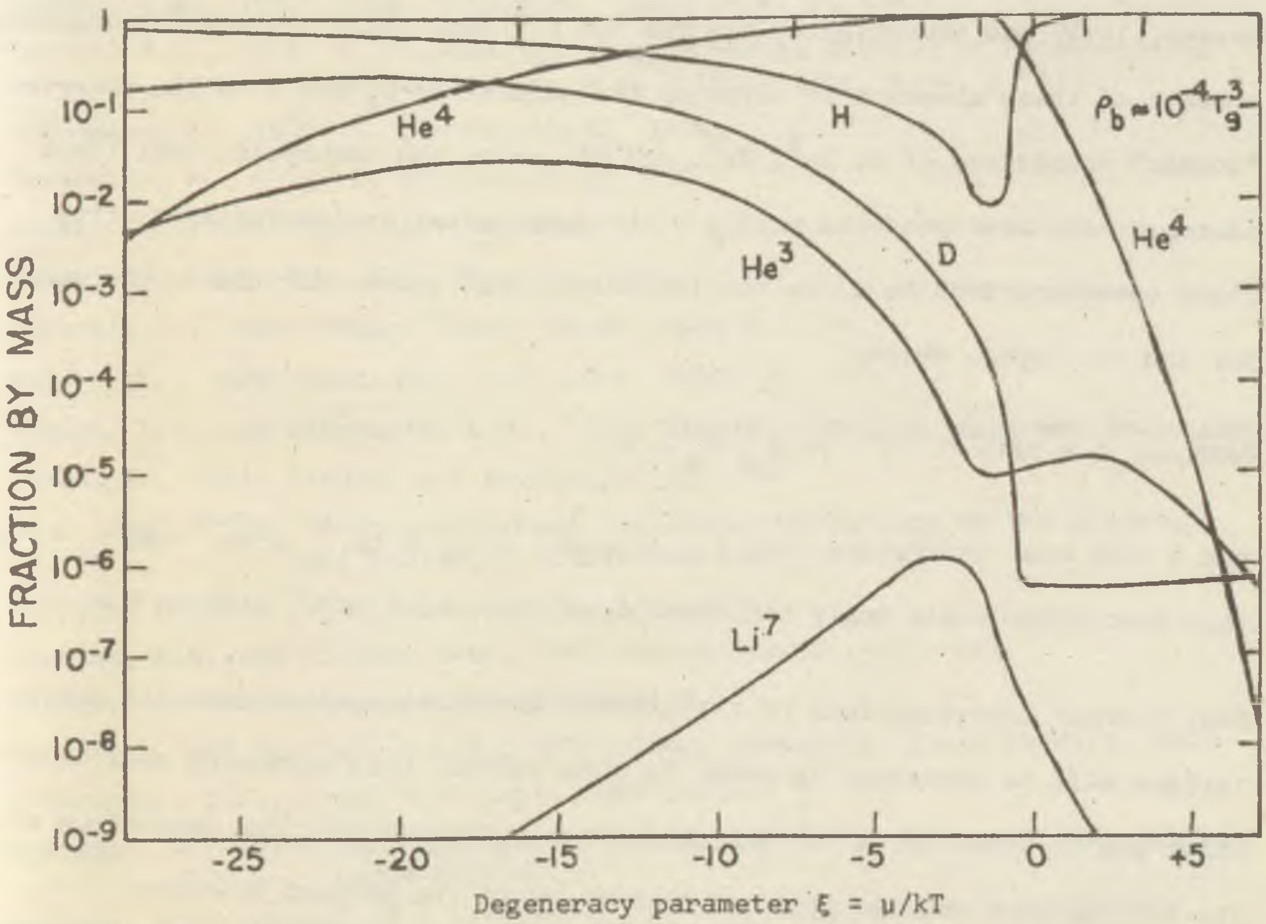

Fig. 2 Element production in a Universe with density $\rho \approx 2\times10^{-30}$ g-cm^{-3} for different values of the degeneracy parameters for electron-neutrinos (Fowler, 1971)

For $\xi < 0$ the opposite happens: $n_{\bar{\nu}} > n_{\nu}$ and n/p increases. At $\xi \approx -1.5$ this ratio becomes equal to one and the helium abundance reaches a maximum. Calculations of Wagoner et al. (1967) gave the results for the element abundances which are illustrated in Figure 2.

This figure gives the resulting abundances of light isotopes as a function of the degeneracy parameter $\xi = \mu/kT$ for a present matter density $\rho_B \approx 2\text{x}10^{-30}$ g-cm^{-3} (with $h \approx 10^{-4}$ and baryon number $B \approx 3\text{x}10^{-9}$). The same figure can also be given for other values of the density. In this way we find how the present abundance of the isotopes D, He^3, He^4, and Li^7 would depend on both parameters B and L_e, if these abundances originated from the Hot Big Bang with further the usual properties (like isotropy and homogeneity).

Reeves (1972) has investigated how the Hot Big Bang could produce the observed amounts of these elements by adapting the parameters ρ_B and ξ to the observed "cosmic" abundances of D, He^3, He^4, and Li^7 under the assumption that these isotopes have been produced mainly with cosmological nucleosynthesis. From these considerations he finds the following limit values for the baryon number and the lepton number:

(71) $\quad B < 3\text{x}10^{-9} \quad ; \quad -0.3 < L_e < 4$

and a very weak indication that $B = 1.5\text{x}10^{-9}$ or $\rho_B \approx 10^{-30}$ g-cm^{-3} and $L_e \approx -0.18$. This last value would imply different densities for ν and $\bar{\nu}$ with $n_{\bar{\nu}} > n_{\nu}$.

Many further investigations of the element abundances and other model calculations will be necessary in order to show whether this asymmetry does indeed exist and to establish a further relationship between the free parameters of the Hot Big Bang and the physical properties of the present Universe.

The author is highly indebted to Prof. G. Marx for the invitation to the Europhysics Conference Neutrino '72 and for the kind hospitality in Hungary.

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COSMOLOGICAL LIMIT ON NEUTRETTO MASS (DISCUSSION REMARK)

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By making use of the general relativistic cosmology, the rest mass of the muon neutrino could be obtained, if the Hubble constant, the deceleration parameter and the photon temperature were accurately known. In this note the decoupling temperature of neutrinos is discussed. From the available astronomical data the upper limit of muon neutrino rest mass has been obtained to be 130 eV.

It was stressed by Zel'dovich as first, that the empirical cosmology offers an approach to obtain the value or at least the upper limit of the neutrino masses [1,2]. Zel'dovich considered a homogenous and isotropic expanding universe, which was dominated by electromagnetic and neutrino radiation. The number density of the relict neutrinos has been determined by the present photon temperature and by the thermal history of the Big Bang. An upper limit on the average mass density can be obtained from the geometric history of the expansion and from the fact, that the universe is older than the oldest moon rock /4,5 Gy, where $1Gy=10^9$ years/. By dividing this mass density limit with the neutrino number density Zel'dovich has obtained an upper limit for the neutrino rest mass.

Our aim was to make a rather complete computer simulation of the thermal and geometric history of the universe, in order to explore the reliability of this approach and to arrive at an empirical neutretto mass.

After the Hadron Era /$kT < m_\pi c^2$/ the universe was made of photons, leptons and a negligible amount of nucleons. In this Lepton Era the muons started annihilating as first. The neutrettos are coupled to the muons / $\nu_\mu + \mu \to \mu + \nu_\mu$ / via conventional weak interactions [3] / and in a smaller amount to the electrons / $\nu_\mu + e \to e + \nu_\mu$ via the electromagnetic form-factor of the neutretto [4] and via the hypothetical W^0 boson coupling [5]/. The decoupling from muons happened at $T_0 = 1.2.10^{11}$K [3]. If one uses a cut-off parameter Λ =300 GeV, according to our numerical calculations the electromagnetic decoupling from the electron plasma occured even earlier, at T=4,1.10^{11}K. If we had a $\bar{\nu}_\mu \nu_\mu \to W^0 \to \bar{e}e$ weak interaction [5], the decoupling would happen somewhere in the temperature range $10^{11}\mathrm{K} > T > 10^{10}$K. The decoupling of electron neutrinos from the electron plasma came later, at T=1,8.10^{10}K [3]. Finally the electron pairs annihilated and the photons were decoupled from the surviving hydrogen-helium plasma.

Before decoupling all the components were in thermodynamical equilibrium having a common temperature T /t/, and the cooling was adiabatic. After having decoupled, the wave length of each particle changed proportional to the radius R /t/ of the universe. We used vanishing rest mass for the photon and for the neutrino, the known rest mass for electron and positron, and we left the neutretto rest mass as a free parameter.

The computer calculation started at the neutretto decoupling temperature T =1,2.10^{11}K, with an equilibrium distribution of photons, electron pairs, neutrinos and neutrettos. The corresponding radius R_0 was also considered as free parameter. The geometric history is described by the Einstein equation

$$\left(\frac{\dot{R}}{R}\right)^2 + K\frac{c^2}{R^2} = \frac{8\pi G}{3}\rho$$

/K=+1 for closed space./ ϱ and p has been given for photons and leptons in terms of T /t/ before decoupling, in terms of R /t/ after decoupling. The connection between R /t/ and T /t/ has been given by the adiabatic equation /which is a direct consequence of the Einstein equations/:

$$\frac{d}{dt}(\varrho c^2 R^3) + p\frac{d}{dt}(R^3) = 0$$

/After the decoupling point of the photons one has had evidently R /t/. T /t/=const./ The integration was performed with a CDC 3300 computer.

In the first period the kinetic energy dominated / $\varrho = 3p \sim T^4 \sim R^{-4}$ /, one had a radiation-filled universe. At about the temperature $kT_x = m_\nu c^2$ the neutretto rest energy $m_\nu c^2$ became the dominant component / $\varrho \sim R^{-3}$, $p=0$ /. The neutrettos were so abundant, that the atomic mass density was negligible during the whole history.

The integration stopped, when the photon temperature reached the present observed value

$$T_\gamma(t_0) = 2{,}7 \pm 0{,}3\ K$$

At this point the solution has to obey the empirical restrictions [6]:

$$t_0 > 4{,}5\ Gy,$$

$$R(t_0)/\dot{R}(t_0) = H_0^{-1} = 18{,}4 \pm 2\ Gy,$$

$$-R(t_0)\ddot{R}(t_0)/\dot{R}(t_0)^2 = q_0 = 1 \pm 1.$$

These restrictions are fulfilled only for special choises of m_ν and R_0. Our numerical results are collected in the Table. The equalities are valid, if below the temperature T_x the neutretto rest mass density has been dominating indeed. If our actual universe contained a considerable amount of unknown forms of matter /gravitons, black holes, intergalactic plasma/, this would make upper limits from the neutretto mass values.

The limit m_ν=130 eV for the neutretto rest mass obtained from cosmology is by four orders of magnitude lower, than the value quoted in the Review of Particle Properties /$m_\nu \leq$1,2 MeV / [7].

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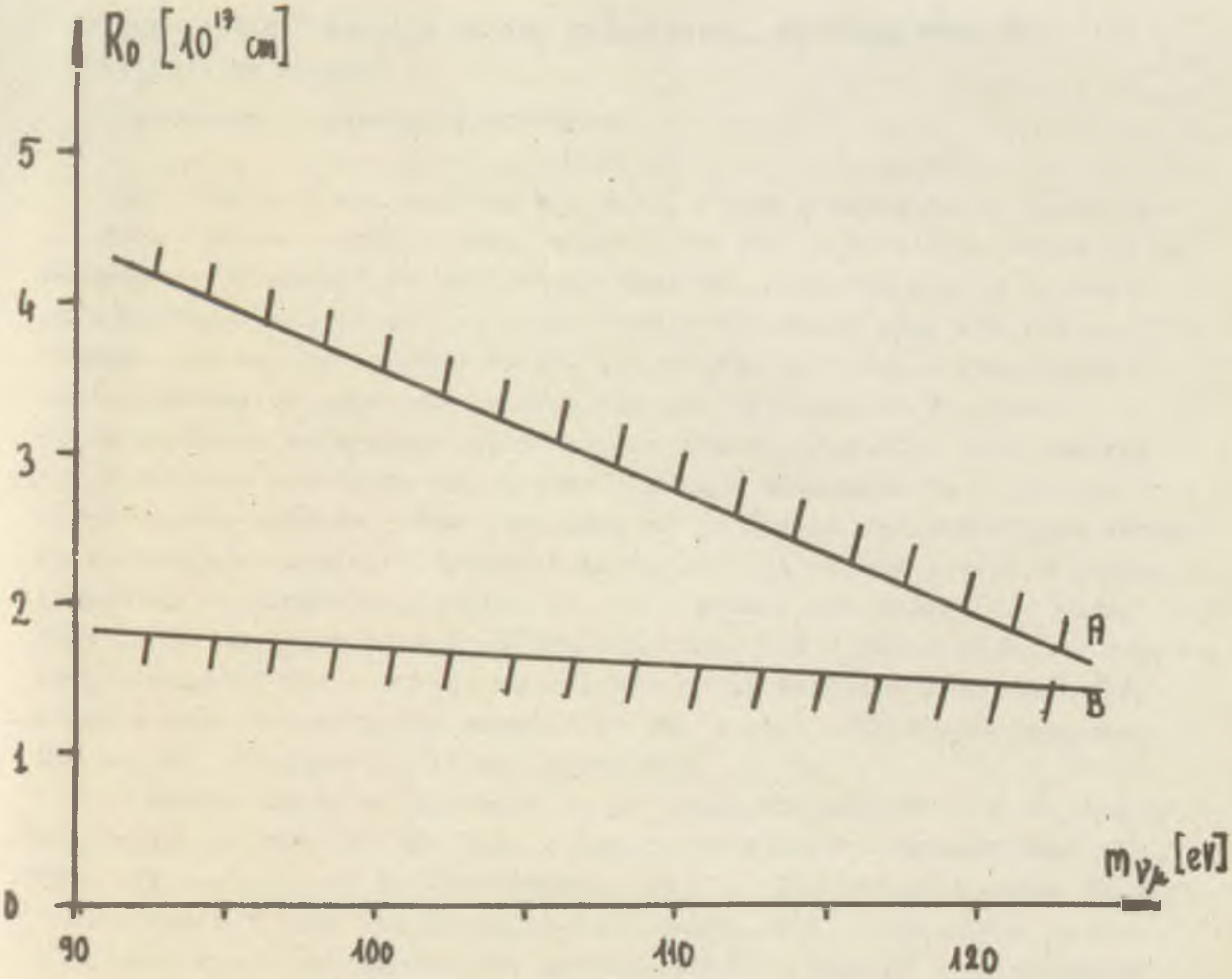

Curve A corresponds to the value H_0= 75 km/s.Mpc

The expansion of the Universe stops under curve B before reaching 2.7 $^{\circ}$K.

The allowed area for the parameters is between the two curves, so the upper limit to the ν_μ mass is 130 eV.

FIELD-PARTICLE ASPECTS OF THE COSMOLOGICAL NEUTRINO PROBLEM (DISCUSSION REMARK)

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Not too much was said on the role of the gravitational field in treating the neutrino; a brief summary of the recnt status seems to be necessary. It should be emphasized that the neutrino may be treated as a particle and as a field, and both a classical approach and a quantum one may be applied to it. The first, and almost completely exploited way of approach is to treat the neutrino as a classical point particle in a given external gravitational field. This was the way I used to calculate the neutrino escape hindrance from general relativistic spheres which were applied as models for superdense stars. This classical neutrino dynamics is equivalent to the search for zero geodesics confined to a region of space around the superdense star; this is in the case of a massless neutrino. The results might be more complicated if the neutrino were not a point particle with no higher moments than the monopole moment (as for bodies with higher moments the motion, in general, is not geodetic).

A second level of approach is to treat the neutrino in a quantum way, with the help of the Weyl equation in a given gravitational field. This way was applied by WheelerBrill et al. The use of movable Cartan frames may but need not be applied. Results on the metastable states of a neutrino in an equivalent gravitational potential were obtained a long time ago (when applications to the geon theory were sought for). Recently, among the most important results on these lines are those of Pestov on the connection between duality rotations for electromagnetic and for neutrino fields (as different representations of the same transformation group), and of Novello and Rotelli on a link between gravitation and weak interactions. The latter result, for homogeneous and isotropic cosmological models, is that the axial vector current in the UFI gives varying contributions when compared with the vector current at different cosmological times. In consequence, neutrinos and antineutrinos produced are admixtures of both left and right polarized states, and the mixing ratio depends on the time they were created; this may have important consequences for the cosmological neutrino sea. Along this line of research it has been proved, too, that when we start with the Dirac equation for a massless neutrino in curved space-time, an induced mass term appears in the final equations.

Up to now only the problem of a neutrino in a given gravitational field was considered. But what would result if the gravitational field would have as its source the energy-momentum tensor of the neutrino field.

This approach leads to a study of the coupled Einstein-Weyl equations. It has been developed recently, with the help of modern tools (ray optics, spin coefficient formalism etc.), and has given some interesting results. It has given some kind of Rainich-type already unified field theory of the neutrino and gravitation. Several classes of neutrino fields were investigated, especially those with a positive energdy density (class E_+). It was found that the principal null congruence of any neutrino field of this class is geodesic, and its shear σ and twist ω are restricted by the condition: $|\sigma|^2 - 4\omega^2 \leq 0$. If a neutrino field of this class is twistfree, it must be also shearfree. For neutrino pure radiation fields, with an energy-momentum tensor of the form: $T_{\mu\nu} = \lambda^2 k_\mu k_\nu$, where the tetrad vector k_μ is the neutrino flux vector, the theorem has been obtained :

All radiation solutions of the Weyl equation possess a shearfree and twistfree geodesic congruence which is also hypersurface-orthogonal. This is a stronger theorem than the corresponding Mariot-Robinson theorem for null electromagnetic fields.It was proved also that spherical symmetry of a space-time is not compatible with the presence of a neutrino field subject to the restriction that the energy flow vector of the neutrino field be timelike or null.

Other important results in this area are: the Plebański classification of the neutrino energy-momentum tensor, necessary and sufficient conditions on the Ricci tensor to admit a neutrino field of some classes,,asymptotic domination of radiation terms in the energy-momentum tensor, finally some exact solutions. The most contributions in this area have been obtained by Wainwright, Trim, Griffiths, Newing, Audretsch and Graf.

While in this appraoch the neutrino field was treated in a classical way, without the inclusion of the Pauli principle, a possible, fourth line of approach is to study a quantized neutrino field, possibly in a quantized space-time.

The general scheme of developments, and possible further directions of approach, may be as below:

	Classical approach		Quantum approach
Particle aspects	Geodesic motion of a ν in a given grav.field (also nongeod.motion, in some future)	$\longrightarrow$	Metastable states of a ν in a grav. potential;modified UFI theory
Field aspects	Already unified field theory of ν and gravit.	$\longrightarrow$	Necessity to include second quantization of the ν field; what more?

REVIEW OF THE EXPERIMENTAL SITUATION ON NEUTRAL CURRENTS

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I) INTRODUCTION

The weak interactions are presently understood in terms of a current-current interaction. All experimentally observed phenomena so far require only the charge-changing currents $(e^-, \bar{\nu}_e)$, (e^+, ν_e), $(\mu^-, \bar{\nu}_\mu)$, (μ^+, ν_μ), $(n, \bar{p})$, $(\Lambda, \bar{p})$ etc.; there is no evidence for the existence of the weak neutral currents $(\nu_e, \bar{\nu}_e)$, $(\nu_\mu, \bar{\nu}_\mu)$, (e^+, e^-), (μ^+, μ^-), $(p, \bar{p})$, etc.

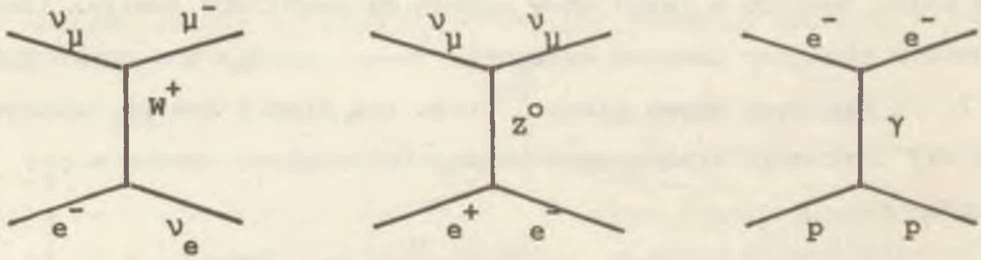
Several recently proposed models of the weak and electromagnetic interactions, in particular the model of Weinberg [1)], that show promise of being renormalizeable, predict the existence of neutral currents. However, the extremely sensitive experiments searching for neutral currents in K decays have shown that strangeness changing neutral currents do not exist down to a level many orders of magnitude smaller than the strangeness changing charged currents. These results are summarized in Table I. It has been shown since [2)] that the theory can be arranged in such a way that only strangeness conserving neutral currents are required.

The effects of the neutral current terms (e^+, e^-) and (μ^+, μ^-) are masked by the electromagnetic interactions. The $(\nu_e, \bar{\nu}_e)$ and $(\nu_\mu, \bar{\nu}_\mu)$ terms would, however, give rise to effects that are experimentally detectable in the interactions of neutrinos with nucleons and electrons.

A number of experiments have been completed to search for such effects. No evidence for neutral currents was found. The approximate upper limits obtained in these experiments are summarized in Table II. These limits are of the same order of magnitude as the ratios predicted by the Weinberg model.

The purpose of this talk is to discuss the experimental limits and the predictions of the Weinberg model in more detail, so that a valid comparison can be made. Two points in particular come to attention :

1. The model does not predict a unique value for the rate of the neutral current reactions, but rather a range of values, depending on a free parameter in the model which is related to θ, a mixing angle in the following way. The model hypothesizes, among other things, an isotriplet of intermediate vector bosons, two charged and one neutral, the W^+, W^-, and W^o, and an isosinglet vector boson B^o. The $W^{\pm}$ as usual, mediate the charge changing part of the weak interactions. The W^o and B^o mix to produce the Z^o that mediates the neutral current weak interactions, and the photon that mediates the electromagnetic interactions, as, for example, in the diagrams below.



$\nu_\mu + e^- \rightarrow \mu^- + \nu_e$ $\qquad$ $\nu_\mu + e^- \rightarrow \nu_\mu + e^-$ $\qquad$ $e^- + p \rightarrow e^- + p$

The mixing can be written in terms of some mixing angle θ as

$$Z^o = \cos\theta\ W^o + \sin\theta\ B^o$$
$$\gamma = -\sin\theta\ W^o + \cos\theta\ B^o$$

If the coupling constant of the isotriplet intermediate bosons to the lepton current is written as g, and the coupling of the isosinglet to the lepton current as g',
then the mixing angle is defined as

$$\sin\theta = g'/\sqrt{g^2 + g'^2}$$
$$\cos\theta = g\ /\sqrt{g^2 + g'^2}$$

and the electromagnetic coupling constant is

$$e = gg'/\sqrt{g^2 + g'^2}$$

The rates for the neutral current reactions depend on the ratio of the electromagnetic to the isotriplet coupling constants e^2/g^2, which from the above can be seen to the

$$e^2/g^2 = \sin^2\theta$$

and is thus limited between 0 and 1.

$$0 < e^2/g^2 < 1.$$

The masses of the intermediate bosons are related to these coupling constants as

$$M^2_{W^\pm} = \frac{\sqrt{2}\ g^2}{8\ G}$$

$$M^2_{Z^o} = \frac{\sqrt{2}\ (g^2 + g'^2)}{8\ G} = \frac{M^2_{W^\pm}}{(1 - e^2/g^2)}$$

where $G \cong 10^{-5}/m_p^2$ is the usual beta decay coupling constant. The numerical values of these masses as a function of e^2/g^2 are shown in Fig. 1.

2. The energy spectra of the final state electron or hadrons in the neutral current processes depend quite strongly on e^2/g^2. Because of the detection methods used and the presence of various backgrounds, the experiments searching for the neutral current reactions are sensitive only to a limited range of the final state electron or hadron energies. Thus the detection efficiencies of the experiments, and therefore the rate or upper limit obtained, depend on e^2/g^2 in the Weinberg model (or on some other, possibly more complicated, parameters in other models). The experiments will thus be discussed individually with some attempt to calculate the detection efficiencies and the expected rates in the Weinberg model.

The predictions of the model for the purely leptonic processes are straight forward and free of any uncertainties related to the poorly understood strong interactions of the hadrons. As the complexities of the semileptonic processes increase from the elastic ($\nu_\mu + p \rightarrow \nu_\mu + p$) to single pion production ($\nu_\mu + N \rightarrow \nu_\mu + N + \pi$) and finally the inclusive inelastic process ($\nu_\mu + N \rightarrow \nu_\mu$ + hadrons, the uncertainties associated with the theoretical predictions of the model increase rapidly. Thus a clearcut test of the model is best carried out in the purely leptonic processes.

II) PURELY LEPTONIC PROCESSES

II.1 Search for $\nu_\mu + e^- \rightarrow \nu_\mu + e^-$

The process

$$\nu_\mu + e^- \rightarrow \nu_\mu + e^- \tag{1}$$

can proceed only via neutral currents. It has been looked for in the CERN neutrino experiments, both in the heavy liquid bubble chamber [4)] and the spark chamber detectors [3]. The event is characterized as a single e^- produced within a few degrees of the ν beam direction, without any accompanying hadrons or nuclear fragments. The most sensitive search was in the 1963-64 run in the spark chambers [3]. To insure good electron detection and to eliminate the low energy electron background, a sample of 30 events with electrons over 1 Bev was selected. These events were probably due to the reaction

$$\nu_e + (n \text{ or } p) \rightarrow e^- + \text{hadrons} \tag{2}$$

Of this sample only one event had an electron within 4^o of the ν beam direction and had no accompanying hadrons and could thus be an example of reaction (1). Since this event could also be due to reaction (2), this one event can be used to set an upper limit only. In reference (3) this event was used to set a limit on the reaction

$$\nu_e + e^- \rightarrow \nu_e + e^- \tag{3}$$

which has the identical signature as reaction (1). The normalization given in that paper is calculated for the ν_e spectrum. To estimate the normalization for the ν_μ induced reaction (1), one could use the ratio of the calculated ν_μ and ν_e fluxes. However, in view of the difficulty of estimating the absolute neutrino fluxes, it seems preferable to normalize to the total number of ν_μ induced reactions, and use the measured [6] ν_μ total cross section, $(0.8 \pm 0.2)\, E_\nu \times 10^{-38}$ cm^2/nucleon. The detection efficiency used in ref. 3 was calculated for reaction (3), which in the usual V-A theory is uniformly distributed in E_e/E_ν. This detection efficiency gives the limit listed in Table II.

The differential cross section [7] for the neutrino-electron scattering process is

$$\frac{d\sigma}{dE_e} = \frac{G^2 M_e}{2\pi}\left[(g_V + g_A)^2 + (g_V - g_A)^2\left(1 - \frac{E_e}{E_\nu}\right)^2 + \frac{M_e E_e}{E_\nu^2}(g_A^2 - g_V^2)\right] \tag{4}$$

where E_e, E_ν are the laboratory energies of the e^- and ν, respectively, and M_e is the electron mass. In the V-A theory, $g_A = g_V = 1$. In the Weinberg model they take on the following values for the four neutrino-electron elastic scattering processes :

Process	g_V	g_A	
$\nu_e + e^- \rightarrow \nu_e + e^-$	$+ \frac{1}{2} + 2\, e^2/g^2$	$+ \frac{1}{2}$	
$\bar{\nu}_e + e^- \rightarrow \bar{\nu}_e + e^-$	$+ \frac{1}{2} + 2\, e^2/g^2$	$- \frac{1}{2}$	
$\nu_\mu + e^- \rightarrow \nu_\mu + e^-$	$- \frac{1}{2} + 2\, e^2/g^2$	$- \frac{1}{2}$	(5)
$\bar{\nu}_\mu + e^- \rightarrow \bar{\nu}_\mu + e^-$	$- \frac{1}{2} + 2\, e^2/g^2$	$+ \frac{1}{2}$	

Fig. 2 shows the cross sections, which rise linearly with E_ν, expected for these four processes in this model as a function of the coupling constant ratio e^2/g^2. The total cross section for reaction (1) varies from 0.1 to 1.0 x 10^{-41} E_ν cm^2/electron. The cross section for the process $\nu_e + e^- \rightarrow \nu_e + e^-$ in the V-A theory, 1.6 x 10^{-41} E_ν cm^2/electron, is also indicated on Fig. 2.

To estimate the detection efficiency for reaction (1) in this experiment, we have written a Monte Carlo program, using the differential cross section (4) and the calculated flux [8] of ν_μ in the experiment. With an $E_e > 1.0$ BeV cut, the detection efficiency varies from 20 % to 40 %, as a function of e^2/g^2, as shown in Fig. 3. Using this detection efficiency, and normalizing to the total number of ν_μ interactions, and recalling that there are ~ ½ as many electrons as nucleons in Aluminium; we arrive at a 90 % confidence level upper limit on reaction (1) from this experiment of

$$\sigma(\nu_\mu + e^- \rightarrow \nu_\mu + e^-) \leqslant (1.8 \text{ to } 3.6) \times E_\nu \times 10^{-41} \text{ cm}^2/\text{electron}$$

The upper limit is plotted as a function of e^2/g^2 in Fig. 4, which also shows the cross sections predicted by the Weinberg model. As can be seen, the limits are not inconsistent with the model.

II.2 Search for $\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-$

The process

$$\bar{\nu}_e + e^- \to \bar{\nu}_e + e^- \qquad (6)$$

can proceed in the usual V-A theory with a total cross section of $0.54 \times 10^{-41}\ E_\nu\ cm^2$/electron. The contribution of the neutral currents to this process in the Weinberg model can either depress or enhance this cross section, depending on the value of e^2/g^2. Using the differential cross section (4) with the coupling constants (5), we calculate a cross section varying from $0.14 \times E_\nu \times 10^{-41}\ cm^2$/el to $2.9 \times E_\nu \times 10^{-41}\ cm^2$/el, as shown in Fig. 2. Thus a measurement of the cross section for this process, even though it is allowed via the usual charged currents, can yield some information relevant to neutral currents.

An experimental search for this process has been carried out [(9)] using a fission reactor as a source of low energy $\bar{\nu}_e$. The recoil electrons were detected in a ~ 8 kg plastic scintillator that provided a ± 2 % measurement of the electron energy. The detector was surrounded by an elaborate array of background shielding and anticoincidence counters. In order to reduce the remaining background, the energy of the electrons accepted as candidates for reaction (6) were restricted to $3.6 < E_e < 5.0$ MeV. No signal corresponding to reaction (6) was observed after a background subtraction, the latest value of the upper limit has been reported in a recent article [(9)] as $\sigma(\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-) < 1.9\ \sigma_{V-A}$. This is a one standard deviation limit. The 90 % confidence level upper limit corresponds to

$$\sigma(\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-) \leqslant 3\ \sigma_{V-A}$$

This limit is based on the detection efficiency calculated assuming the electron energy spectrum as predicted by the V-A theory. We have calculated the detection efficiency in the Weinberg model, using the differential cross section (4) with the coupling constant (5) and the $\bar{\nu}_e$ spectrum calculated by Avignone [10]. The detection efficiency (probability of getting an electron with $3.6 < E_e < 5.0$ MeV) varies by a factor of ~ 6 over the range of e^2/g^2, from $\sim 1/3$ % to ~ 2 %, as shown in Fig. 5. This rapid variation is due to the fact that the bulk of the $\bar{\nu}_e$ spectrum is below 5 MeV, and thus the efficiency is critically dependent on the shape of the electron energy spectrum near its maximum. The upper limit, using this detection efficiency, is plotted against e^2/g^2 in Fig. 6, together with the cross section predicted by the Weinberg model. As can be seen from this figure, if this experimental result is correct, the upper region of e^2/g^2 is ruled out. The 90 % confidence level upper limit [11] on e^2/g^2 is

$$e^2/g^2 \leqslant 0.35$$

The corresponding limit on the mixing angle is

$$\theta \leqslant 36^{\circ}$$

Another order of magnitude decrease in this experimental upper limit would be needed to rule the Weinberg model out altogether.

III) SEMI-LEPTONIC PROCESSES

III.1 The Elastic Process $\nu_\mu + p \rightarrow \nu_\mu + p$

The elastic neutrino-proton scattering process

$$\nu_\mu + p \rightarrow \nu_\mu + p \tag{7}$$

can proceed only via neutral currents. This process was searched for in the CERN heavy liquid bubble chamber neutrino experiment [4].

The signature for this reaction is a single proton. The main experimental problem was the large number of neutron induced background. This background was reduced by requiring that the proton have a kinetic energy between 150 and 500 MeV, (this corresponds to a cut of $0.3 < q^2 < 1.0$ BeV) and that the incident neutrino energy, computed from the recoil proton angle, lie between 1 and 4 BeV. A total of 4 such events, were found. Since these could also be neutron induced events, only an upper limit on this process can be set. Using the sample of $\nu_\mu + n \rightarrow \mu^- + p$ events, with the same kinematic cuts on the proton, the 90 % confidence level upper limit on their ratio is

$$\frac{\nu_\mu + p \rightarrow \nu_\mu + p}{\nu_\mu + n \rightarrow \mu^- + p} \leqslant 0.24 \tag{8}$$

This limit assumes that the proton energy spectrum is identical for the two processes, and thus their detection efficiencies with the same kinematic cuts are the same. This is not the case in the Weinberg model for values of e^2/g^2 other than 0.

The prediction of the Weinberg model for this ratio, as well as the relative detection efficiency, can be calculated using the differential cross section for these processes

$$\frac{d\sigma}{dq^2} = \frac{G^2}{8\pi}\left\{ (g_A^2 - g_V^2)\frac{q^2}{2E_\nu^2} + (g_A + g_V)^2 + (g_A - g_V)^2\left(1 - \frac{q^2}{2ME_\nu}\right)^2 + \left[(4M^2 + q^2)\, f_V - 4M f_V g_V\right]\left(2 - \frac{q^2\,(2E_\nu + M)}{2ME_\nu^2}\right)\right\} \tag{9}$$

where E_ν is the neutrino lab energy, M is the nucleon mass, and q^2 is the square of the four momentum transfer.

The values of the coupling constants are given by Weinberg [12)] as

	$\nu_\mu + n \to \mu^- + p$	$\nu_\mu + p \to \nu_\mu + p$
g_V	$4.7\ F_V\ (q^2)$	$(2.35 - 5.58\ e^2/g^2)\ F_V\ (q^2)$
$g_V - 2\ Mf_V$	$1.0\ F_V\ (q^2)$	$(\frac{1}{2} - 2\ e^2/g^2)\ F_V\ (q^2)$
g_A	$1.2\ F_A\ (q^2)$	$0.6\ F_A\ (q^2)$

The vector and axial vector form factors (using CVC) can be written as

$$F_V\ (q^2) = \frac{1}{(1 - q^2/0.84^2)^2}$$

$$F_A = (q^2)\ \frac{1}{(1 - q^2/M_A^2)^2}$$

The only undetermined quantity is M_A, which can be taken to be $M_A = 0.84$ BeV (this is consistent with the experiments studying $\nu_\mu + n \to \mu^- + p$). Using a Monte Carlo program, with the 1967 CERN ν_μ spectrum [14)], the prediction of the model for the ratio (8) and the relative detection efficiency for the kinematic cuts used in the experiment, have been calculated, and are shown in Figure 7, as a function of e^2/g^2. The experimental upper limit, using the calculated relative detection efficiency, is also indicated. As can be seen from this figure, the experimental upper limit falls below the predicted ratio only for $e^2/g^2 \geqslant 0.9$. However, in view of the assumptions and uncertainties in the knowledge of the nucleon form factors, no strong conclusions can be drawn from this. For example, the predicted cross section changes by about a factor of 2 as M_A is varied by a factor of 2.

III.2 Single Pion or Δ Production

Several experiments have looked for the effects of neutral currents in single pion production by neutrinos.

1. The 1967 CERN neutrino experiment in the propane bubble chamber looked for the reaction [4]

$$\nu_\mu + p \rightarrow \nu_\mu + n + \pi^+ \tag{10}$$

The signal for this reaction was a single π^+ in the final state identified by a visible interaction or decay; 4 such events were found. The reaction

$$\nu_\mu + p \rightarrow \mu^- + p + \pi^+$$

was used for comparison; 51 such events were found, using the same criteria as for reaction (10) to identify the π^+. Since the four candidates for reaction (10) could also be due to the neutron background, they were used to give the 90 % confidence level upper limit of

$$\frac{\nu_\mu + p \rightarrow \nu_\mu + n + \pi^+}{\nu_\mu + p \rightarrow \mu^- + p + \pi^+} < 0.16 \tag{11}$$

In the limit of complete Δ (1238) dominance, a prediction for this ratio at $e^2/g^2 = 0$ was given by Weinberg [12] as 1/9. The dependence on e^2/g^2 was estimated by B.W. Lee [15], and is shown in Fig. 8. Since the predicted cross section falls off from 1/9 as e^2/g^2 increases, the experimental limit (11) does not constrain the model.

2. The neutral current processes $\nu_\mu + n \to \nu_\mu + n + \pi^0$ and $\nu_\mu + p \to \nu_\mu + p + \pi^0$ were looked for in a recent re-analysis of the 1963 Columbia-Brookhaven neutrino experiment[5)] using an optical spark chamber detector. All events with one or more γ showers but no muons (a muon was defined as a track that penetrates at least 3 interaction lengths of chamber without a visible interaction) were taken as candidates. Five such events were found with fairly energetic showers, which were probably due to electron showers from the reaction $\nu_e + N \to e^- +$ hadrons. Since the π's in the single pion processes were expected to be less energetic, a cut which corresponds roughly to $160 < E_{\pi^0} < 400$ MeV was imposed. No candidates remained after this cut. The reaction

$$\nu_\mu + n \to \mu^- + p + \pi^0 \qquad (12)$$

was used as a comparison. Candidates for this reaction were selected as events with a muon track and one or more showers, using the same cut on the energy of the showers as in the case of the neutral current events. 12 such events were found; 9 of these were classified as examples of reaction (12) and 3 as double pion production. However, lack of visible detail in the spark chambers and the effects of the aluminum nucleus make this classification subject to some assumptions. A further problem is that in the reaction $\nu_\mu + p \to \mu^- + p + \pi^+$, which might be 9/2 as frequent as reaction (12), the π^+ could charge exchange and produce a π^0, which could be responsible for some of the 9 events assigned to reaction (12). Neglecting these problems, a 90 % upper limit was derived for the ratio

$$R = \frac{(\nu_\mu + n \to \nu_\mu + n + \pi^0) + (\nu_\mu + p \to \nu_\mu + p + \pi^0)}{2\,(\nu_\mu + n \to \mu^- + p + \pi^0)} < 0.14 \qquad (13)$$

The prediction of the Weinberg model for this ratio is very difficult to estimate because of the poor understanding of the strong interactions at the hadronic vertex. A calculation, making a number of assumptions, was carried out by B.W. Lee [15] in the Weinberg model, predicting the ratio R to vary from 0.2 to 1.0 as a function of e^2/g^2, as shown in Fig. 9. The upper limit of 0.14 is inconsistent even with the lowest value of 0.2. Furthermore, if the region of $e^2/g^2 > 0.35$ is excluded by the limit [(9)] on $\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-$, as discussed above (see Fig. 6), then the predicted range of R shrinks to 0.4 to 1.0, making the disagreement with experiment even worse.

The experimental limit (13) is based on the assumption that the angular and energy distributions of the π^0's in the reactions involved in the ratio R are sufficiently similar that their detection efficiencies are the same. This might not be a good assumption for values of e^2/g^2 other than 0. We have made no attempt to evaluate this detection efficiency as a function of e^2/g^2 as we have done for the other reactions discussed above, because of the complexity and uncertainty of the theory of the single pion production process.

Many calculations of the single pion production process by neutrinos via the conventional charged current have been carried out in the past. These calculations disagree with each other and experiment by factors of two or three. Some of the uncertainties, but not all of them, may cancel out in the ratio R for values of e^2/g^2 other than 0.

In view of these difficulties it is not clear how serious the disagreement between the experimental upper limit (13) and the Weinberg model is.

IV. FUTURE PROSPECTS

IV.1 In the 1971-1972 neutrino experiment at the Argonne National Laboratory a total of 700'000 pictures have been taken in the 12 foot bubble chamber, about half with Hydrogen and half with a Deuterium fill. A search for the neutral current reactions

$$\nu_\mu + p \rightarrow \nu_\mu + n + \pi^0 \tag{14}$$

$$\nu_\mu + p \rightarrow \nu_\mu + p + \pi^0 \tag{15}$$

is under way. Reaction (14) is identified by a π^+ stopping and decaying in the chamber, which is expected to happen ~ 25 % of the time. Reaction (15) is identified by one or more $\gamma \rightarrow e^+ e^-$ pairs associated with a single positive track. The probability of at least one of the γ's from the π^0 converting is estimated to be 12 %. The total exposure is expected to yield 200 events of the reaction $\nu_\mu + p \rightarrow \mu^- + p + \pi^+$. Thus if no events of the type (15) or (16) are found, a limit of around 0.15 could be set on the ratio $(\nu_\mu + p \rightarrow \nu_\mu + n + \pi^+ + \nu_\mu + p \rightarrow \nu_\mu + p + \pi^0) / (\nu_\mu + p \rightarrow \mu^- + p + \pi^+)$. This ratio can be expected to be around 1/3 at $e^2/g^2 = 0$ in the Weinberg model, with an e^2/g^2 dependence similar to that in Figs. 8 and 9. Results from this experiment should be forthcoming in a few months.

IV.2 In the 1971 CERN neutrino experiment using the large heavy liquid bubble chamber Gargamelle filled with freon, 500'000 photographs have been taken, half with an incident ν_μ beam and half with $\bar{\nu}_\mu$. Searches for various neutral current reactions like $\nu_\mu + p \rightarrow \nu_\mu + p + \pi^0$, $\nu_\mu + p \rightarrow \nu_\mu + n + \pi^+$, and the inclusive inelastic process ν_μ + (n or p) $\rightarrow \nu_\mu$+hadrons are under way. The signature for these processes are events with interacting pions, π^0 showers, or other identified hadron final states without the presence of a muon. Since the radiation length in freon is 11 cm and the interaction length is 70 cm, and Gargamelle is 1.95 m diameter x 4.8 m long, the efficiency of identifying these final states should be quite high.

The most serious experimental problem is the possible presence of a sizeable neutron background that would produce hadrons without a muon in the chamber. Since the chamber is over six interaction lengths long, the spatial distribution of any observed events might distinguish neutral current events from the neutron background. The total number of neutrino induced events are expected to be between 5000 and 10'000 in the run already completed, so that sensitivities of the order of 5 % compared to the corresponding charged current reactions might be achieved.

Another search under way in the Gargamelle experiment is for the purely leptonic neutral current processes $\nu_\mu + e^- \to \nu_\mu + e^-$ and $\bar{\nu}_\mu + e^- \to \bar{\nu}_\mu + e^-$. The signature for these events is a single e^- within a few degrees of the ν beam direction without any additional hadrons or nuclear fragments. Preliminary results show that this search is very background free down to an electron energy of 300 MeV. With such a low energy cut, and the relatively high ν energy of the new CERN beam (the spectrum peaks around 2 BeV) the detection efficiency for these processes should be 70 to 85 %. The total number of ν_μ events expected if the Weinberg model is correct is between 1 ½ and 15, depending on the value of e^2/g^2, as shown in Fig. 10.

Another neutrino run of 500'000 pictures is scheduled in Gargamelle for the fall of 1972. This run would double the sensitivities mentioned above.

IV.3 The anticipated neutrino experiments at NAL should greatly improve the sensitivity of neutral current searches in neutrino interactions. For example, the search for the purely leptonic process $\nu_\mu + e^- \to \nu_\mu + e^-$ should be quite straightforward in the 15 foot bubble chamber filled with neon. The radiation length in neon is around 25 cm, so that unambiguous electron identification with high detection efficiency should be possible. The detailed visibility of the interaction vertex possible in a bubble chamber should allow the rejection of the major background, $\nu_e + N \to e^- +$ hadrons, by the requirement that no hadrons or nuclear fragments accompany the electron. In a 10^6 picture run, assuming 5×10^{12} protons per pulse at 350 BeV, and a 20 m^3 fiducial volume, the range of the number of events expected in the Weinberg model is between 100 and 1000, as shown in Fig. 11.

Table I

Summary of upper limits on branching ratios involving neutral currents in K decays. (From the summmary of the Particle Data Group, Physics Letters 39 B, 1 April 1972)

Decay	Branching Ratio
$K_L^o \to e^+e^-$	$< 1.6 \times 10^{-9}$
$\mu^+\mu^-$	$< 1.9 \times 10^{-9}$
$e^\pm\mu^\mp$	$< 1.6 \times 10^{-9}$
$K^\pm \to \pi^\pm e^+e^-$	$< 0.4 \times 10^{-6}$
$\pi^\pm\mu^+\mu^-$	$< 2.4 \times 10^{-6}$
$\pi^\pm\nu\bar{\nu}$	$< 1.2 \times 10^{-6}$

Table II

Summary of Neutral Current Searches in Neutrino Interactions

Cross section Ratio	Ref.	Approx. upper limit [(a)]
$\dfrac{\nu_\mu + e^- \to \nu_\mu + e^-}{\text{V-A theory for } \nu_e + e^- \to \nu_e + e^-}$	3	1.0 (90 % conf. level)
$\dfrac{\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-}{\text{V-A theory for } \bar{\nu}_e + e^- \to \bar{\nu}_e + e^-}$	9	1.9 (1 stand. dev. limit)
$\dfrac{\nu_\mu + p \to \nu_\mu + p}{\nu_\mu + n \to \mu^- + p}$	4	$0.12 \pm .06$
$\dfrac{\nu_\mu + p \to \nu_\mu + n + \pi^+}{\nu_\mu + p \to \mu^- + p + \pi^+}$	4	$0.08 \pm .04$
$\dfrac{(\nu_\mu + n \to \nu_\mu + n + \pi^0) + (\nu_\mu + p \to \nu_\mu + p + \pi^0)}{2\ (\nu_\mu + p \to \mu^- + p + \pi^0)}$	5	0.14 (90 % conf. level)
$\dfrac{\nu_\mu + N \to \nu_\mu + \text{hadrons}}{\nu_\mu + N \to \mu^- + \text{hadrons}}$		see footnote (b)

FOOTNOTES TO TABLE II

a) The upper limits quoted depend on the detection efficiencies for the final state electron or hadrons in the neutral current reactions. The above limits are based on the assumption that the energy spectra of the detected particles in the neutral current reactions are identical to those in the corresponding charged current reactions. As discussed in sections II and III, this assumption could be quite wrong, and the upper limits could vary by factors of two or three, depending on the detailed structure of the neutral current reactions.

b) Upper limits varying from 0.01 to 0.05 for this ratio can be found in the literature. However, all of these limits have been rescinded by the authors, and the only reliable statement that can be made at the moment is that the ratio is not larger than 1.

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11. This conclusion is consistent with a similar analysis carried out by H.H. Chen and B.W. Lee - University of California at Irvive,Preprint.

12. S. Weinberg - Phys. Rev. D 5, 1412 (1972).

13. The predicted range for this ratio given in Weinberg's paper, 0.15 to 0.25, (ref. 12) is only approximate. I would like to thank G. Myatt for pointing this out to me.

14. C. Franzinetti - Proceedings of the Topical Conference on Weak Interactions - CERN 1969.

15. B.W. Lee - National Accelerator Laboratory - Preprint, May 1972.

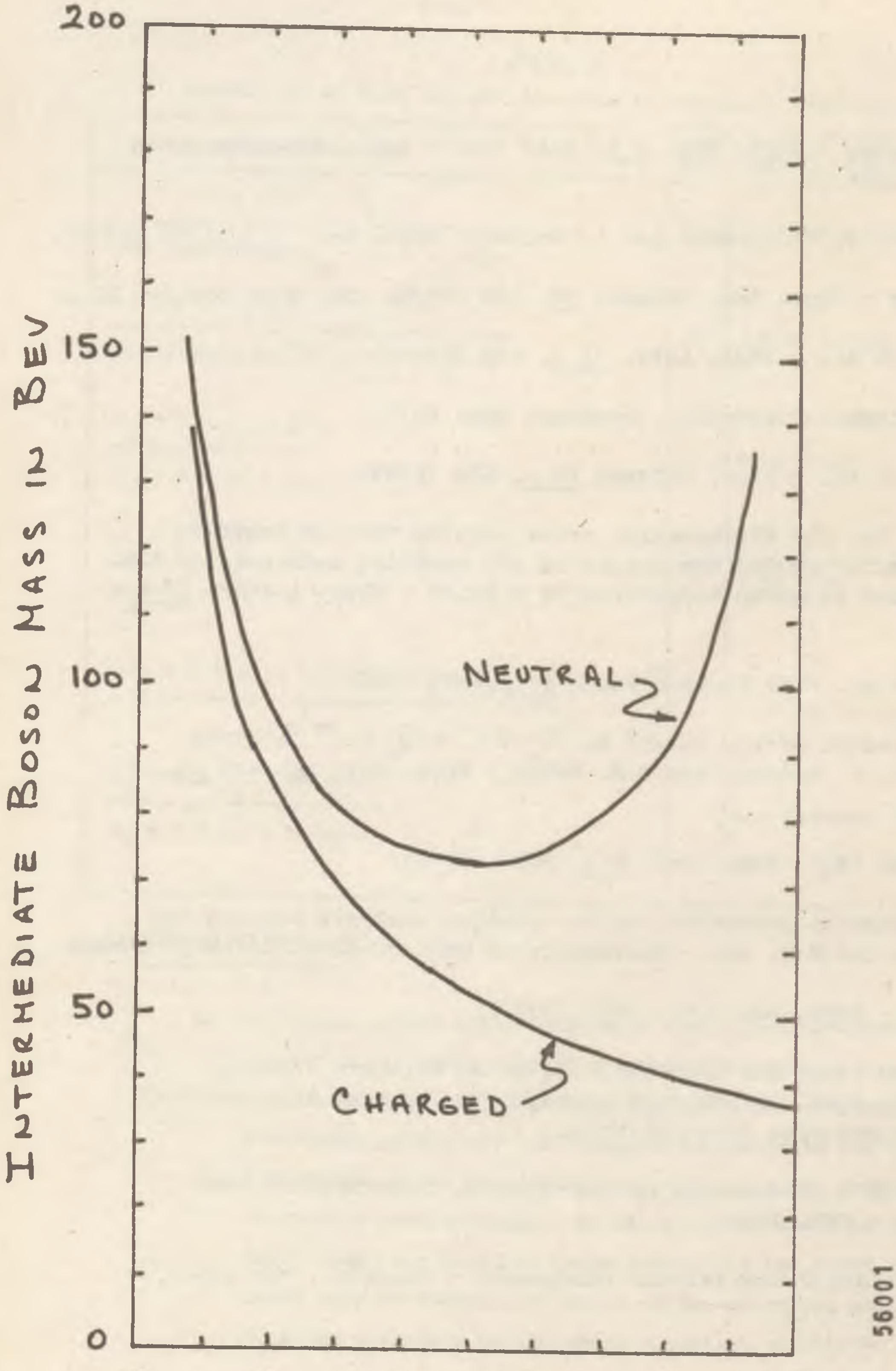
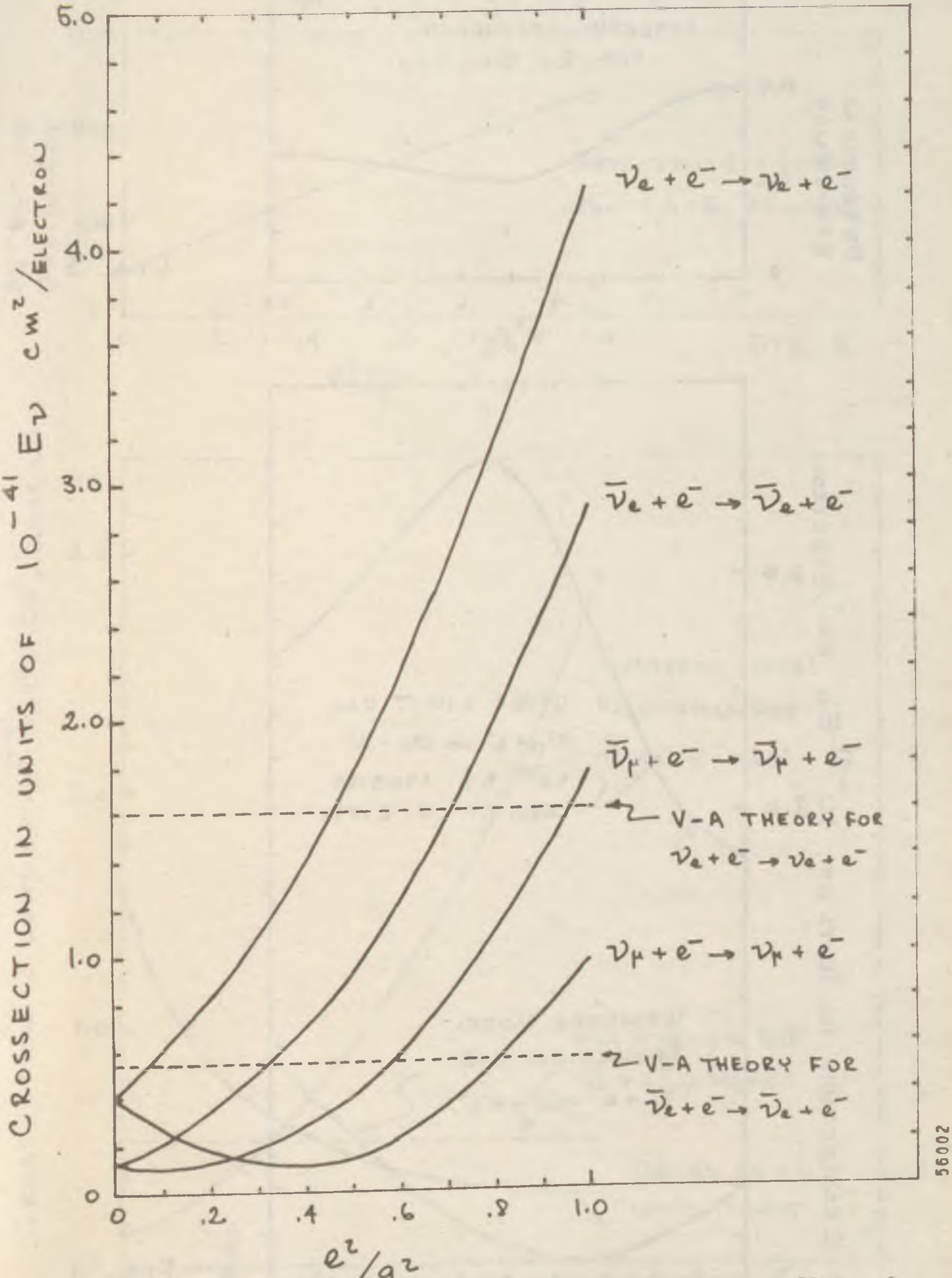

200
150
100
50
0
Intermediate Boson Mass in BeV
Neutral
Charged
0
.2
.4
.6
.8
1.0
e^2/g^2
56001

FIG.



5.0
4.0
3.0
2.0
1.0
0
0
.2
.4
.6
.8
1.0
CROSSECTION IN UNITS OF 10^{-41} E_ν cm^2/ELECTRON
e^2/g^2
$\nu_e + e^- \to \nu_e + e^-$
$\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-$
$\bar{\nu}_\mu + e^- \to \bar{\nu}_\mu + e^-$
V-A THEORY FOR $\nu_e + e^- \to \nu_e + e^-$
$\nu_\mu + e^- \to \nu_\mu + e^-$
V-A THEORY FOR $\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-$
56002

FIG. 2.

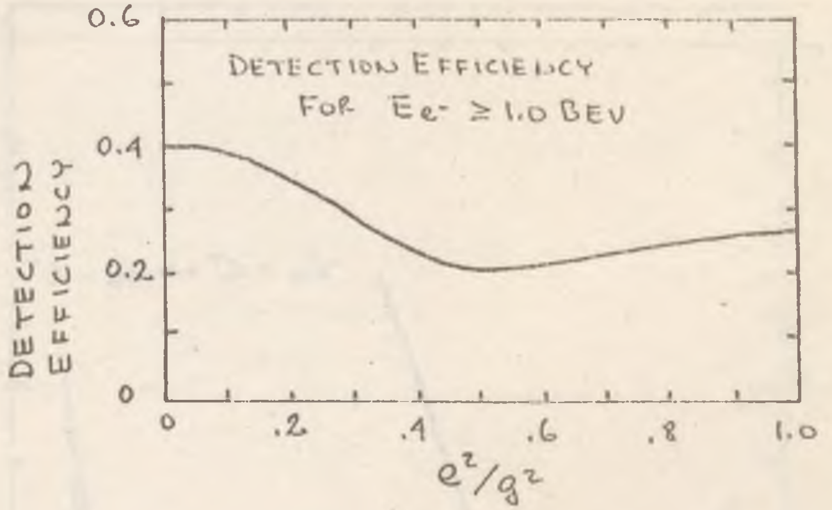

DETECTION EFFICIENCY
FOR Ee- ≥ 1.0 BEV
DETECTION EFFICIENCY
0.6
0.4
0.2
0
0
.2
.4
.6
.8
1.0
e^2/g^2

FIG. 3

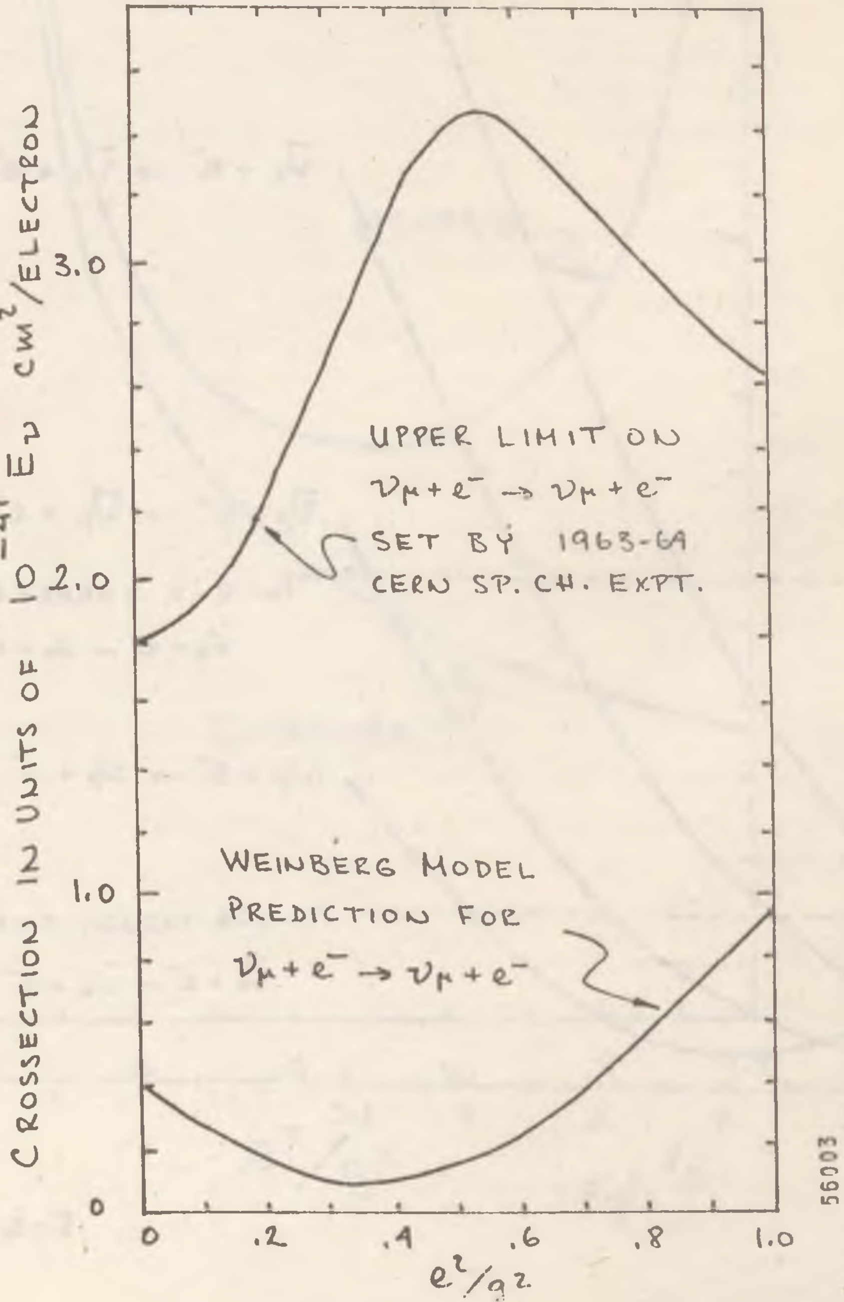

CROSSECTION IN UNITS OF 10^{-41} E_ν cm^2/ELECTRON
3.0
2.0
1.0
0
UPPER LIMIT ON
$\nu_\mu + e^- \rightarrow \nu_\mu + e^-$
SET BY 1963-64
CERN SP. CH. EXPT.
WEINBERG MODEL
PREDICTION FOR
$\nu_\mu + e^- \rightarrow \nu_\mu + e^-$
0
.2
.4
.6
.8
1.0
e^2/g^2
56003

FIG. 4

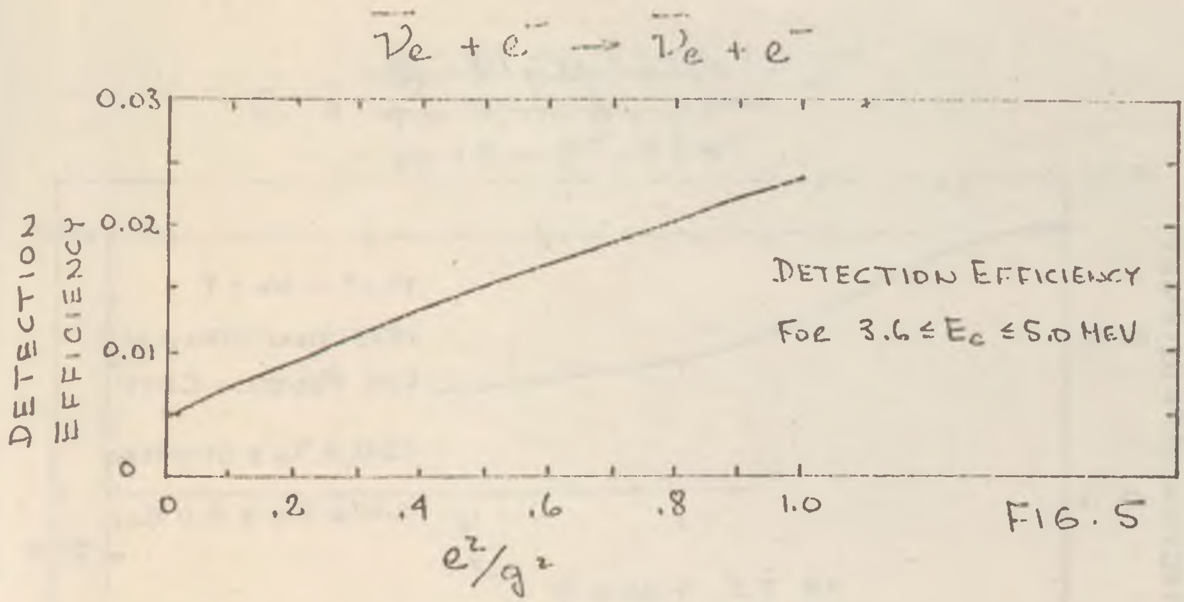

$\bar{\nu}_e + e^- \rightarrow \bar{\nu}_e + e^-$
0.03
0.02
0.01
0
DETECTION EFFICIENCY
DETECTION EFFICIENCY FOR $3.6 \leq E_e \leq 5.0$ MEV
0
.2
.4
.6
.8
1.0
e^2/g^2

FIG. 5

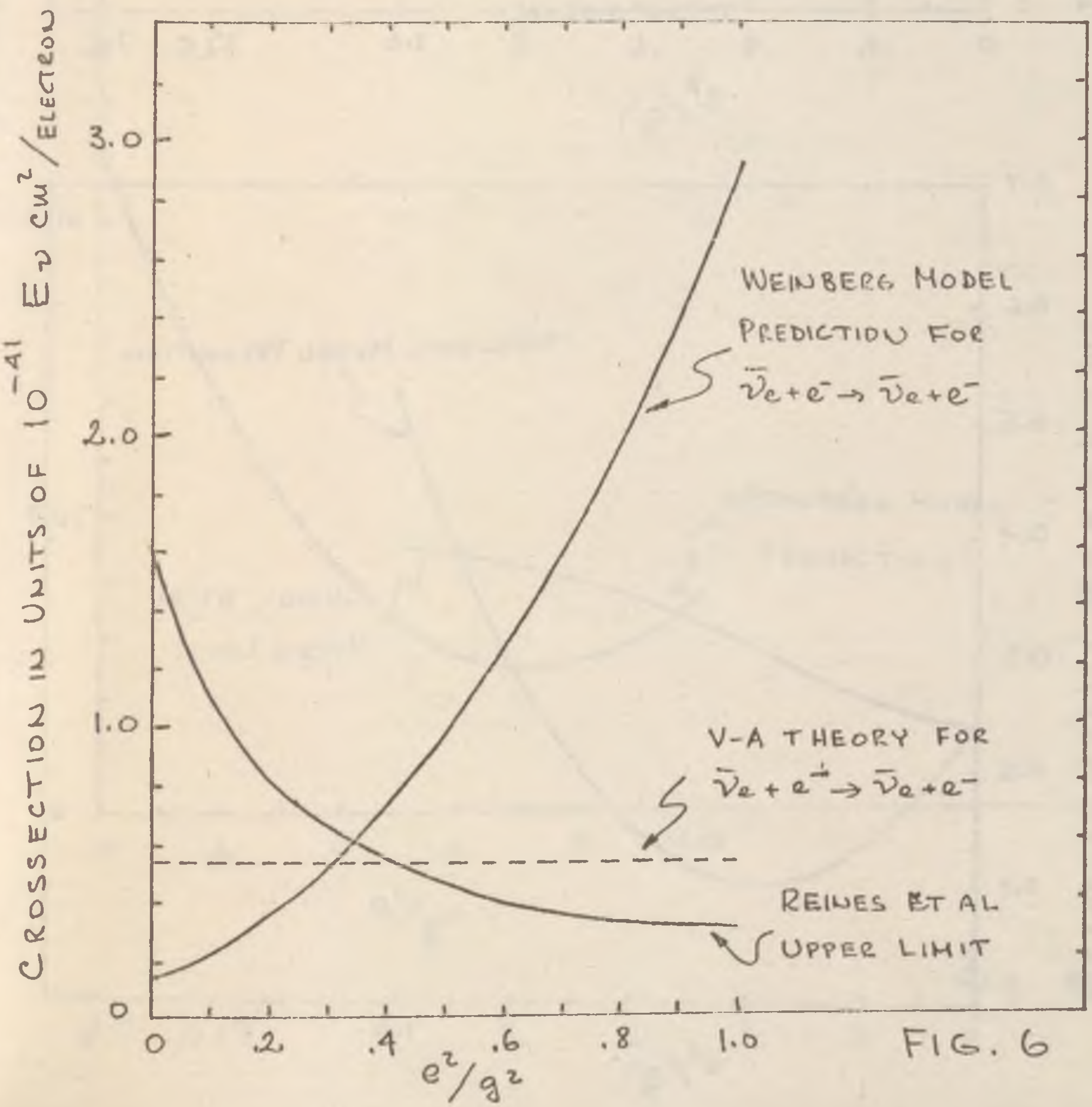

CROSSECTION IN UNITS OF 10^{-41} E_ν cm^2/ELECTRON
3.0
2.0
1.0
0
WEINBERG MODEL PREDICTION FOR $\bar{\nu}_e + e^- \rightarrow \bar{\nu}_e + e^-$
V-A THEORY FOR $\bar{\nu}_e + e^- \rightarrow \bar{\nu}_e + e^-$
REINES ET AL UPPER LIMIT
0
.2
.4
.6
.8
1.0
e^2/g^2

FIG. 6

$$R_\nu = \frac{\nu_\mu + p \rightarrow \nu_\mu + p}{\nu_\mu + n \rightarrow \mu^- + p}$$

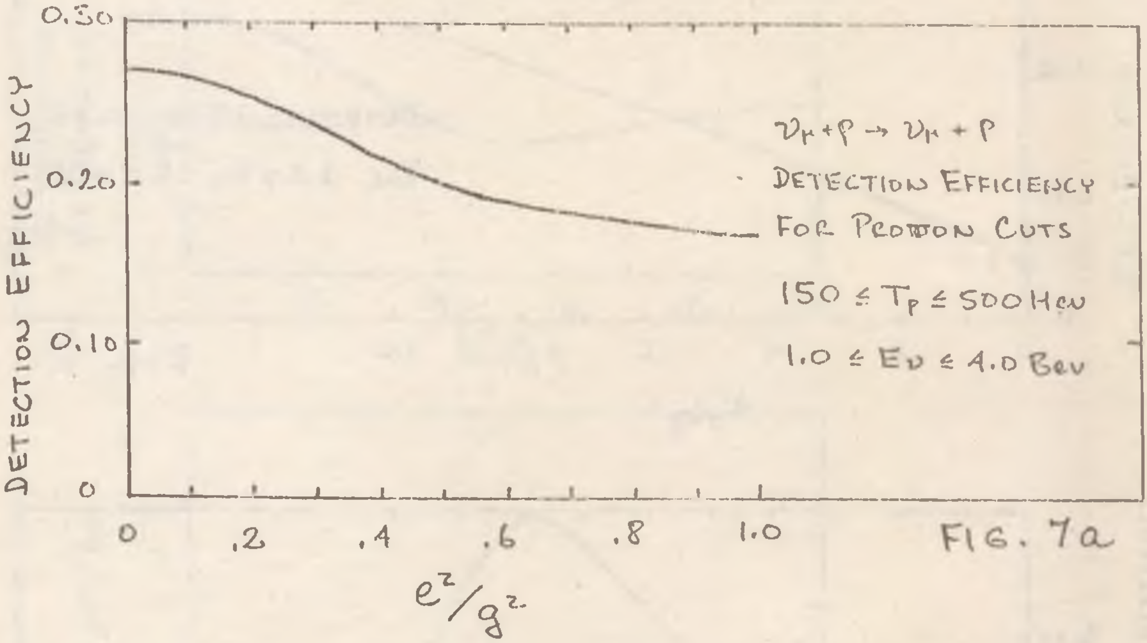


FIG. 7a

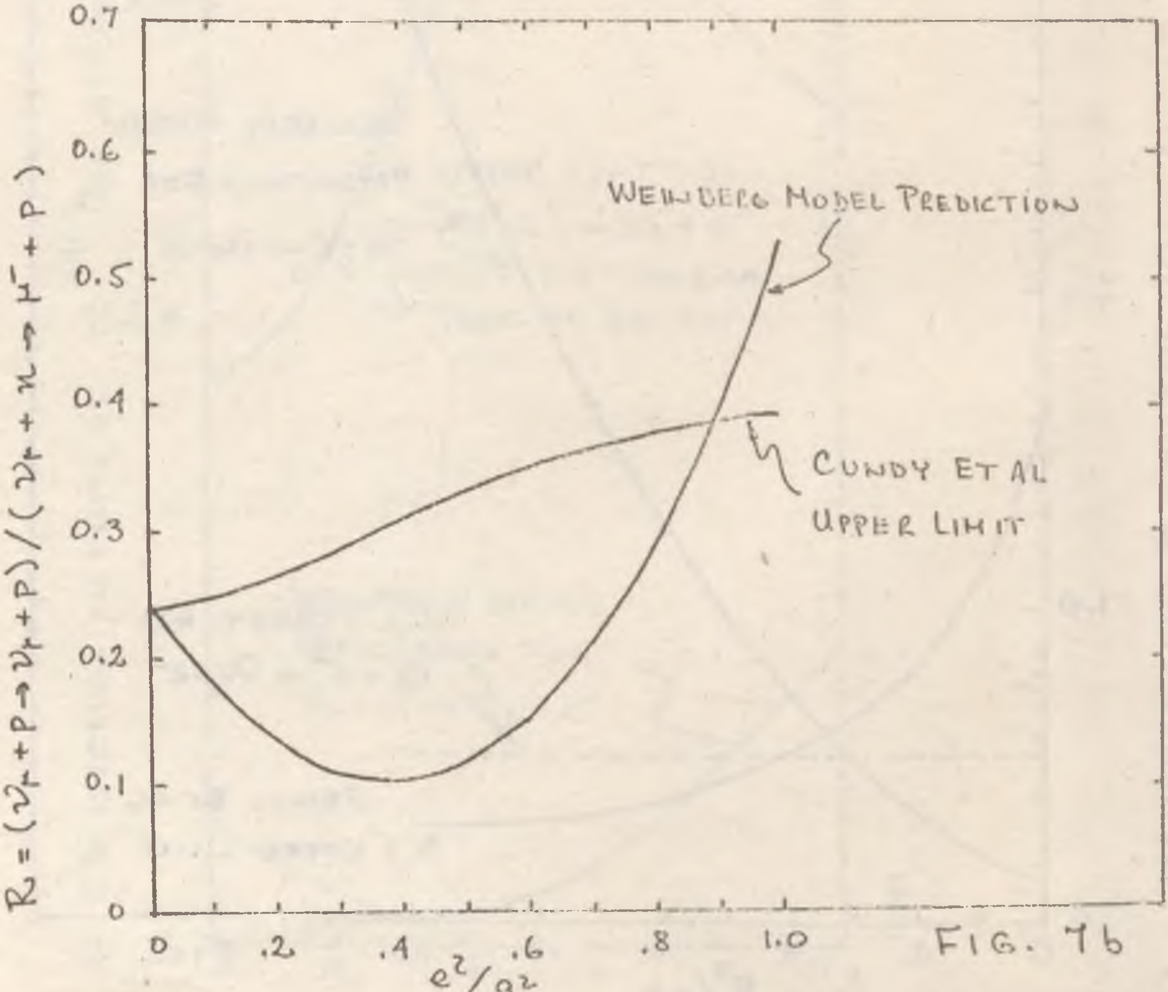


FIG. 7b

$$R = \frac{\nu_\mu + p \to \nu_\mu + n + \pi^+}{\nu_\mu + p \to \mu^- + p + \pi^+}$$

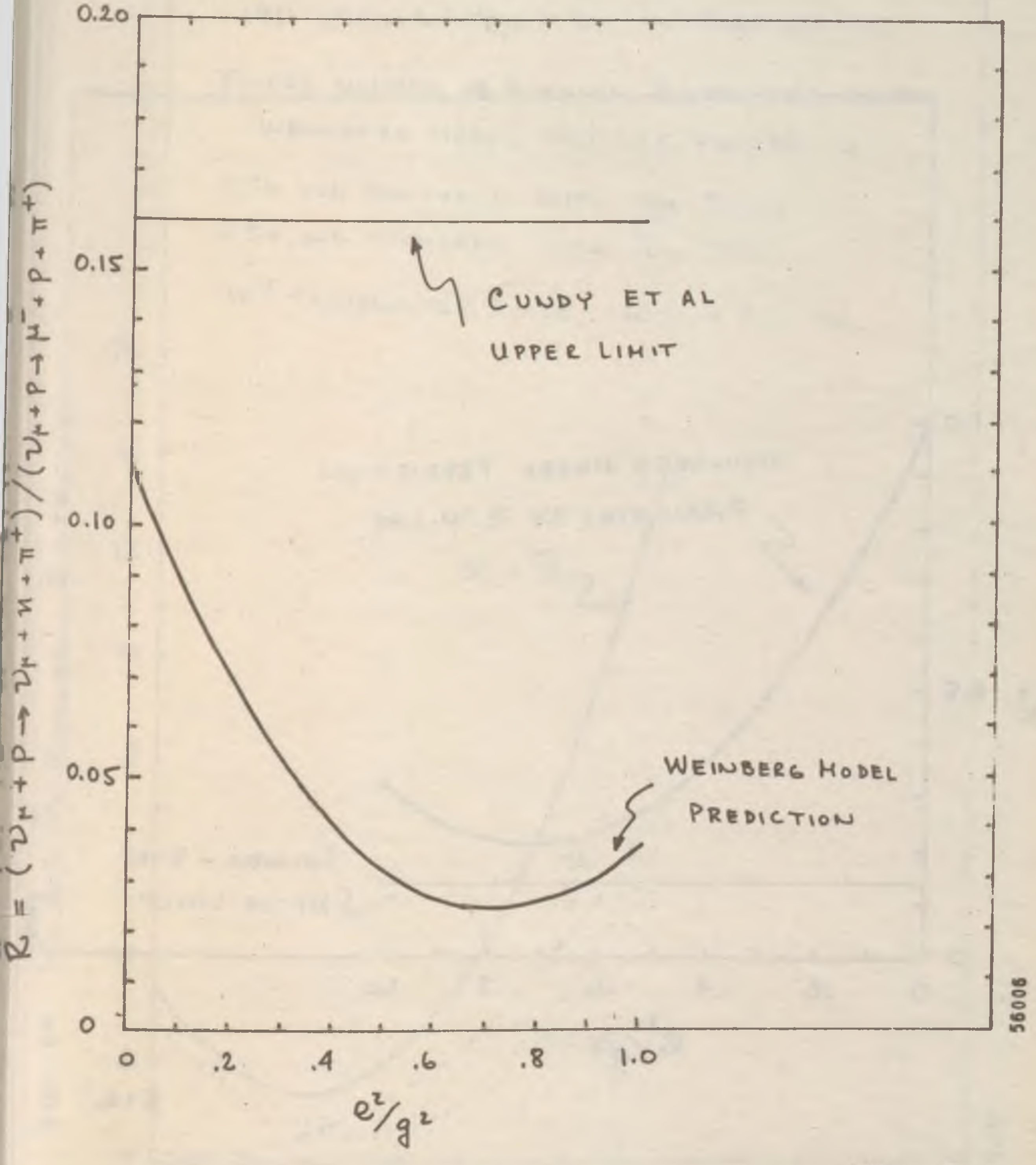


FIG. 8

$$R = \frac{(\nu_\mu + p \to \nu_\mu + p + \pi^0) + (\nu_\mu + n \to \nu_\mu + n + \pi^0)}{2\,(\nu_\mu + n \to \mu^- + p + \pi^0)}$$

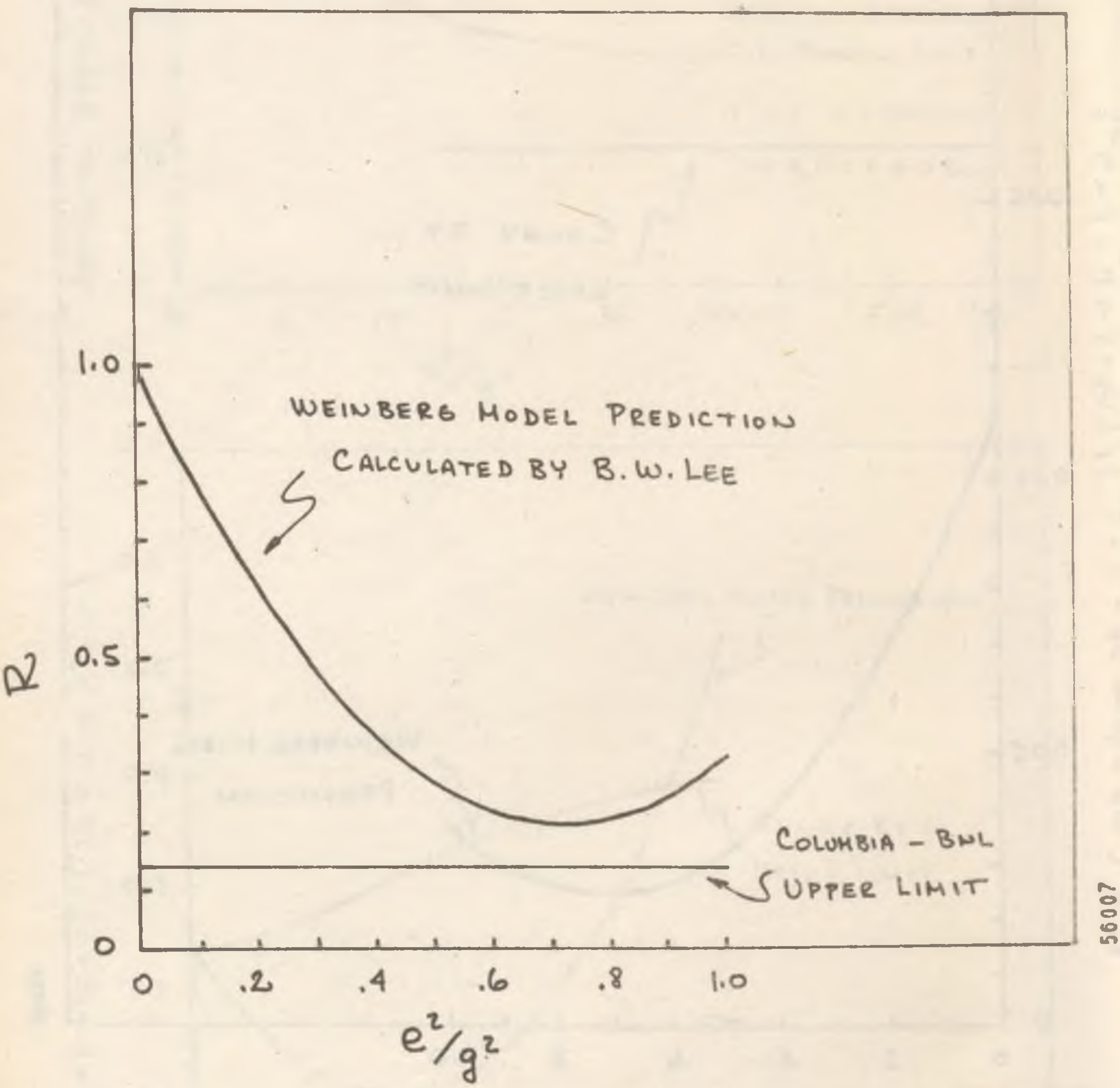


FIG. 9

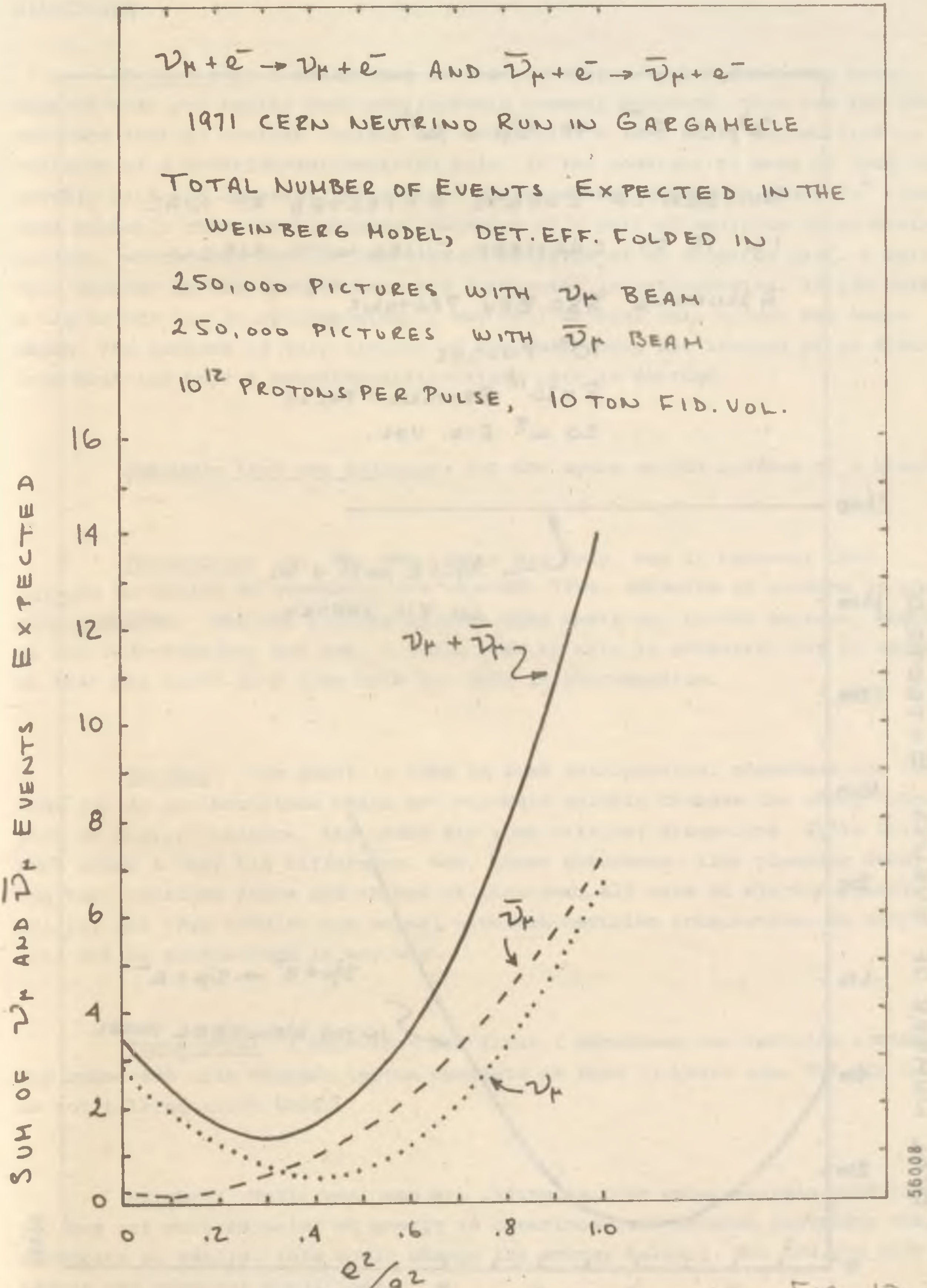
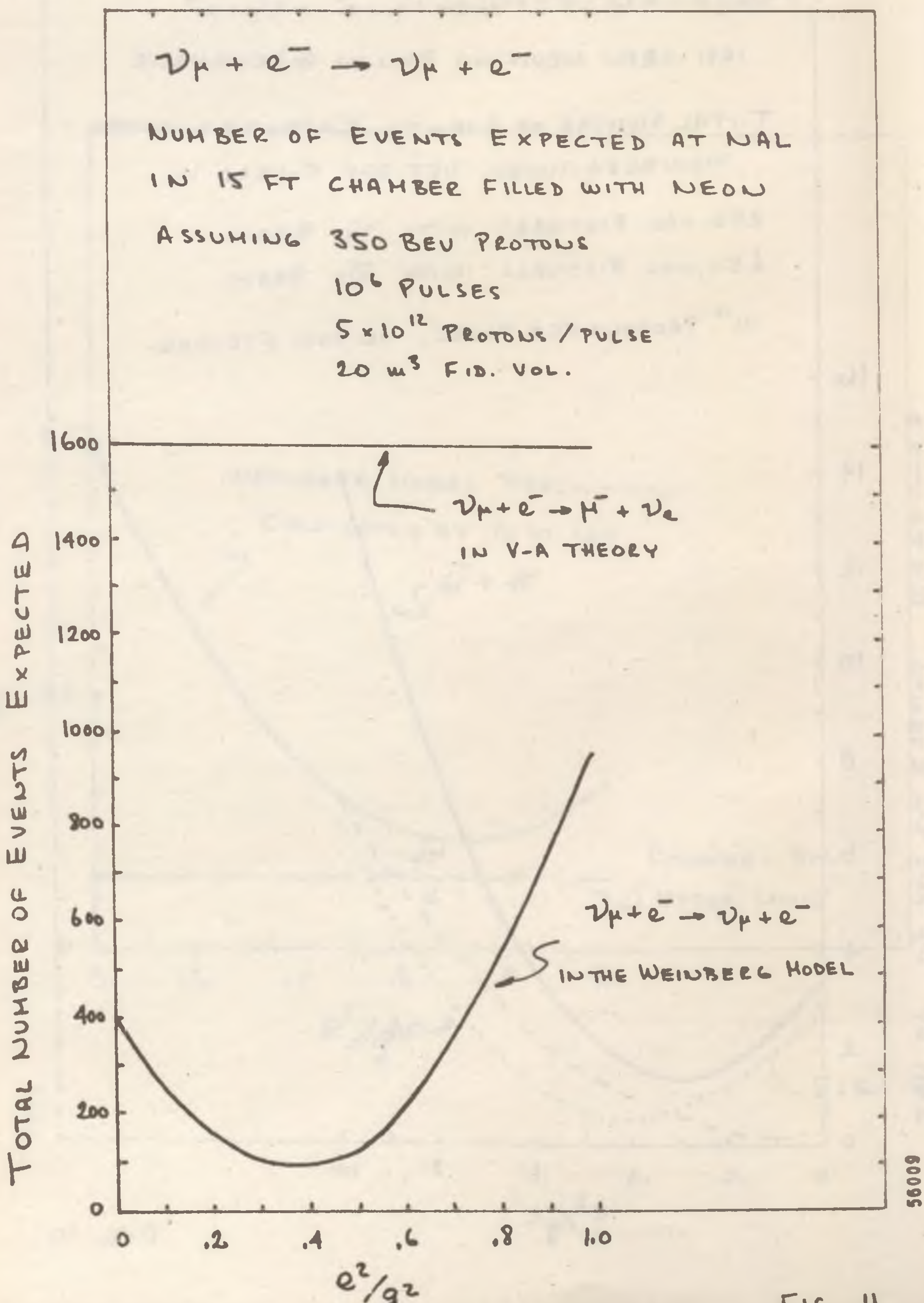

$\nu_\mu + e^- \to \nu_\mu + e^-$ AND $\bar{\nu}_\mu + e^- \to \bar{\nu}_\mu + e^-$
1971 CERN NEUTRINO RUN IN GARGAMELLE
TOTAL NUMBER OF EVENTS EXPECTED IN THE WEINBERG MODEL, DET. EFF. FOLDED IN
250,000 PICTURES WITH ν_μ BEAM
250,000 PICTURES WITH $\bar{\nu}_\mu$ BEAM
10^{12} PROTONS PER PULSE, 10 TON FID. VOL.
SUM OF ν_μ AND $\bar{\nu}_\mu$ EVENTS EXPECTED
$\nu_\mu + \bar{\nu}_\mu$
$\bar{\nu}_\mu$
ν_μ
e^2/g^2
0
2
4
6
8
10
12
14
16
.2
.4
.6
.8
1.0
56008

FIG. 10



$\nu_\mu + e^- \rightarrow \nu_\mu + e^-$
NUMBER OF EVENTS EXPECTED AT NAL
IN 15 FT CHAMBER FILLED WITH NEON
ASSUMING 350 BEV PROTONS
10^6 PULSES
5×10^{12} PROTONS / PULSE
20 m^3 FID. VOL.
$\nu_\mu + e^- \rightarrow \mu^- + \nu_e$
IN V-A THEORY
$\nu_\mu + e^- \rightarrow \nu_\mu + e^-$
IN THE WEINBERG MODEL
TOTAL NUMBER OF EVENTS EXPECTED
1600
1400
1200
1000
800
600
400
200
0
0
.2
.4
.6
.8
1.0
e^2/g^2
56009

FIG. 11

DISCUSSION

Pontecorvo: I would like to draw attention to the following point. Suppose that you really have semileptonic neutral currents. This has the consequence that an excited nucleus may decay with a very small probability by emission of a neutrino-antineutrino pair. If the constant is more or less comparable with the constant for charged currents this effect is about 10^4 times more probable then the "ordinary" emission of a pair of neutrino by an excited nucleus, which goes through the virtual emission of an electron pair. I said this because so many people here are interested in astrophysics. If you have a big object as in astrophysics, it may well be that this effect has importance. The process is very similar to the beta decay but instead of an electron-neutrino pair a neutrino-antineutrino pair is emitted.

Somebody from the audience: For the atoms on the surface of a star?

Pontecorvo: No. No. You have a big body, and it happens, that baryons or nuclei or something are excited. True, emission of photons is much more probable. But the photons have to find their way to the surface, whereas the neutrinos just get out. I don't know if this is relevant, but it might be that one would find some role for this in astrophysics.

Marshak: The point is that in some astrophysical phenomena the fact that you do get neutrinos which get out very quickly changes the whole transport of energy balance. And there are some critical situations where this fact makes a very big difference. Now, those phenomena, like plasmons decaying into neutrino pairs and things of that sort, all have an electromagnetic origin; and they involve the normal electron-neutrino interaction. So they will not be accelerated in any way.

Pontecorvo: I know that and first I mentioned the neutrino luminosity connected with charged lepton currents at Kiev 13 years ago. But now I am not talking about that.

Marshak: Well, what you are saying is that under certain conditions you may get more emission of energy in neutrinos, and because neutrinos can penetrate so easily, this could change the energy balance. But now you have to invent the physical conditions.

Pontecorvo: I am not able to invent them, and what I want to say is that may be there will be collisions in which neutrinos are emitted directly (without electromagnetic interactions) although I do not know under which conditions. I am not going to discuss now the effect of pure leptonic neutral currents. This might change somewhat the picture of the "ordinary" neutrino luminosity (through electromagnetic + Fermi lepton interaction). In particular pairs of muon neutrinos (in addition to electron neutrinos) will be emitted, a fact which was stressed here by Radicati.

SEARCH FOR NEUTRAL CURRENTS IN "GARGAMELLE"

ν collaboration: Aachen, Bruxelles, CERN, Ecole Polytechnique, Milano
Orsay, University College of London
/presented by A.Pullia/

I have to report about search for neutral currents in "Gargamelle", the large heavy liquid bubble chamber exposed to the ν and $\bar{\nu}$ beams at CERN during 71.

The experiment is in a very preliminary stage; therefore I can report only about the first qualitative indications and about future possibilities. I will first briefly summarize the parameters of the beam(Fig. I); the calculated flux of ν and ν are reported in fig. 2. The chamber was filled with freon (CF_3Br).

Concerning neutral currents I divide my talk in 2 parts:

I) leptonic neutral currents

2) semileptonic neutral currents

<u>LEPTONIC NEUTRAL CURRENTS /</u>

we are searching for events composed by I electronic track and without hadrons, both in ν and in $\bar{\nu}$ films.

Our goal is to test the theory of weak interactions of Weinberg(I) who foresees the existence of neutral intermediate bosons W_0 and hence the existence of reactions of the type:

$$(\mathrm{I})\quad \begin{cases} \nu_\mu + e^- \rightarrow \nu_\mu + e^- \\ \bar{\nu}_\mu + e^- \rightarrow \bar{\nu}_\mu + e^- \end{cases}$$

G. t' Hooft (2) worked out calculations of cross sections for reactions (I) as a function of the parameter e^2/g^2

(e= electromagnetic coupling constant; g=coupling constant of the isotriplet intermediate bosons to the lepton current). e^2/g^2 can range between 0 and I.
If we adapt these calculations to our experimental situation we obtain the number of expected events shown in fig. 3. To reduce background due to "Compton" electrons associated to γ rays incoming in the chamber, we search only for high energy electrons (> 300 Mev) and in the direction of the beam ($\vartheta < 5°$).
In this condition (after scanning of I/2 of the films) no candidate was found; we cannot make quantitative com= parison with the theory now, for our scanning efficiency must be checked by a particular rescanning (in progress); we can however conclude that we have a qualitative indi= cation against high values of e^2/g^2.

SEMILEPTONIC NEUTRAL CURRENTS

The experimental study of semileptonic neutral currents is strictly connected with the problem of backround from neutrons incoming in the detector.
Possible source of neutrons are:
I)Interactions of cosmic rays near the chamber
2)Hadronic cascade (p,π,K) originated in the μ shielding
3)Interactions of ν ($\bar{\nu}$) in the shielding
4)Interactions of ν and $\bar{\nu}$ in the iron of the magnet around the chamber.

There are however many advantages by using a large bubble chamber filled with an heavy liquid. It is possible indeed, to observe the exponential decrease (of a neutron beam) along the beam axis (case 2,3) ; Gargamelle in fact is long 4.8 m (corresponding to 7-8 interactions lengths of neutrons in freon) and its radius is 90 cm. (about I,5 interaction lenghts). We can consider for instance the axial distribution of the interaction points of events composed by a short recoil proton; they are due to the flux of low energy neutrons from the shielding of μ. This distribution is well fitted by an exponential curve (fig. 4). There are other advantages , not strictly connected with the neutron beckround , but connected with the problem of neutral currents:

a)Clear distinction of pions from muons; the potential path of produced mesons in the chamber is about 2 interaction lenghts of a

b)High γ rays detection efficency; this fact allows to study cleanly, possible channels with π^0 ; for instance:

$$(2)\quad \begin{cases} \overset{\nu}{\bar{\nu}} + n \longrightarrow N^{*0} + \overset{\nu}{\bar{\nu}} \\ \qquad\qquad\quad \hookrightarrow \pi^0 + n \\ \overset{\nu}{\bar{\nu}} + p \longrightarrow N^{*+} + \overset{\nu}{\bar{\nu}} \\ \qquad\qquad\quad \hookrightarrow \pi^0 + p \end{cases}$$

Events of type 2 are very clean because they are composed only by protons and γ rays and the possibility of confusion with muonic events does not exist.

Very preliminary results:

I) $\bar{\nu}$ films : we started with the analysis of events with at least I charged or neutral π associated. They are recognized by the presence of γ rays or charged interacting particles. The rate of these events as compared with the corresponding muonic events is $\simeq$40%.

Let me consider only the channel with π^0 production.

In fig. 5 is shown the diplot of x (coordinate of the interaction point along the beam axis) versus P_r (Component of the momentum along the radial axis). We can see:

I) General decrease of density of interaction points along the beam axis (neutron flux incoming from the entrance surface of the chamber)

2)Strong asymmetry around $P_r=0$. Events with $P_r<0$ are dominant; this fact indicates the presence of neutron flux from peripheral parts of the chamber (from the magnet).

We can conclude that we have no evidence for presence of neutral currents , at least for this channel.

II) ν films: in this case the situation is more profitable. The rate of hadronic events found in the scanning (with at least 1π associated) is indeed 12% as compared with the corresponding muonic events. Unfortunately the measurements of these events are not yet ready and it is impossible to reach to day any conclusion about them.

R E F E R E N C E

I) S.Weinberg P.R.L. I9,I264 (67)

2) G. t'Hooft P.L. 37B,I95 (7I)

ν beam parameter

Proton beam intensity	$\sim 1.2 \; 10^{12}$
Decay length	70 m.
Shielding (iron)	27 m.
Horn 1 current	337 KA
Horn 2 current	400 KA

	ν	$\bar{\nu}$
n° of pictures	$\simeq$ 250 K	$\simeq$ 250 K

Fig. 1

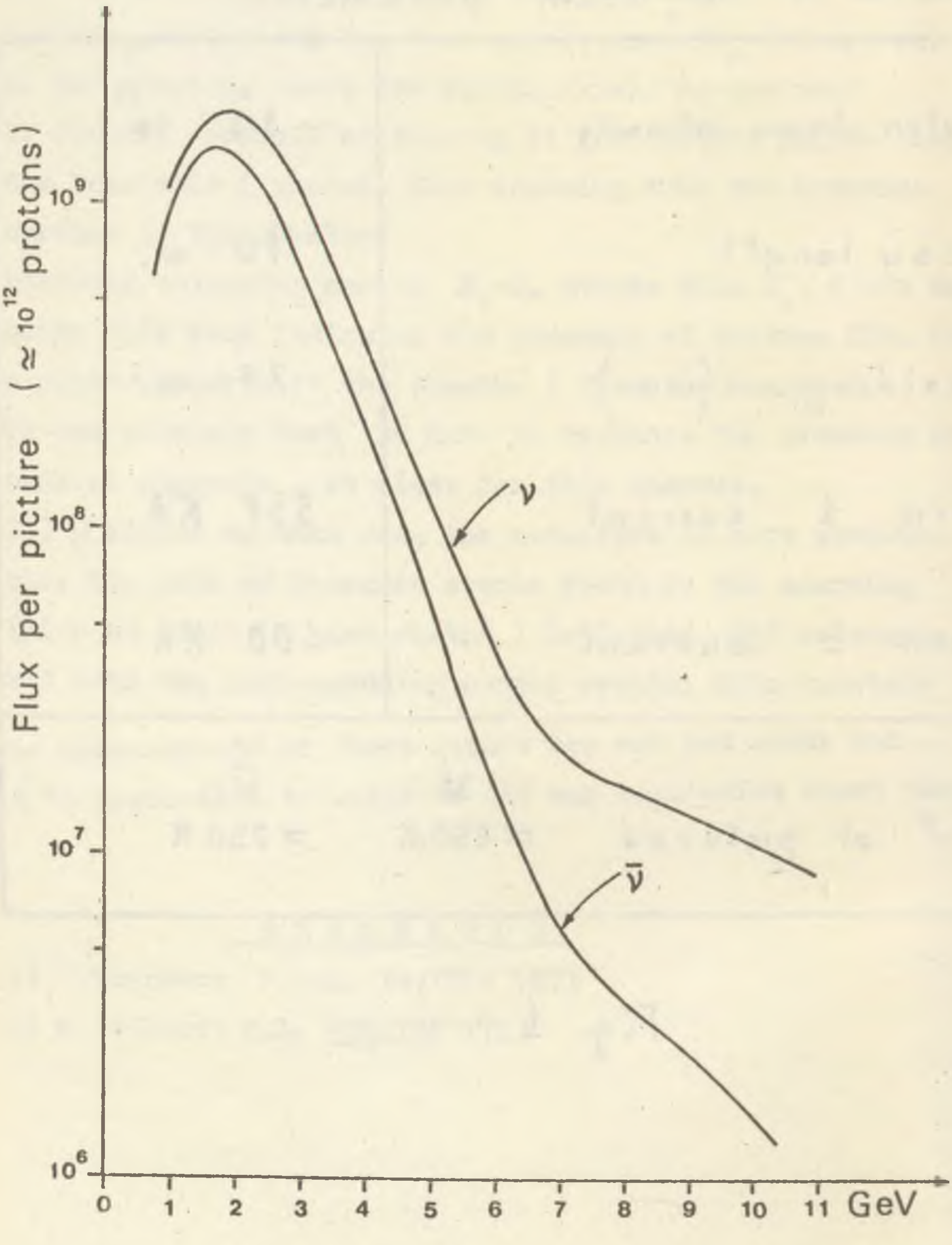

Flux per picture ($\simeq 10^{12}$ protons)
10^9
10^8
10^7
10^6
0 1 2 3 4 5 6 7 8 9 10 11 GeV
ν
$\bar{\nu}$

Fig. 2

$\nu_\mu + \bar{\nu}_\mu$ events for 250.000 pictures of each

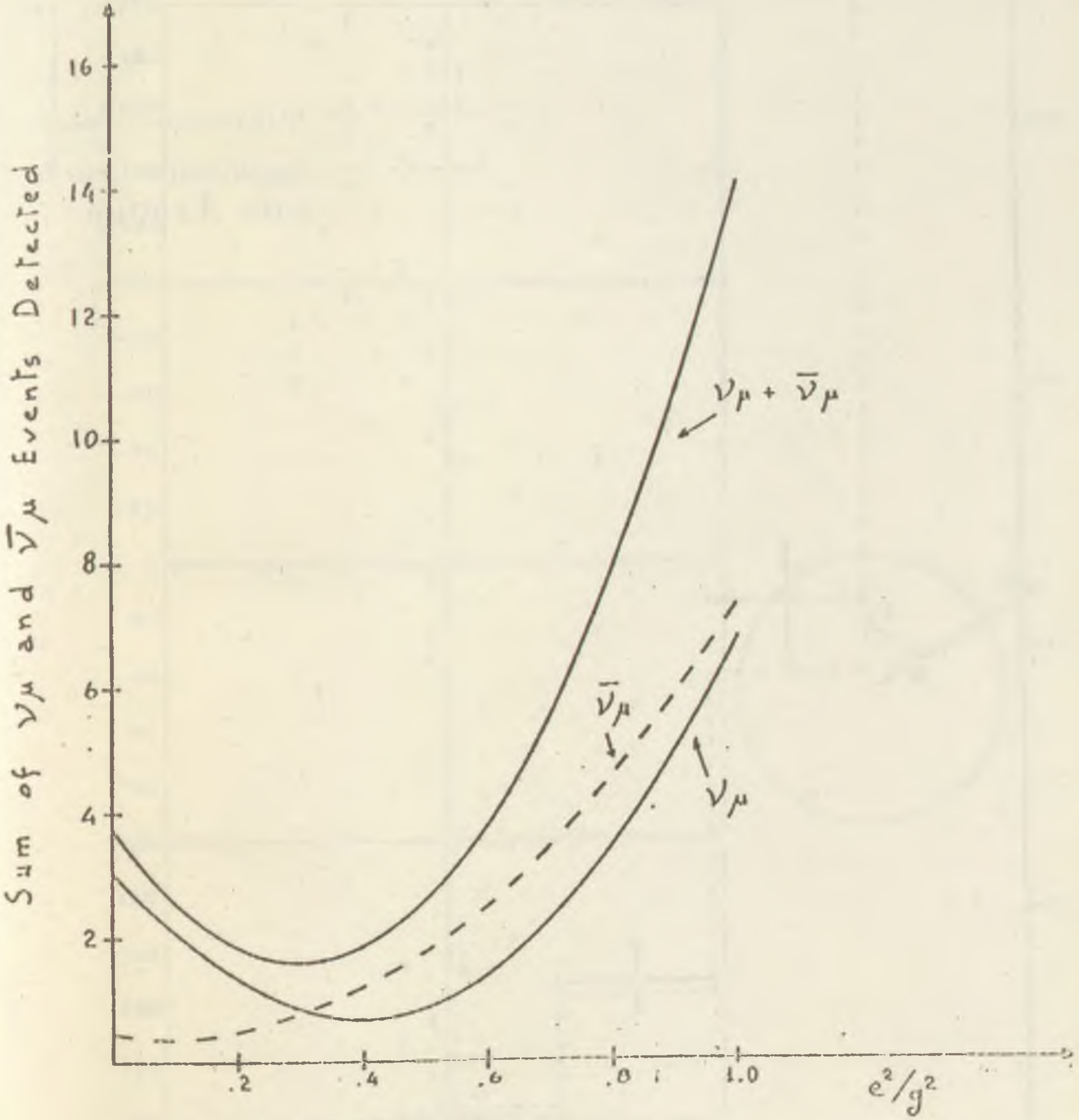
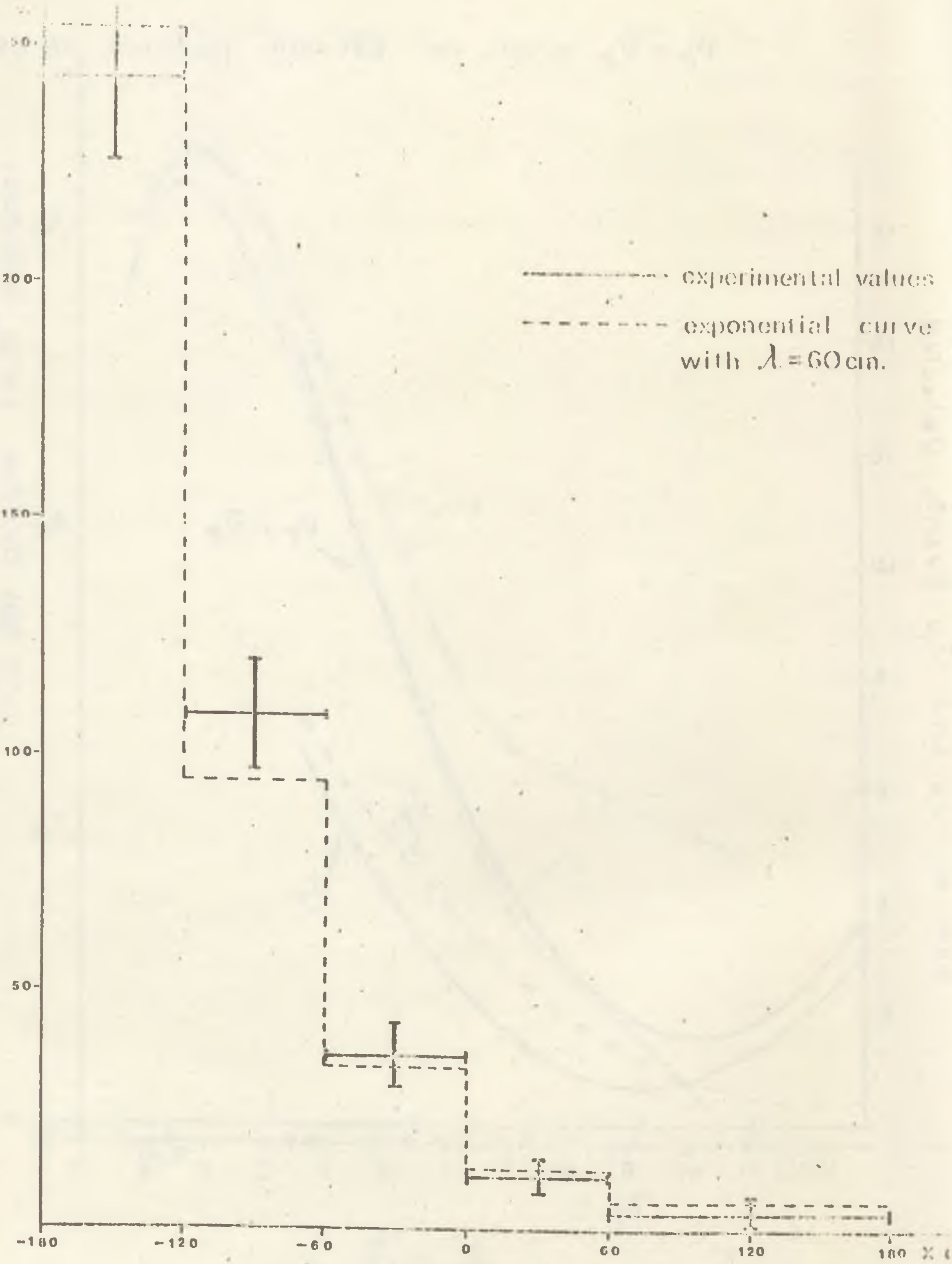


Fig. 3



250
200
150
100
50
-180
-120
-60
0
60
120
180
X
experimental values
exponential curve with λ = 60 cm.

Fig. 4

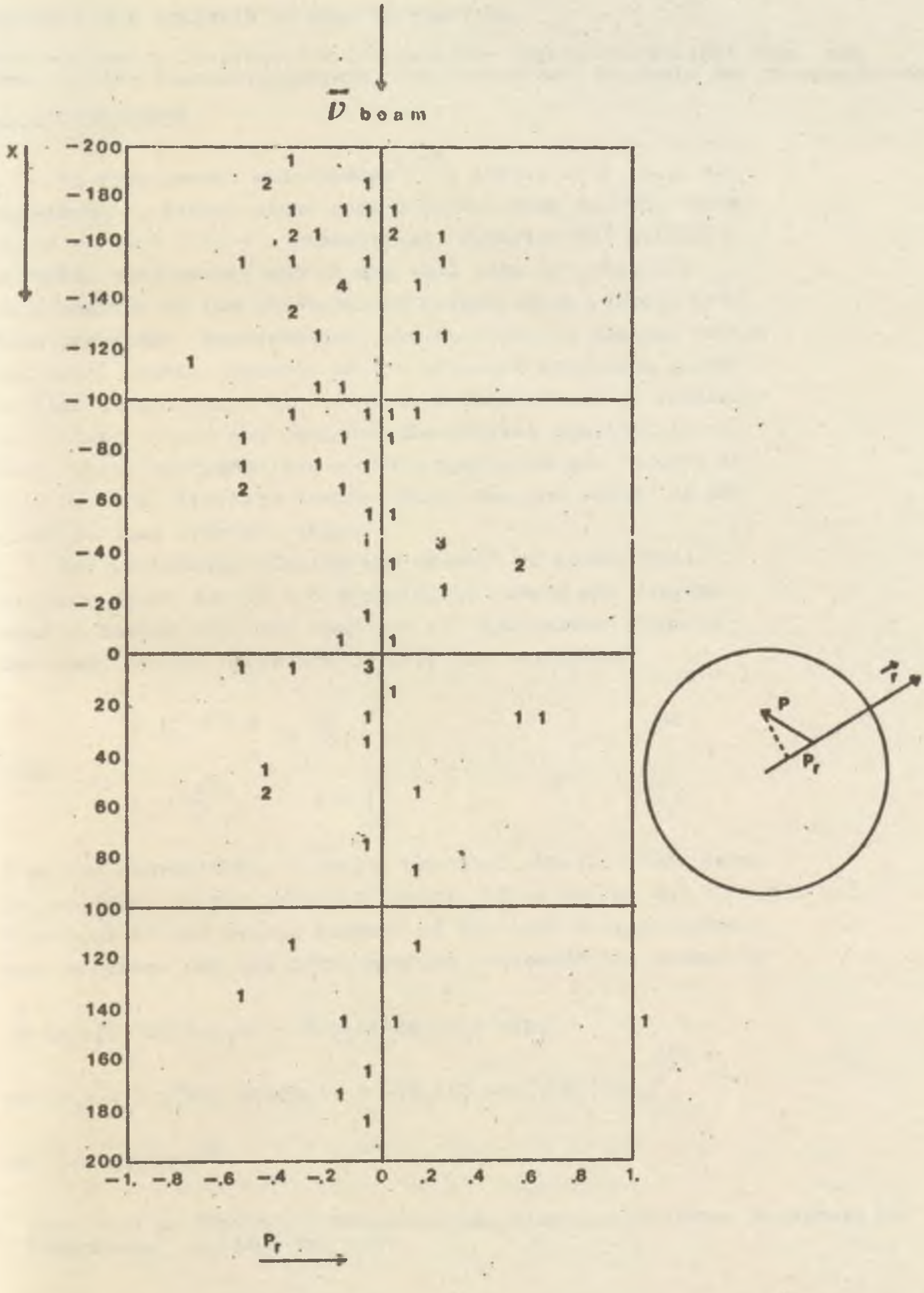

$\bar{\nu}$ beam
X
-200
-180
-160
-140
-120
-100
-80
-60
-40
-20
0
20
40
60
80
100
120
140
160
180
200
-1. -.8 -.6 -.4 -.2 0 .2 .4 .6 .8 1.
P_r
P
P_r
r

Fig. 5

SECOND CLASS CURRENTS IN WEAK INTERACTIONS +

H.Pietschmann, Institut für Theoretische Physik, Universität Wien and Institut für Hochenergiephysik, Österreichische Akademie der Wissenschaften

1. Introduction

In some recent experiments[1,2], indications as to the existence of second class currents have been detected both in $\Delta Y = 0$ and $|\Delta Y| = 1$ transitions. Experimental evidence is still very scanty and it may well turn out that the announcement of the existence of second class effects had been premature. Nevertheless, the question of whether vector and axial vector currents of the standard structure alone suffice to construct the total weak Hamiltonian is intrinsically interesting and deserves theoretical considerations. Thus, three contributions to this symposium are devoted to this subject; they are complementary and the reader is advised to look into all three.

Let us briefly clarify the concept of second class currents first for $\Delta Y = 0$ transitions, there are various ways to define them but they are all equivalent. Suppose the weak current operators satisfy the relations

$$G\, \overline{j^i_\lambda}\, G^{-1} = \varepsilon_i\, j^i_\lambda \tag{1}$$

with

$$\varepsilon_i = \begin{cases} +1 \\ -1 \end{cases} \qquad i = \begin{cases} V \\ A \end{cases} \tag{2}$$

G is the conventional G-parity operator. Eq.(1) holds true, for example, in the pure V-A theory. If we define the vertex functions of the matrix element of the weak current between baryon states and the corresponding G-transformed states by

$$\begin{aligned} \langle B'(p_2)|j^i(0)|B(p_1)\rangle &= \bar{u}(p_2)\Gamma^i_\lambda(p_1,p_2)\, u(p_1) \\ \langle B'(p_2)|G^{-1}\, j^i_\lambda(0)\, G|B(p_1)\rangle &= \bar{v}(p_1)\Gamma^{i'}_\lambda(-p_2,-p_1)v(p_2) \end{aligned} \tag{3}$$

it then follows[3)]

+ Supported by "Fonds zur Förderung der wissenschaftlichen Forschung in Österreich", Project Nr. 1497.

$$\Gamma_\lambda^i(p_1,p_2) = -\varepsilon_i \; C \; \Gamma_\lambda^{i'}(-p_2,-p_1)C^{-1} \tag{4}$$

where C is the 4×4 charge conjugation matrix with the properties

$$C^{-1}\gamma_\mu C = -\gamma_\mu^T$$

$$\bar{u}\, C^T = v^T \tag{5}$$

$$C^T = -C$$

For the case of transitions within isomultipletts, the primed and unprimed vertex functions of eqs.(3) coincide and we obtain a restriction on the vertex function. For transitions between isomultipletts, Eq.(4) connects G-parity related transitions. The standard example for the first case is n-decay, where Eq.(4) requires the absence of F_3 and G_3 in the decomposition of the vertex function

$$\Gamma_\lambda(p_1,p_2) = F_1(q^2)\gamma_\lambda + F_2(q^2)i\sigma_{\lambda\nu}q^\nu + F_3(q^2)q_\lambda +$$

$$+ G_1(q^2)\gamma_\lambda\gamma_5 + G_2(q^2)q_\lambda\gamma_5 + G_3(q^2)i\sigma_{\lambda\nu}q^\nu\gamma_5 \tag{6}$$

with

$$q = p_2 - p_1 \tag{7}$$

The standard example for the second case is $\Sigma^\pm\Lambda$ decay, where Eq.(4) gives the equality for the two relevant vertex functions. Hence the ratio

$$R_{\Sigma\Lambda} = \frac{\Gamma(\Sigma^+\to\Lambda e^+\bar{\nu}_e)}{\Gamma(\Sigma^-\to\Lambda e^-\nu_e)} \tag{8}$$

is given by phase space only, i.e.

$$R^{o}_{\Sigma\Lambda} = \left(\frac{m_{\Sigma^+} - m_\Lambda}{m_{\Sigma^-} - m_\Lambda}\right)^5 = 0,6 \tag{9}$$

We shall assume that G-parity is conserved in strong interactions. Therefore, any experimental detection of form factors violating Eq.(4) requires additional current operators which do not obey Eq.(1). These are called second class currents.

Unfortunately, experiments can never give us direct information on the structure of these second class current operators. We only obtain the vertex function (6). Therefore, the method of trial and error has to be applied. We have to write down all possible forms of second class current operators, work out all their experimental consequences and in this way try to obtain upper bounds for them. It is the hope, that eventually most of them or even all of them can be ruled out. It is very clear, that we are only on our way to this goal; much work is still necessary before we can rest and contemplate final results.

2. The Induced Pseudotensor Current of Second Class.

The straightforward way to turn a second class form factor into a current operator is to write the following quark current[4)]

$$j^{PT}_\nu = i\partial^\mu(\bar{q}\gamma_5 i\sigma_{\nu\mu}\lambda q) \tag{10}$$

where λ is a 3×3 SU_3 matrix which selects the wanted transformation properties. It has first been pointed out by P. Hertel[5)], that any second class current operator can give contributions to all three form factors of its parity (vector or axial vector). It turns out, that second class contributions to axial vector and pseudoscalar form factors G_1 and G_2 vanish in the SU_3 limit. (See also R. Oehme, ref.4).

A pseudotensor form factor contribution changes $R_{\Sigma\Lambda}$ (Eq.8) into

$$R_{\Sigma\Lambda} = R^{o}_{\Sigma\Lambda} \frac{G_1 - \frac{4}{3}G_3(m_\Sigma - m_\Lambda)}{G_1 + \frac{4}{3}G_3(m_\Sigma - m_\Lambda)} \tag{11}$$

It is of the utmost importance to improve on the rather poor experimental value of

$$R^{exp}_{\Sigma\Lambda} = 0{,}63 \pm 0{,}22 \quad . \tag{12}$$

3. The Mesonic Current of Second Class.

The current (10) contains derivatives and is thus suppressed kinematically in the vertex function. A non-derivative second class current can be constructed out of mesons[6,7)]

$$j^M_\lambda = i\, V_\lambda (D_1 + iD_2) P + \alpha\, \omega_\lambda\, \pi^- + \beta\, \rho^-_\lambda\, \eta' \quad , \tag{13}$$

where α and β are parameters showing possible different strengths for the 3 contributions. D_i are the well-known 8×8 matrices of the d-type SU_3 coupling

$$(D_i)_{jk} = d_{ijk} \tag{14}$$

If the current (13) is coupled to the lepton current ℓ_λ by

$$\mathcal{H}_{II} = G_d\, j^M_\lambda\, \ell^\lambda + h.c. \tag{15}$$

it gives a prominent contribution to $R_{\Sigma\Lambda}$ via the graphs of Fig. 1. Taking strong interaction vertices from SU_3

and neglecting orders of q^2 gives the following second class contribution to the axial vector form factor of $\Sigma\Lambda$ decay[7])

$$G_1^{II} = \pm\frac{1}{8\pi^2}\frac{1}{\sqrt{3}}\frac{G_d f_V}{G_V m_{K^*}^2}\log\frac{\lambda}{m_\Sigma}(3f_P(m_\Sigma - m_\Lambda) + 2d_P(m_\Xi - m_N)) \quad . \tag{16}$$

Here small f's and d are the strong coupling constants and λ is a cut-off.

Because of the poor experimental value (12), the coupling constant G_d can presently only be restricted to be about one order of magnitude below the weak four fermi coupling for any reasonable cut-off λ.

4. Hypercharge Changing Transitions.

For $|\Delta Y| = 1$ transitions, second class currents could be induced by symmetry breaking effects without separate current operators like Eqs.(10) or (13). However, the naive Cabibbo-model, which works so surprisingly well in its straightforward application, would tell us that they are absent.

A possible way to check the existence of G_3 (or F_3) experimentally is the μ-e ratio. To first order in the baryon mass difference, it is given by[8])

$$\frac{\Gamma(A\to B+\mu+\bar{\nu}_\mu)}{\Gamma(A\to B+e+\bar{\nu}_e)} = r(x) - \frac{G_1G_3+F_1F_3}{(F_1)^2+3(G_1)^2}(m_A-m_B)\,t(x) \tag{17}$$

with

$$r(x) = \frac{1}{2}(1-x^2)^{3/2}(2-7x^2) + \frac{15}{4}x^4(L - 2\sqrt{1-x^2})$$

$$t(x) = 5x^2[\frac{1}{2}\sqrt{1-x^2}(2+13x^2) - \frac{3}{4}x^2(x^2+4)L] \tag{18}$$

$$L = \log \frac{1 + \sqrt{1-x^2}}{1 - \sqrt{1-x^2}}$$

$$x = m_\mu/(m_A - m_B) \quad .$$

Hyperon beams and larger bubble chambers should soon decrease experimental uncertainties in the μ-e ratio enough to allow for a check on the presence of F_3 or G_3 in hyperon decays.

5. A Decuplett Current of Second Class.

If we want to keep an open mind for all possibilities, we can think of second class currents outside the conventional octett assumptions[9]. We do not attempt to propose a particular form for such a current. In table 1, we collect upper limits on G_3 from decays which are forbidden for octett currents only.

Table 1: Limits on G_3 from decuplett transitions

$\Sigma^+ \to n e^+ \nu_e$	$\lvert G_3 \rvert < 0{,}26$
$\Xi^- \to n e^- \nu_e$	$\lvert G_3 \rvert < 1{,}47$
$\Xi^0 \to p e^- \nu_e$	$\lvert G_3 \rvert < 0{,}41$

Another way to investigate the possible existence of peculiar second class currents shall be opened by future neutrino experiments[10]. It is of great importance to use all possible means for ruling out or establishing the existence of any other current than the well-known V and A currents. This question relates to the most fundamental one of the structure of the interactions which we observe in elementary particle reactions.

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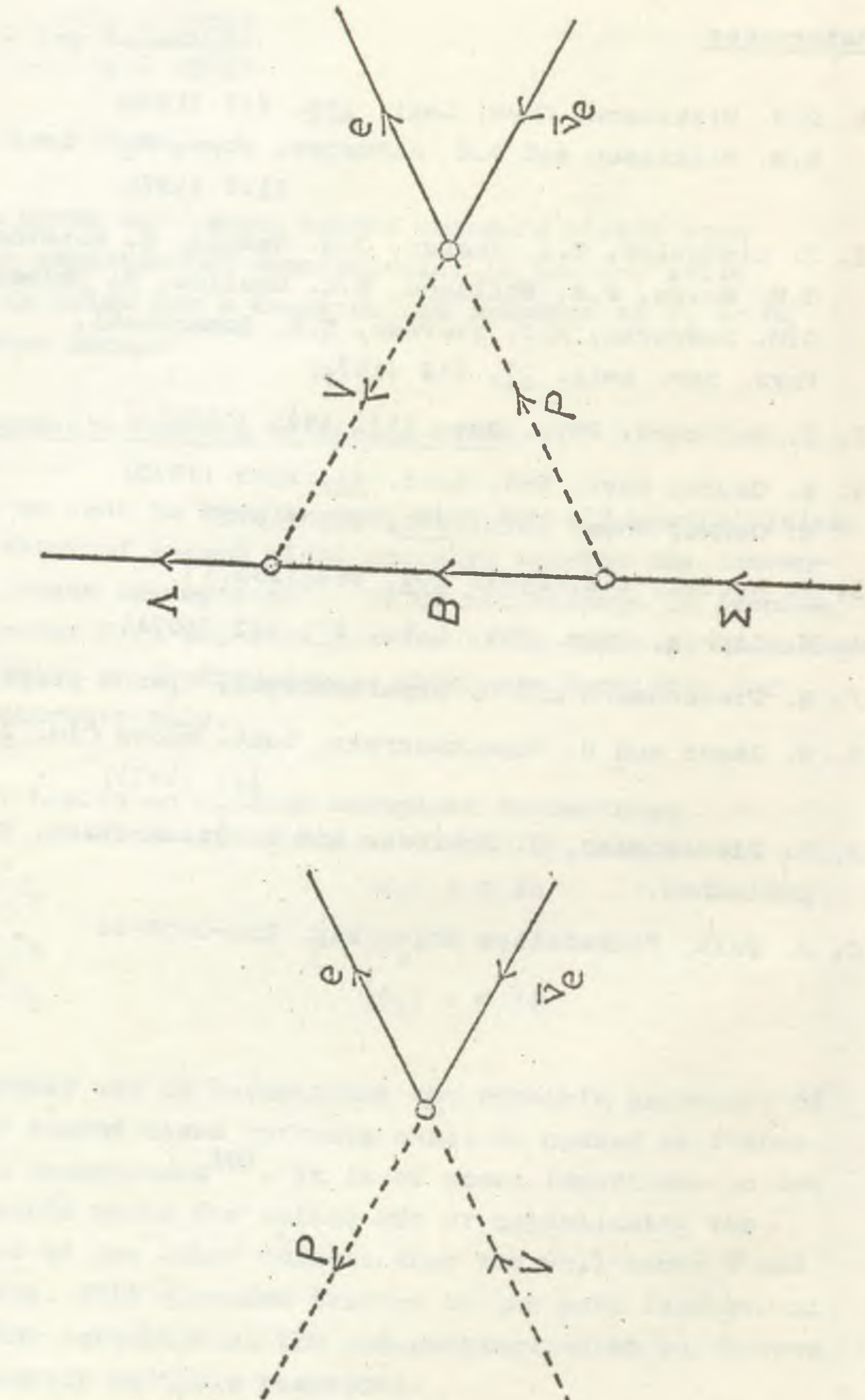

FIG. 1

INDUCED TENSOR INTERACTION AND SECOND-CLASS CURRENTS

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I. INTRODUCTION

Soon after the advent of the V-A theory of weak interactions, the realization that the proposed weak hadronic currents imply special isospin and Hermitian properties led to investigations. The concept of first- and second-class currents was formed /1/ and the first possible nuclear beta-decay experimental tests were discussed. Searching for second-class current effects very soon singled out the so-called induced-tensor term in the axial-vector current interaction /1/, which was later investigated in nuclear beta-decay theory /2,3,4,5/. Although it is well known that second-class currents (SCC) can lead to many additional effects, contributing even to the interaction terms normally due to first-class currents /6/, in this contribution we intend to concentrate on the induced-tensor contribution. The contributions of SCC to other axial-vector form factors are due to the breaking of symmetry and therefore might not be very important in the case of nuclear beta decay.

Recent papers of Wilkinson /14/ have initiated renewed interest in the question of SCC /7,8,9,10,11, 12,13/. This question can still be considered opened at present in spite of some negative evidence /15,16, 17/ which has been found. Since the induced tensor was mentioned and investigated in our earlier papers /3,18,19,20,21,22/, we want to make our previous work up to date and fully relevant to the present level of knowledge. The induced-tensor term has practically been the only one investigated in beta-decay calculations so far, so a thorough and complete approach is strongly indicated. The impulse approximation was not quite correctly carried out in many of the previously published beta-decay calculations. It turned out that an important contribution was unjustifiably neglected /2, 3,4/, thus leading to statements about different induced-tensor forms /10/.

After discussing this problem, we present results of the analysis concerning ft values, β-γ angular correlations, and parity-violating processes in heavy nuclei /23/. Besides these results, which are to be published soon, we present evidence for forbidden beta-decay transitions ($0^- \to 0^+$ and unique) and briefly discuss implications for β^+/K ratios /24/. Depending on the interpretation, ft values can be taken as a strong indication in favour of the existence of the induced tensor.

Such an interpretation has run into experimental difficulties /17/; besides, it appears to be in blatant disagreement with experimental β^+/K ratios /15/. Other evidence is inconclusive because of both interpretational difficulties and theoretical or/and experimental uncertainties. It appears, therefore, natural to turn to alternative SCC, namely, to mesonic ones /11,25,52/.

The existence of SCC has been searched for, as it should be, but the final answer is still forthcoming. Obviously, it will depend strongly on the solution of the intricate Coulomb-interaction interplay of nucleons among themselves, and of nucleons with leptons.

II. IMPULSE APPROXIMATION

When performing theoretical calculations for nuclear beta decay or related processes, one is always forced to use the impulse approximation as a last resort. Suitable form-factor formulations can sometimes present some general ideas in a very adequate manner /9,13/, but when actual calculation commences they always have to be culled down to the nonrelativistic approximation (NRA). In this scheme, questions concerning the relevance of the induced tensor and its connection with SCC's /11/ have been successfully cleared up /12/. The numerical calculations involving the induced tensor have so far been limited to two forms, which are equal for nucleons on the mass shell

$$(-)f_T(q^2)\ \frac{1}{2M}\ \bar{u}(p')\sigma_{\mu\nu}(p-p')_\nu\ \gamma_5 u(p)\ , \tag{2.1a}$$

$$(-)f_T(q^2)\ \frac{1}{2M}\ \bar{u}(p')\,i(p_\mu+p'_\mu)\gamma_5 u(p)\ . \tag{2.1b}$$

Expression (2.1a) is the one used in previous beta-decay calculations /2,3,4,5,18,19,20,22/ with the weak-interaction Hamiltonian given in the NRA by Case B*)

$$H_{int} = (g_A \mp \frac{Y}{2M}E_o - \frac{Y}{2M}2\xi)\vec{\sigma}\cdot\vec{L}_4 \pm$$
$$\pm \frac{Y}{2M}\vec{\sigma}\cdot(-i\vec{\nabla})L_4\ . \tag{2.2}$$

The notation has the usual meaning.**)
Equation (2.2) was derived from the following "relativistic" expression for the induced tensor

$$IT = -\frac{Y}{2M}\frac{1}{2}(\gamma_\mu\gamma_\nu-\gamma_\nu\gamma_\mu)\gamma_5\partial_\nu\tilde{L}_\mu\ , \tag{2.3a}$$

$$\tilde{L}_\mu = \Psi_e^\dagger\gamma_4\gamma_\mu\gamma_5\Psi_\nu\ . \tag{2.4}$$

In deriving the NRA of Eq. (2.2) the contribution coming from the combination of indices

---

*) This name is introduced for later reference.

**) g_A is the axial vector coupling constant, $Y = f_T(0)$,

$$\mu = j; \quad \nu = k \qquad\qquad j,k = 1,2,3 \tag{2.5}$$

was neglected. This contribution, denoted by D, can be written in the form

$$D = D_1 + D_2 , \tag{2.6}$$

$$D_1 = -\frac{\gamma}{2M} \Psi_f^\dagger \gamma_4 |(\vec{\sigma}\,\vec{p})(\vec{\alpha}\,\vec{L}) + (\vec{\alpha}\,\vec{L})(\vec{\sigma}\,\vec{p})|\Psi_i , \tag{2.7}$$

$$D_2 = \frac{\gamma}{2M} \Psi_f^\dagger \gamma_4 \gamma_5 |(\vec{p}\,\vec{L}) + (\vec{L}\,\vec{p})|\Psi_i . \tag{2.8}$$

Here Ψ_i and Ψ_f refer to the initial and final "relativistic" nucleon wave functions, respectively, and the impulse p is taken to be an operator. The NRA follows from these functions by approximating

$$\Psi_n \simeq \begin{pmatrix} -(\vec{\sigma}\,\vec{p})/2M & \chi_n \\ & \chi_n \end{pmatrix} . \tag{2.9}$$

Here χ_n are nuclear shell-model functions satisfying

$$p^2\chi_n = 2M(E_n - V)\chi_n . \tag{2.10}$$

Nucleons are interacting only with the shell-model potential V. One then obtains

$$\begin{aligned} D_1 &\cong -\frac{\gamma}{(2M)^2} \chi_f^\dagger |p^2(\vec{\sigma}\,\vec{L}) - (\vec{\sigma}\,\vec{L})p^2|\chi_i = \\ &= \frac{\gamma}{2M}(E_f - E_i)\chi_f^\dagger \vec{\sigma}\vec{L}\chi_i . \end{aligned} \tag{2.11}$$

This contribution just cancels the second and the third term appearing in Eq. (2.2). $E_f - E_i$ in Eq. (2.11) goes into $E_o + 2\xi$.

In the case of forbidden beta transitions, additional terms are of importance. One can extract from the term D_2, Eq. (2.8), the term

$$D_{2b} \cong \frac{\gamma}{2M^2}(-i\vec{\sigma}\,\nabla)\vec{L}_4\,\vec{p} , \tag{2.12}$$

and an additional contribution comes from the combination of indices in Eq. (2.3)

$$\mu = 4 \qquad \nu = j . \tag{2.13}$$

This contribution, denoted by C, is

$$\begin{aligned} C &\cong \frac{\gamma}{2M}(-i\vec{\sigma}\,\nabla\, L_4) + \frac{\gamma\,\Omega}{4M^2} L_4\,\vec{\sigma}\,\vec{p} , \\ \Omega &= E_o + 2\xi . \end{aligned} \tag{2.14}$$

E_o is the maximal lepton energy, $\xi = Z\alpha/(2r_o)$, $\vec{L}_4 = \Psi_e^\dagger\,\vec{\sigma}\,\Psi_\nu$, $L_4 = \Psi_e^\dagger\gamma_5\Psi_\nu$.

The alternative form (2.1b) of the induced tensor leads to the effective beta-decay interaction

$$(IT) = -\frac{Y}{2M}\,\Psi_f\gamma_5 i\,|\,p_\mu\tilde{L}_\mu + \tilde{L}_\mu p_\mu\,|\,\Psi_i\ , \qquad (2.15)$$

Expression (2.15) in the NRA is completely equivalent to Eq. (2.3a). The selection of indices $\mu = j$ leads to the term D_2, Eq. (2.8), while $\mu = 4$ in the NRA leads to the terms appearing in Eqs. (2.2) and (2.11).

It appears that forms (2.1a) and (2.1b) lead to the same NRA in the usual nuclear shell-model impulse-approximation limit. They are, however, not equivalent when exchange effects, i.e., two-body contributions to the weak-interaction Hamiltonian, are considered.

Let us illustrate this by a simple one-pion-exchange diagram. The nucleon propagator has the form

$$P = \sum_{s,E_s \lessgtr 0} \frac{u^s(p_s)\,\overline{u}^s(p_s)}{E_s + \omega - E_1}\ . \qquad (2.16)$$

Its $E_s > 0$ part contributes to the nonrelativistic perturbation series, while its $E_s < 0$ part leads to the two-body exchange contribution

$$\sum_s \overline{u}(p_2)\gamma_5\tau_j\,\frac{1}{2}\,(1+\beta)\,O_T\,u(p_1)R\ . \qquad (2.17)$$

The contributions of the type (2.17) were estimated in Ref. 10 for a special case of the conserved second-class axial vector current. In this special model, the induced tensor of the form (2.1a) receives no significant contribution, leading to Case A (2.18). Two-body contributions from the induced tensor (2.1b), when approximated through the effective single-body operators, are significant. They add with the single-particle contribution (2.18) in such a way that Eq. (2.2) results. This leads to Case B_2 (2.19) in the NRA.

On the basis of this analysis, Cases A and B_2 have to be confronted with experiment, for which purpose we summarize our results regarding the NRA.

Case A

$$H_{int} = g_A\,\vec{\sigma}\,\vec{L}_4 + \frac{Y}{2M}\,\vec{\sigma}\,(-i\nabla)L_4 + \qquad (2.18)$$

$$+ \frac{Y}{2M}\,\{\,\frac{1}{2}\,(E_o + 2\xi)L_4\,\frac{\vec{\sigma}\,\vec{p}}{M} + (-i\vec{\sigma}\nabla)\,\frac{\vec{L}_4\vec{p}}{M}\,\}\ .$$

Here all signs refer to β^- decay. The terms inside the wavy brackets are important only in forbidden decays.

Case B_2

$$H_{int} = H_{int}(2.18) \quad \text{with} \quad g_A \to g_A - \frac{Y}{2M}(E_o + 2\xi) \tag{2.19}$$

III. ft VALUES AND MIRROR TRANSITIONS

As comparisons of ft values have aroused renewed interest in SCC´s, it is suitable to use these values as a reference frame for other analysis. The characteristic parameter is the ratio

$$b = Y|g_A|^{-1} . \tag{3.1}$$

Instead of merely quoting the values for b following from the analysis /14/, we recalculated them using Bhalla and Rose´s wave functions /26/ for electrons and taking into account the vector-axial-vector interference corrections. The correction factors necessary to the spectrum shape were published a long time ago /3/. We calculated the experimentally meaaured half-lives t by the formula

$$t \int_{o}^{E_o^2-1} p^2 q^2 F(Z,E)\, \tilde{C}^{\beta}(E,b)\, dp = \frac{(2\pi)^3 \ln 2}{g_A^2 |\langle \vec{\sigma} \rangle|^2} , \tag{3.2}$$

where E_o is the maximal energy, E is the electron energy, p and q are the electron and neutrino momenta, respectively. F(Z,E) is the Fermi function, while $\tilde{C}^{\beta}(E,b)$ is the spectrum-shape correction factor depending on the induced tensor through the b of Eq. (4.1).

The ft values for electron and positron mirror transitions are expected to be approximately equal. The deviation from this equality

$$\delta = \frac{(ft)^+}{(ft)^-} - 1 \tag{3.3}$$

can be interpreted /14/ as an indication for the existence of the induced tensor. We calculated δ_o by putting b = 0 in Eq. (3.3). The value of b, (i.e., Y) was selected in such a way as to give δ = 0 when using the full expression for $\tilde{C}^{\beta}$.

The results for Case 1 in Table I were obtained by neglecting the cross terms with the vector coupling, (i.e., weak magnetism etc.). The values for b in Case 2 in Table I were found taking into account all corrections. The calculations were performed both for Case A (2.18) and for Case B (2.19).

In Table II we recalculated the results of Table I taking into account particular /27/ corrections to the matrix element ratios, which are actually the most favourable ones for the existence of the induced tensor.

Table 1

A	E_o^+ / E_o^- (MeV)	t^+ / t^- (sec)	1: δ_o	1: b_A	1: b_B	2: δ_o	2: b_A	2: b_B
8	14.05	0.774	0.072	6.6	-2.9	0.067	6.2	-2.9
	13.10	0.849						
12	16.32	0.0116	0.122	9.6	-4.8	0.111	9.2	-4.8
	13.37	0.0210						
20	11.26	0.51	0.046	7.3	-3.7	0.045	7.3	-3.7
	5.40	11.03	-0.064	-9.2	4.6	-0.061	-9.9	5.1
24	8.737	26.12	-0.018	-4.8	2.2	-0.017	-4.8	2.2
	1.392	53.80	0.37					
28	11.537	0.52	0.282	45.9	-22.4	0.270	44.8	-21.3
	2.856	136.4						

Table II

A	δ_W [a]	$\delta_o-\delta_W$ [b]	b_A^W	b_B^W
8	0.048	0.020	1.8	-0.9
12	0.098	0.013	1.0	-0.44
20	0.045	0.00	0	0
24	0.040	-0.058	-14.7	7.0
28	0.051	0.22	35.2	-16.9
18	0.007	0.026	23.9	-12.9
30	0.005	-0.057	-77	42

[a] δ_W from Ref. 27 is defined as Λ^-/Λ^+-1.

[b] δ_o is the value from our case 2 in Table 1.

As is to be expected, the values of the induced-tensor constant b are generally reduced. The additional corrections seem to be responsible for most or even for the whole (A = 20) of the apparent effect. This gives additional weight to the outcome of the Li^8 and B^8 experiment /17/ and to the discrepancy in the β^+/K ratio /15/, to which we turn shortly.

IV. CAPTURE TO β^+ RATIOS

The analysis already published in Ref. 15 used the NRA, denoted by us as Case B_2 (2.19). As indicated in this reference, the same b could satisfy neither $(ft)^+/(ft)^-$ nor K/β^+ ratios. It seems important to point out that Case A (2.18) in the NRA does not change this conclusion. The correction factor ratio for Case B_2 obtained by the simplest approximation is

$$C^K/C^+_\beta = 1 - \frac{4}{3}\, b_B(4\xi - E_o^+) \quad . \qquad (4.1)$$

Case A (2.18) in the same approximation leads to

$$C^K/C^+_\beta = 1 + \frac{2}{3}\, b_A(4\xi - E_o^+) \quad . \qquad (4.2)$$

We see from Tables I and II that the ft values can be fitted only with $b_A > 0$, which leads to an increase of the C^K/C^+_β ratios with increasing atomic number Z, in complete disagreement with experiment /15/.

V. SPECTRUM-SHAPE FACTORS FOR $0^- \to 0^+$ AND UNIQUE NUCLEAR β TRANSITIONS

The spectrum shape for forbidden decays does depend on the induced-tensor contribution, in contrast to the case of allowed decays.

In the past, analyses were made /18,19,20,22/ for the approximation case B . Let us compare these results with those obtained for the approximation case A.

$0^- \to 0^+$ nuclear beta transitions were analyzed /24/ using the formalism already published in Ref. 22. The fits to the Daniel-Kashl measurements /28/ are summerized in Table III. In this table χ^2 measures the

Table III

	Approx. A				Approx. B			
	f	b	χ^2	$\frac{\langle\vec\sigma\cdot\vec r\rangle}{r_o}$	f	b	χ^2	$\frac{\langle\vec\sigma\cdot\vec r\rangle}{r_o}$
^{144}Pr	+10.4	-13.0	2.68	1.38	+11.6	-13.7	2.68	1.15
^{166}Ho	+13.6	-8.7	2.56	0.34	+14.6	-8.3	2.56	0.34

goodness of the fit, and f is the matrix element ratio

$$f = \left\langle \frac{i\,\vec\sigma\cdot\vec p}{M} \right\rangle / \langle \vec\sigma\cdot\vec r \rangle \; .$$

There is not much difference between approximations A and B which is easily explainable when theoretical expressions are studied. It appears that the respective contributions differ only in the constant additive terms. $b < 0$ is clearly favoured in both cases. This conclusion, however, can be changed if large variations of the electron wave functions inside the nuclei are allowed. In fact, it is necessary to introduce two additional parameters, which are varied independently inside physically acceptable limits. The values $b > 0$

may be found, but in the case of ^{144}Pr the fit is much worse, χ^2 being 28.2.

Unique nuclear beta transitions were analyzed using a formalism described in great detail in Ref. 19. The spectrum shape depends on the five-nuclear-matrix-element ratios calculated in the shell-model approximation. The resulting fits thus depend only on b, and are summarized in Table IV. It seems that the positive value for the

Table IV

	B		A	
	b	χ^2	b	χ^2
^{42}K	17.0	1.02	14.9	1.04
^{86}Rb	2.65	0.22	2.30	0.22
^{90}Sr	1.50	3.13	1.25	3.20
^{90}Y	-4.72	2.47	-4.42	2.45
^{142}Pr	0.73	1.27	0.55	1.28
^{166}Ho	34.0	1.26	18.5	1.26

induced tensor, i.e., b > 0 is favoured, in contrast to the indication for $0^- \to 0^+$ transitions. An extensive analysis of a newer measurement of $0^- \to 0^+$ and $0^- \to 2^+$ transitions together with a β-γ correlation measurement, both performed for Ho^{166} nuclei /30/, also allowed for some variations in nuclear-matrix-element ratios. The result was b > 0, with b = 0 being also acceptable for reasonable values of the parameters.

Generally speaking, it appears that $0^- \to 0^+$ and unique transitions are not very consistent with the existence of the induced tensor. However, experiments are uncertain, the experimental results of the different groups being sometimes contradictory. We have mostly /29/ used measurements of a particular group, in order to achieve an overall comparison among various nuclei.

VI. β-γ ANGULAR CORRELATIONS

In the description of β decay the induced tensor appears as an additional small correction, which has to compete with large main contributions and with other possible small corrections. Obviously, it would be useful to investigate experiments where the large main term does not contribute, since it is forbidden by the selection rules. Thus it is worth while to turn our attention to β-γ angular correlation measurements in the case of allowed β decays. In this case the whole effect is directly proportional to the induced terms and the so-called "second-order corrections". The sign of the induced-tensor contribution changes going from β^- to β^+ decays, depending also on the γ transition. In general, the whole effect depends strongly on the sign

TABLE V. Dependence of the angular-correlation coefficients on the induced coupling constant δ.

Transition	$\frac{W_0}{W}$	$10^3 A_{2\,exp}$	$10^3 A^{\beta}_{2\,th}$ Case A $\delta=-0.002$	Case A $\delta=0$	Case A $\delta=0.002$	Case A $\delta=0.004$	Case B $\delta=-0.002$	Case B $\delta=0$	Case B $\delta=0.002$	Case B $\delta=0.004$
^{20}F $2^+(\beta^-)2^+$	11.6 >9	-10 ± 3 [a]	3.7 ± 0.4	-7.5 ± 1	-19 ± 2	-30 ± 4	8 ± 1	-3.2 ± 0.2	-14 ± 1	-25 ± 2
^{24}Na $4^+(\beta^-)4^+$	3.15 3.1	-0.4 ± 0.8 [b]	1.1 ± 0.1	-2.2 ± 0.1	-5.5 ± 0.2	-8.8 ± 0.3	1.7 ± 0.1	-1.64 ± 0.05	-4.9 ± 0.2	-8.2 ± 0.3
^{22}Na $3^+(\beta^+)2^+$	2.06 1.78	1.7 ± 2 [c] 3.0 ± 0.5 [b] 25 ± 5 [d] 0.5 ± 0.8 [e] 1.7 ± 8 [f]	0.30 ± 0.02	-0.19 ± 0.01	-0.69 ± 0.04	-1.18 ± 0.08	0.32 ± 0.01	-0.17 ± 0.01	-0.66 ± 0.04	-1.15 ± 0.08
^{56}Co $4^+(\beta^+)4^+$	3.935 3.348	-1.8 ± 4.1 [c] -7 ± 6 [g] -25 ± 4 [e] 2 ± 7 [f]	-2.2 ± 0.3	1.5 ± 0.15	5.0 ± 0.5	8.6 ± 0.9	-3.5 ± 0.4	0.19 ± 0.08	3.8 ± 0.5	7.6 ± 1
^{46}Sc $4^+(\beta^-)4^+$	1.4 1.35	-0.4 ± 0.8 [b]	0.38 ± 0.04	-0.78 ± 0.08	-1.9 ± 0.2	-3.1 ± 0.3	0.79 ± 0.09	-0.37 ± 0.03	-1.54 ± 0.16	-2.7 ± 0.3
^{58}Co $2^+(\beta^+)2^+$	1.95 1.68	0.5 ± 2.4 [c]	-0.93 ± 0.07	0.61 ± 0.05	2.16 ± 0.16	3.70 ± 0.28	-1.54 ± 0.10	0.03 ± 0.06	1.55 ± 0.13	3.09 ± 0.24
^{60}Co $5^+(\beta^-)4^+$	1.613 1.4	0.4 ± 0.3 [h] 0.7 ± 1.1 [c] 0.4 ± 0.7 [b] -0.7 ± 2.2 [i]	-0.15 ± 0.02	0.30 ± 0.04	0.76 ± 0.09	1.2 ± 0.2	-0.36 ± 0.05	0.09 ± 0.01	0.55 ± 0.07	1.0 ± 0.14
^{110}Ag $6^+(\beta^-)6^+$	2.035 1.78	0.6 ± 1 [c]	0.50 ± 0.03	-1.02 ± 0.07	-2.5 ± 0.1	-4.1 ± 0.3	1.7 ± 0.1	0.23 ± 0.02	-1.3 ± 0.1	-2.8 ± 0.2
^{134}Cs $4^+(\beta^-)4^+$	2.292 1.78	-1.5 ± 2.2 [c] 37 ± 7 [j] -2 ± 7 [f]	0.46 ± 0.03	-0.94 ± 0.07	-2.3 ± 0.1	-3.7 ± 0.2	1.7 ± 0.1	0.37 ± 0.03	-1.03 ± 0.07	-2.4 ± 0.2
^{154}Eu $3^-(\beta^-)2^+$	2.135 1.88	0.7 ± 2 [c] 7 ± 6 [h] 1.7 ± 3 [k]	-0.12 ± 0.01	0.25 ± 0.03	0.62 ± 0.08	1.0 ± 0.1	-0.52 ± 0.07	-0.14 ± 0.02	0.22 ± 0.02	0.60 ± 0.07
^{124}Sb $3^-(\beta^-)3^-$	2.215 1.88	-2.2 ± 2.8 [c] 6 ± 7 [f] 6 ± 8 [i]	0.51 ± 0.06	-1.0 ± 0.1	-2.6 ± 0.3	-4.2 ± 0.06	1.9 ± 0.3	0.32 ± 0.05	-1.2 ± 0.2	-2.8 ± 0.4
^{152}Eu $3^-(\beta^-)3^-$	2.37 >1.88	-17 ± 12 [h]	0.45 ± 0.12	-1.0 ± 0.2	-2.4 ± 0.4	-3.9 ± 0.6	2.0 ± 0.4	0.59 ± 0.11	-0.86 ± 0.13	-2.3 ± 0.4

TABLE V *(Continued)*

Transition	$\frac{W_\theta}{W}$	$10^3 A_{2exp}$	$10^3 A^{\beta}_{2th}$ Case A $b=-0.002$	$b=0$	$b=0.002$	$b=0.004$	Case B $b=-0.002$	$b=0$	$b=0.002$	$b=0.004$
^{152}Eu $3^-(\beta^-)2^-$	1.39 >1.31	16 ± 20 [h]	−0.08 ± 0.01	0.15 ± 0.01	0.39 ± 0.02	0.63 ± 0.02	1.4 ± 0.1	1.55 ± 0.11	1.8 ± 0.2	2.2 ± 0.3
^{160}Tb $3^-(\beta^-)3^-$	2.108 2.1	90 ± 30 [h] 40 ± 20 [h] 170 ± 30 [k]	0.44 ± 0.05	−0.89 ± 0.10	−2.2 ± 0.2	−3.5 ± 0.5	1.9 ± 0.2	0.57 ± 0.06	−0.76 ± 0.06	−2.1 ± 0.2

of the induced-tensor term, as can be seen from the results presented in Table V. This should, in principle, enable us to distinguish Case A from Case B defined in Sec. II. In both cases, only the term

$$\frac{Y}{2M}\vec{\sigma}\cdot(-i\vec{\nabla})L_4 \qquad (6.1)$$

can contribute. It, therefore, shows a certain similarity with other angular correlations, such as electron-neutrino or electron-nuclear spin /10,31/.

The standard expression for β-γ angular correlations is

$$W(\theta) = 1 + A_2(\beta)F_2(LLI_fI)P_2(\cos\theta) , \qquad (6.2)$$

where P_2 is the Legendre polynomial and F_2 is the well-known angular-momentum-dependent spectroscopical factor. The factor $A_2(\beta)$ depends on the particulars of weak interactions, and in the simplest approximation in which the so-called "second-order-correction" contributions are neglected and plane waves are used for electrons it is given by

$$A_2(\beta) = \frac{4}{3}\frac{p^2}{W}\left[\frac{1}{2M} \pm \frac{g_V}{g_A}\frac{(\mu_p-\mu_n)}{2M} \pm \frac{b}{2M}\right] F_2(11I_iI) , \qquad (6.3)$$

where p and W are the momenta and the energy of the electron (positron), respectively.

Although several measurements of β-γ angular correlations have appeared in the literature /32-43/, most of them are still associated with large uncertainties and experimental errors. We decided to calculate $A_2(\beta)$ for all of them on the basis of the full formula using Bhalla and Rose's wave functions /26/ for electrons and including all "second-order corrections". Final results were obtained by averaging over the experimental range of electron energy. The errors quoted with the theoretical results correspond to the difference between β-γ correlation values for the minimal or maximal experimental energies.

It seems that $b > 0$ is favoured by most of the experiments. If anything, the results seem to be an additional piece of evidence for the existence of weak magnetism, the induced tensor itself not being particularly needed.

VII. PARITY VIOLATION IN HEAVY NUCLEI

It was pointed out a long time ago /44/ that weak parity-violating nuclear forces are a natural consequence of the current-current model of the weak Hamiltonian. In the lowest so-called "factorization" approximation the weak parity violating (PV) internucleon two-body potential is proportional to

$$V_{PV} \sim \langle N_2'|J_{\mu_5}(x)|N_2\rangle \langle N_1'|J_\mu(x')|N_1\rangle f(x,x') , \qquad (7.1)$$

where J_{μ_5} and J_μ are the axial-vector and the vector hadronic weak current, respectively. It follows from (7.1) that when applied to PV nuclear forces, the two cases of the NRA, Eq. (2.18) and Eq. (2.19), are equivalent.

One can use in the first approximation the potential already published in /45,21/

$$V_{IT} = \frac{YG}{4\sqrt{2}\,M} (\vec{\sigma}_1+\vec{\sigma}_2)\ |\vec{p}_{12}, f(r)|_- T^{(-)} , \tag{7.2}$$

where

$$T^{(-)} = \tau_{1+}\tau_{2-} - \tau_{1-}\tau_{2+} \tag{7.3}$$

and

$$f(r) \simeq \frac{1}{4\pi} \frac{e^{-m_\rho r}}{r} , \tag{7.4}$$

m_ρ being the ρ meson mass.

Some discussion on the value of Y, based on circular polarization measurements, has already been presented /46/. As the approximation used in Ref. 46 overestimates the effect /47/, some change in the conclusions is to be expected.

One group of experimental results for the circular polarization of γ emission in the 482-keV decay of Ta^{181} /48/ leads to an average

$$P_{exp} = (-5.0 \pm 1.5)\times 10^{-6} .$$

An essentially different sign is given by Kuphal, Daum, and Kankeleit /49/, where

$$P_{exp} = (2.0 \pm 4.0)\times 10^{-6} .$$

In the wide spectrum of the published experimental results one can also find some giving larger P than those quoted above /50/. They would naturally set higher bounds on Y, and therefore we are actually dealing with the most sensitive cases.

Looking for such combinations of numbers which would fit the quoted results, we obtain the limits for the conventional weak Hamiltonian

$$-1.7 \leq Y \leq 7 .$$

If additional residual nucleon-nucleon interactions are introduced in theoretical calculations, in the way outlined by Vinh-Mau /51/, the theoretical value for P is decreased, so we can estimate

$$-15 \leq Y \leq 19 .$$

For other weak-Hamiltonian models, the situation is either approximately the same as for the conventional model, or the induced-tensor contribution is relatively in-significant in comparison with the strongly enhanced one-pion exchange contribution. In that case even larger limits on Y result. None of these limits actually contradicts the estimates of Y based on the ft value.

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NUCLEAR BETA - DECAY EXPERIMENTS AND SECOND - CLASS WEAK CURRENT

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Abstract: Experimental results on β- decay probabilities connected with the question of second - class currents are discussed.

It was shown in 1958 by Weinberg [1/] that the transformation properties of the induced scalar and tensor /IT/ interactions are opposite to those of the basic V and A interactions, respectively, under G - transformation. The existence of the induced scalar interaction is limited by the CVC theory, therefore, the only way to decide the question of the conservation of G - parity in weak interactions is to see wether the IT interaction does or does not exist. Although different kinds of angular correlation in beta - decay can be used to solve this question [2/] nevertheless comparisons of decay probabilities $\left[(ft)^{+}/(ft)^{-}\right.$ and $\left.K/\beta^{+} \text{ ratios}\right]$ gave more success. In the following only these results will be discussed.

1. <u>$(ft)^{+}/(ft)^{-}$ ratios of mirror β-transitions.</u>

Neglecting all possible differences of mirror β-transitions except the one caused by the IT interaction, the deviation of the $(ft)^{+}/(ft)^{-}$ ratios from unity is expected to be proportional to the sum of the end - point energies of $\beta^{\pm}$ - transitions.

$$\frac{ft^{+}}{ft^{-}} = \frac{C^{-}}{C^{+}} = 1 - \frac{4}{3}\,\frac{C_{IT}}{|C_A|}\,(W_0^{-} + W_0^{+}) \qquad /1/$$

All remaining terms of the IT interaction are an order of magnitude smaller than this, because of the large end - point energies. To prove this proportionality is of great importance, since otherwise the poor knowledge of the nuclear matrix elements does not allow to make a definite conclusion.

The semi - experimental $(ft)^{+}/(ft)^{-}$ ratios at high energies give $C_{IT}/|C_A| = -2\times10^{-3}$ ref.[3,4], but it is impossible to prove the linear dependence /1/ using these results, because the low energy β^{+} and β^{-} decaying mirror pair does not exist / see fig. 1 /. The two positron decaying mirror pairs with A = 18 and 30 do not contradict to the linear dependence, but they cannot be used to prove it either, because of the small energy difference of the transitions [5].

2. K/β^+ ratios. Using the same approximation as in the previous point, the influence of the IT interaction can be given as

$$K/\beta^+ = C_x/\,C^+ = 1 - \frac{4}{3}\ \frac{C_{IT}}{|C_A|}\,(2W_o + 3\xi) \qquad /2/$$

where $\xi = \alpha Z/2R$. The K/β^+ ratios are known in transitions with $W_o \sim 1 \ll \xi \sim 10$, such that W_o can be neglected in the above equation.

As it has been shown in ref. [5/], $C_{IT}/|C_A| \approx - 2\mathrm{x}10^{-3}$ $[m_oc^2]^{-1}$ does not agree with the experimental K/β^+ ratios /fig.2 /. It would be difficult to use these results to give an upper limit for the coupling constant of the induced tensor interaction, but it is clear from fig.2 that the sign and magnitude of the above C_{IT} is not consistent with the K/β^+ ratios. We note that we can rely on the equality of the nuclear matrix elements in EC and β^+ decays, only second order correction terms might change the K/β^+ ratio.

3. Excitation spectra of ^{8}Be in the decays of ^{8}Li and ^{8}B. In these cases the energy dependence of the ratio of $\beta^\pm$ decay probabilities has been determined in one experiment by Wilkinson and Alburger [6/]. The result is nearly independent of the differences of nuclear overlap, therefore, this experiment is especially important.

The ^{8}Be decays into two α- particles from the excited state. The number of α- particles $(N^{\pm})$ with energy E_x is proportional to the total number of $\beta^{\pm}$ - decays with end-point energy $W_o^{\pm} = W_{max}^{\pm} - E_x$, where $W_{max}^{\pm}$ is the maximum available energy of the β - decays, such that $(ft)^{+}/(ft)^{-}$ can be replaced by $N^{-}(E_x)/ f(W_o^{-})$: $N^{+}(E_x)/ f(W_o^{+})$ in eq. /1/. The coupling constant of the IT interaction is then determined from the slope of the $(ft)^{+} /(ft)^{-}$ values vs. $(W_o^{+} + W_o^{-})$. Fig.3 shows, that second - class current with C_{IT} / C_A 2×10^{-3} $[m_oc^2]^{-1}$ can be excluded. The authors give an upper limit for the above ratio as 7×10^{-4} at the 99% confidence limit.

The above experiments shows that there is no evidence of second class currents in nuclear beta-decay. The possibility of the existence of second - class current with $C_{IT}/ C_A \sim 10^{-4}$ $[m_oc^2]^{-1}$ or of a special kind cannot be excluded, but to find it would be extremely difficult at the present accuracy of experimental results and nuclear structure calculations.

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FIGURE CAPTIONS

Fig. 1. $(ft)^{+}/(ft)^{-}$ ratios of existing $\beta^{\pm}$ - transitions do not allow to determine the form of energy dependence. It seems to be independent of energy / full line /, but linear dependence cannot be excluded.

Fig. 2. K/β^{+} ratios. The experimental results are systematically below the full curve representing the effect of IT interaction.

Fig. 3. Experimental result on ^{8}Be excitation in the $\beta^{\pm}$ - decay of ^{8}Li and ^{8}B shows no evidence of second class current with $C_{IT}/|C_{A}| = -2 \times 10^{-3}\ [m_{o}c^{2}]^{-1}$ /full line/.

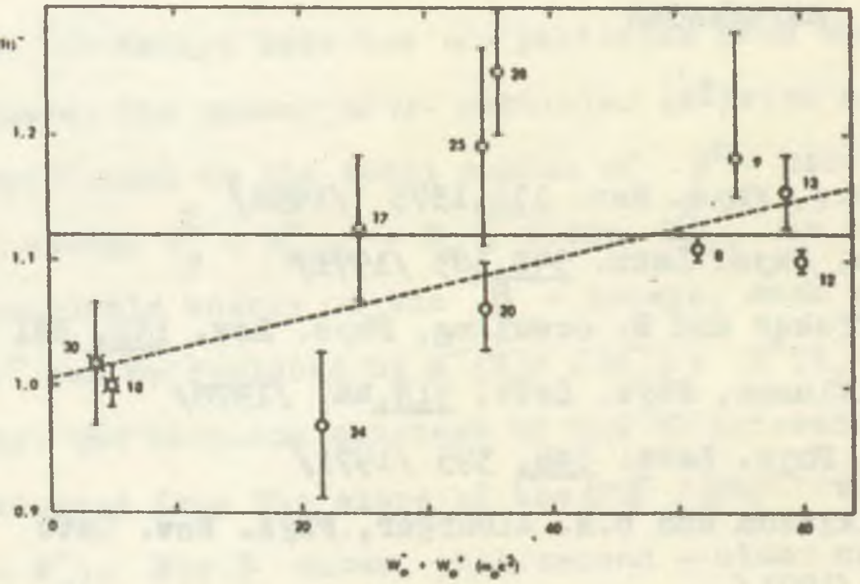

FIGURE 1.

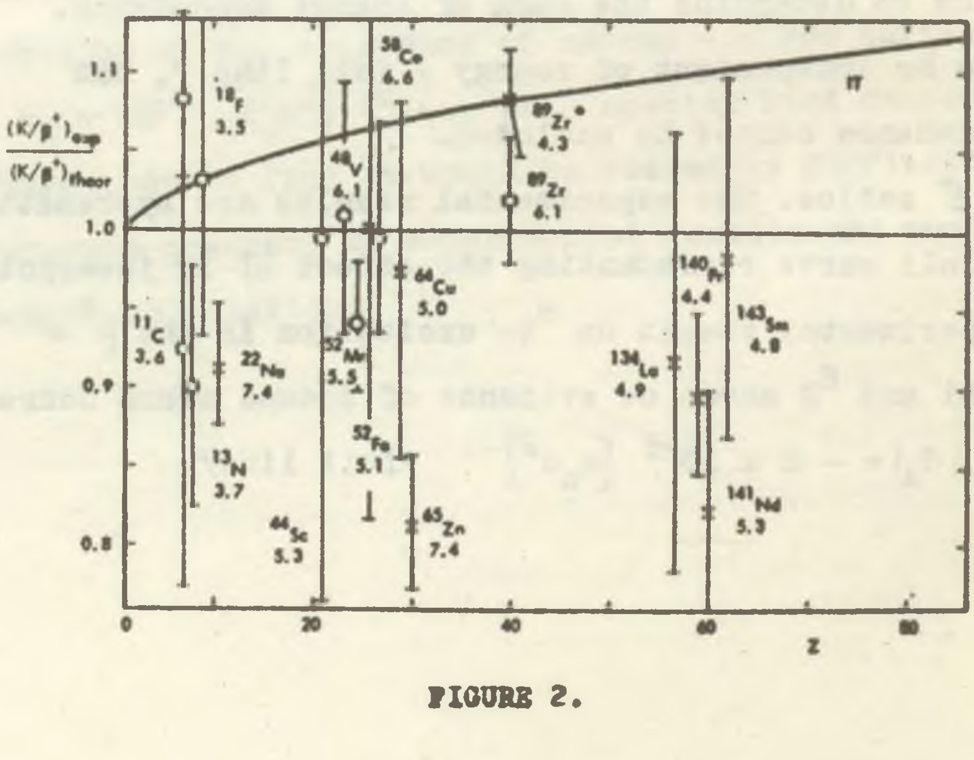

FIGURE 2.

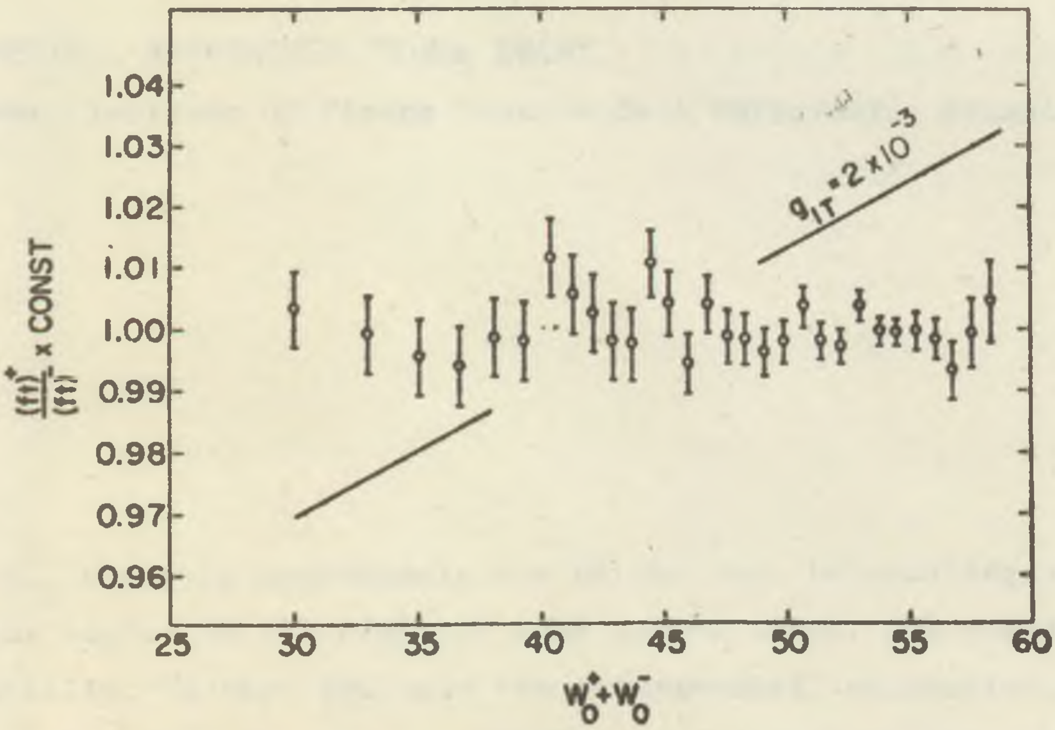
1.04
1.03
1.02
1.01
1.00
0.99
0.98
0.97
0.96
25
30
35
40
45
50
55
60
$\frac{(ft)^+}{(ft)^-}$ x CONST
$W_0^+ + W_0^-$
$g_{1T} = 2 \times 10^{-3}$

FIGURE 3.

THEORETICAL APPROACHES TO $K\ell_3$ DECAY

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Kl_3 decay is undoubtedly one of the most interesting and popular topics in the field of weak interactions. The reason, essentially, is that not only the experimental information on Kl_3 decay allows precise checks of some basic assumptions of the theory of weak interactions, but, also, it represents a sensible test of the theoretical approaches which are available. My talk will be devoted to a short survey on Kl_3 decay, looking at the problem from the point of view of current algebra, light-cone algebra and related techniques.

Let me start by briefly introducing the Kl_3 process. A very gratifying point is that the experimental data strongly recommend the simple Cabibbo description of semileptonic decays, based on V - A hadronic currents with octet behaviour under SU(3). Accordingly, the object of interest for the decay $K\to\pi\ell\tilde{\nu}_\ell$ is the matrix element

$$\langle\pi^0(q)|V_\mu^{(K^+)}|K^-(p)\rangle = \frac{1}{\sqrt{2}}\Big\{f_+(k^2)(p+q)_\mu + f_-(k^2)(p-q)_\mu\Big\}\ , \quad k = p-q \qquad (1)$$

where $V_\mu^{(K^\pm)} = V_\mu^{(4\pm i5)}$ is the vector part of the $\Delta S = \Delta Q$

hadron current, and $m_\ell^2 \leq k^2 \leq (m_K - m_\pi)^2$ in the physical decay region.

A third form factor is usually introduced :

$$f(k^2) = f_+(k^2) + \frac{k^2}{m_K^2 - m_\pi^2} f_-(k^2) \tag{2}$$

which is induced by the SU(3) breaking :

$$i \langle \pi | \partial_\mu V_\mu | K \rangle = (m_K^2 - m_\pi^2) f(k^2) \tag{3}$$

Rather than in terms of f_+, f_-, the process is better described by the form factors f_+, f, with which most of the theoretical predictions are concerned. It seems reasonable to assume a linear variation for all Kl_3 form factors :

$$\begin{aligned} f_\pm(k^2) &= f_\pm(0)\left(1 + \frac{\lambda_\pm}{m_\pi^2} k^2\right), \\ f(k^2) &= f(0)\left(1 + \frac{\lambda_0}{m_\pi^2} k^2\right), \\ \lambda_0 &= \lambda_+ + \xi(0) \frac{m_\pi^2}{m_K^2 - m_\pi^2}, \quad \xi(0) \equiv \frac{f_-(0)}{f_+(0)} \end{aligned} \tag{4}$$

Actually, the assumption of linear dependence for both f_+, f amounts to implicitly neglect λ_-, so that, within our parametrization (4), only $\xi(0)$ and λ_+ (or λ_0 and λ_+) are significant.

How looks our experimental knowledge of $\xi(0)$ and λ_+? The least we can say, to be optimistic, is that it is controversial. The best established case is $K^+_{\mu 3}$, where an overall fit gives [1]

$$\xi(0) = -0.85 \pm 0.20 \; ; \; \lambda_+ = 0.045 \pm 0.012 \tag{5}$$

K_{e3} data, on the other hand, seem to suggest smaller values of λ_+ (typically, 0.029 ± 0.006) so that, by combining $K_{\mu 3}$ and K_{e3} experiments we find [2]

$$\xi(0) = -0.65 \pm 0.20 \; ; \; \lambda_+ = 0.034 \pm 0.006 \tag{6}$$

Finally, the result $\xi(0) = 0.00 \pm 0.18$ can be found in the literature, which has been derived from the branching ratio $\Gamma(K^0_{\mu_3})/\Gamma(K^0_{e_3})$ assuming $\lambda_+ \approx 0.03.$ [1] Anyway, most of the data point towards a negative and sizeable value of $\xi(0)$; since $\xi \equiv 0$ in the SU(3) limit, this result does not seem, at first sight, easy to fit in the current picture of SU(3) as an approximate symmetry of the hadrons.

Now, the thoretical aim is to evaluate the form factors, or, at least, to get predictions on them.

A first possibility is represented by simple theoretical models, like e.g. K^* dominance of f_+, which gives $\lambda_+ = \frac{m_\pi^2}{m_{K^*}^2} \approx 0.024$, in rough agreement with most of the K_{e_3} data.

Alternatively, one can exploit general statements on the weak currents, like current algebra (and light-cone algebra), concerning their transformation and partial conservation properties.

SU(3), of course, is the first candidate, and the information which comes from the algebra of SU(3) charges alone is the well-known relation [3]

$$f(0) \equiv f_+(0) = 1 + O(\varepsilon_8^2) \tag{7}$$

which states that the departure of $f_+(0)$ from its SU(3) symmetric value is of second order in the breaking. This result, although of extreme importance, has not proved enough, so far, to produce a quantitative prediction on Kl_3 decay, since a reliable evaluation of the corrections $O(\varepsilon_8^2)$ is not at hand [4]; the only immediate indication which we can derive, assuming octet dominance, is the limitation $f_+(0) < 1$. Moreover, a relation of the kind in Eq.(7) is not easy to test, since only the product $f_+(0) \times \sin\theta_c$ is determined by experiment, θ_c being the Cabibbo angle.

We are lead therefore to extend our considerations from SU(3) to the SU(3) x SU(3) algebra generated by the vector and axial vector currents, supplemented by the PCAC hypothesis, namely, to soft-pion theory. Soft-pion theorems are, among the predictions of current algebra, the easiest to test experimentally and, as well known, they work in many cases so well that we are inclined to consider the smallness of the pion mass, which is the essence of the PCAC assumption, not as a dynamical accident but, more seriously, as the manifestation of an approximate chiral SU(2) x SU(2) symmetry of the hadron world, exactly realized in the limit $m_\pi = 0$. In our case, the soft-pion result is represented by the Callan-Treiman relation, which states that, in the limit of vanishing mass and momentum of the emitted pion :

$$f_+(m_K^2) + f_-(m_K^2) = \frac{f_K}{f_\pi} \tag{8}$$

or, in terms of the form factor f :

$$f(m_K^2) = \frac{f_K}{f_\pi} \tag{9}$$

Now, it happens that if we assume, according to PCAC, the Callan-Treiman result to hold without sensible variations on the pion mass-shell too, its simple linear extrapolation to the physical region turns out to be inconsistent with the experimental indication, since we find, from Eq.(4), inserting the experimental values of λ_+:

$$\xi(0) \simeq -0.32 \quad \text{and} \quad \xi(0) \approx -0.21,$$

in contrast with the figures (5) and (7).

This is not very pleasant, of course, if we believe that the soft-pion picture has some adherence to the physical world. Then, unless we want to invent some extra effects to explain this failure, the corrections to the PCAC result, due to the finite size of the pion mass, must be taken into account before trying a comparison with the experiment. The prescription to extrapolate soft-pion theorems onto the mass shell is not unique and depends, so to say, on the author's taste. I choose to review three among the most recent examples.

A very simple approach is represented by the rest-frame saturation of equal-time commutation relations of SU(3) x SU(3) charges and divergences, involving the isospin axial charges $\bar{Q}^{(\pi)}$ to which, essentially, the pion is associated [5], and such that, in the spirit of PCAC, $\dot{\bar{Q}}^{(\pi)} = \int d\vec{x}\, \partial_\mu A_\mu^{(\pi)} = O(m_\pi^2)$. The result is [6]

$$\left(1+\frac{m_\pi}{m_K}\right) f[(m_K - m_\pi)^2] = \frac{f_K}{f_\pi} + \frac{2}{m_K m_\pi f_\pi} \langle 0 | [\dot{\bar{Q}}^{(\pi^0)}, Q^{(K^+)}] | K^-(\vec{p}=0)\rangle - \frac{\sqrt{2}\, m_\pi}{m_K}\frac{1}{\pi}\int \frac{dq_0\, \rho(q_0, \vec{q}=0, \vec{p}=0)}{q_0 (q_0 - m_\pi)(q_0 - m_K)}, \tag{10}$$

where

$$\rho(q,p) = \sum_n \delta(q-p_n)\langle 0|\chi^{(\pi^0)}|n\rangle\langle n|\partial_\mu V_\mu^{(K^+)}|K^-(p)\rangle - \text{c.t.} = \int d^4x\, e^{iqx}\langle 0|[\chi^{(\pi^0)}(x), \partial_\mu V_\mu^{(K^+)}(0)]|K^-\rangle \equiv \rho(q^2, k^2), \tag{11}$$

$$k = p - q\ ; \quad (\Box + m_\pi^2)\, \partial_\mu A_\mu^{(\pi)}(x) = m_\pi^2 f_\pi\, \chi^{(\pi)}(x)$$

In shorthand notation

$$\left(1+\frac{m_\pi}{m_K}\right) f[(m_K - m_\pi)^2] = \frac{f_K}{f_\pi} + \delta(m_\pi) \tag{12}$$

It goes without saying that, for $m_\pi = 0$, the Callan-Treiman relation is reproduced. It is easily realized that the extrapolation is performed along a parabola in the (q^2, k^2) plane, represented in Fig. 1, connecting the soft-pion point $q^2 = 0,\ k^2 = m_K^2$ to the nearest point of the physical region, $q^2 = m_\pi^2,\ k^2 = (m_K - m_\pi)^2$: indeed, Eq. (10) is a dispersion relation where both the "mass" q^2 and the momentum transfer k^2 are varying along this line, whose intersections with the "lines of singularities" represent the various contributions to the dispersive sum.

Let us inspect very quickly the corrections to the Callan-Treiman relation, which appear in Eq. (12). The higher commutator $[\dot{\bar{Q}}, Q]$ is essential in order to get rid of the pion crossed mass singularity at $q^2 = m_\pi^2$, $k^2 = (m_K + m_\pi)^2$, proportional to $f[(m_K + m_\pi)^2]$, whose appearance would prevent us the possibility to extrapolate the Callan-Treiman relation to a single physical point. Such a commutator is, of course, outside the current algebra frame; so, this is the ideal case where to test models of chiral SU(3) x SU(3) symmetry breaking. We have, moreover, kinematical corrections (which are seen to work in the right direction but to be not enough) and a dispersive continuum which might be, as a matter of fact, an $O(m_\pi \ln m_\pi)$, since the integration threshold is $O(m_\pi)$. A quantitative discussion of the corrections is not easy: opinions are diverging, and so far, a definite conclusion has not been reached. For example, on the basis of simple, model dependent estimates, Banerjee[7] has claimed that the Callan-Treiman relation(namely, approximate chiral SU(2) x SU(2) symmetry) and the experimental data can be compatible: a dip of $f(k^2)$ somewhere around $k^2 = (m_K - m_\pi)^2$ is required. This conclusion, which has received support from several other authors, is tackled by the opponent party, trying to interpret data rather in terms of near SU(3). Anyway, a good experimental information on $f(k^2)$ near the edge of the physical region would be of much help in clarifying things. The situation concerning $f(k^2)$ is sketched, qualitatively, in Fig. 2.

The second method which I would like to mention starts from the SU(3) x SU(3) light-cone algebra, and has been developed by G. Furlan and myself[8]. Light-cone physics has represented one of the most exciting topics of these recent times and shows nice features of simplicity. Being not willing to enter details, I limit to remind that it seems reasonable to abstract the light-cone current commutation relations from the free quark model[9]. One of the arguments in favour of this assumption is that this model has been already successful in suggesting the SU(3) x SU(3) equal-time algebra of charges (and, perhaps, of divergences).

As a result, introducing the operators $\left(x^{\pm} = \frac{1}{\sqrt{2}}(x^{0} \pm x^{3}),\right.$ $\left.\vec{x}_{\perp} \equiv (x^{1}, x^{2})\right)$:

$$\begin{aligned} Q_a(q, x^+) &= \int dx^- d\vec{x}_{\perp} e^{i(q^+x^- - \vec{q}_{\perp}\vec{x}_{\perp})} V_a^+(x) \\ \bar{Q}_a(q, x^+) &= \int dx^- d\vec{x}_{\perp} e^{i(q^+x^- - \vec{q}_{\perp}\vec{x}_{\perp})} A_a^+(x) \end{aligned} \tag{13}$$

we have at disposal the SU(3) x SU(3) light-cone algebra

$$\begin{aligned} \left[Q_a(q, x^+), Q_b(k, x^+)\right] &= i f_{abc} Q_c(q+k, x^+) \\ \left[Q_a(q, x^+), \bar{Q}_b(k, x^+)\right] &= i f_{abc} \bar{Q}_c(q+k, x^+) \\ \left[\bar{Q}_a(q, x^+), \bar{Q}_b(k, x^+)\right] &= i f_{abc} Q_c(q+k, x^+) \end{aligned} \tag{14}$$

The importance of defining the operators $Q_a, \bar{Q}_a$ in Eq. (13) stems from the fact that at $q^+ = \vec{q}_{\perp} = 0$ they coincide, for conserved currents, with the familiar SU(3) x SU(3) charges $Q = \int d\vec{x} J^0(x)$ and therefore, while being appropriate to light-cone kinematics, they are of use in dealing with partial conservation properties when the symmetry is broken.

The procedure, now, very similar to the previous one, consists in saturating the simple algebraic structure (14). The result which is derived for Kl_3 decay is the modified version of Eq. (12):

$$\begin{aligned} &\frac{1}{2}\left(1 + \frac{m_\pi}{m_K}\right) f\left[(m_K - m_\pi)^2\right] + \tilde{\delta}(m_\pi) = \\ &\qquad = \frac{f_K}{f_\pi} - \frac{f_K}{2 f_\pi(0)} \end{aligned} \tag{15}$$

where $f_\pi(0) = f_\pi(m_\pi = 0)$ is the value of f_π given by the Goldberger-Treiman relation, exactly valid for massless pions, and $\tilde{\delta}(m_\pi)$, whose explicit expression this time I omit, represents, again, a dispersive continuum of order m_π. Having started from a light-cone commutator, the dispersive line is, in this case, the straight line, represented in Fig. 1, of equation

$$q^2 = \frac{m_\pi}{m_\pi - m_K} k^2 + m_\pi m_K \tag{16}$$

Actually, the straight line can be obtained from the parabola, characteristic of the equal-time method, as a degenerate case, by taking an appropriate limit which requires, for consistency, the convergence of form factors at infinity. The straight line, moreover, does not touch the soft-pion point $q^2 = 0,\ k^2 = m_K^2$, which means that in the present method we get the separation of the result into a finite part plus $\tilde{\delta}(m_\pi)$ corrections only after that the soft-pion limit has been imposed, in a sense, like a constraint; this is why $f_\pi(0)$ appears in Eq.(15). Apart from this difficulty in handling the soft-pion limit $m_\pi = 0$, the fact of working along a straight line is really an advantage, since the lines of singularities are met only once (and not twice), and this results in a remarkable simplification in the set of graphs contributing to $\tilde{\delta}(m_\pi)$. Looking at Fig. 1 we see that crossed mass singularities are excluded; as a consequence, the pion crossed contribution, which I mentioned before, is disposed of from the beginning, and we directly get a prediction at a single physical point without introducing higher, model dependent commutators. Furthermore, the q^2-channel (lines $q^2 = M_n^2$) should be strongly depressed, since form factors are evaluated at large (timelike) momentum transfer, of order $1/m_\pi$, while simplifications are possible in the k^2-channel (lines $k^2 = M_m^2$), where momentum transfers are $O(m_\pi)$. Now, to be quantitative, we could try a comparison of Eq.(17) with experimental data, neglecting, as a first approximation, the $\tilde{\delta}(m_\pi)$ effects and exploiting only kinematical coefficients, corrections to the Goldberger-Treiman relation included. The result is that the figures (5) and (6) are pretty well reproduced since we find, according to the alternative values of λ_+ :

$$\xi(0) \approx -0.75 \qquad \lambda_+ = 0.045$$
$$\xi(0) \approx -0.65 \qquad \lambda_+ = 0.034 \tag{17}$$

This is not very conclusive, of course, and the indication coming from a simple-minded estimate is that the $\tilde{\delta}(m_\pi)$ effects might be sizeable, although not dramatic. So, let us consider the figures (17), at least, as an encouraging indication.

An alternative and interesting light-cone approach to Kl_3 decay has been proposed by Brandt and Preparata [10], who find, among other things

$$f_+(0) \approx 0.94 \ ; \ \xi(0) \approx -0.73 \tag{18}$$

In their method light-cone commutators are used in connection with appropriate finite-energy sum rules in the "mass" variable q^2 and play, essentially, a role analogous to that of Regge poles in conventional dispersion relations. The dispersive lines are, in this case, referring to Fig. 1, the vertical lines $k^2 = m_\pi^2$ and $k^2 = m_K^2$. It is very difficult to establish connections (if any) with our approach. We think that our method is more economical and involves, in practice, very little of the light-cone approach, since the simple "charge-charge" algebra (14) is taken as the starting point. Brandt and Preparata, on the other hand, use higher, model dependent commutators, which they abstract from the simple model of Gell-Mann, Oakes and Renner [11]; in their treatment, therefore, the bilocal operators appear to play the fundamental role, and more results are obtained, as a test of more assumptions.

Acknowledgements

I would like to thank Prof. G. Furlan for helpful discussions.

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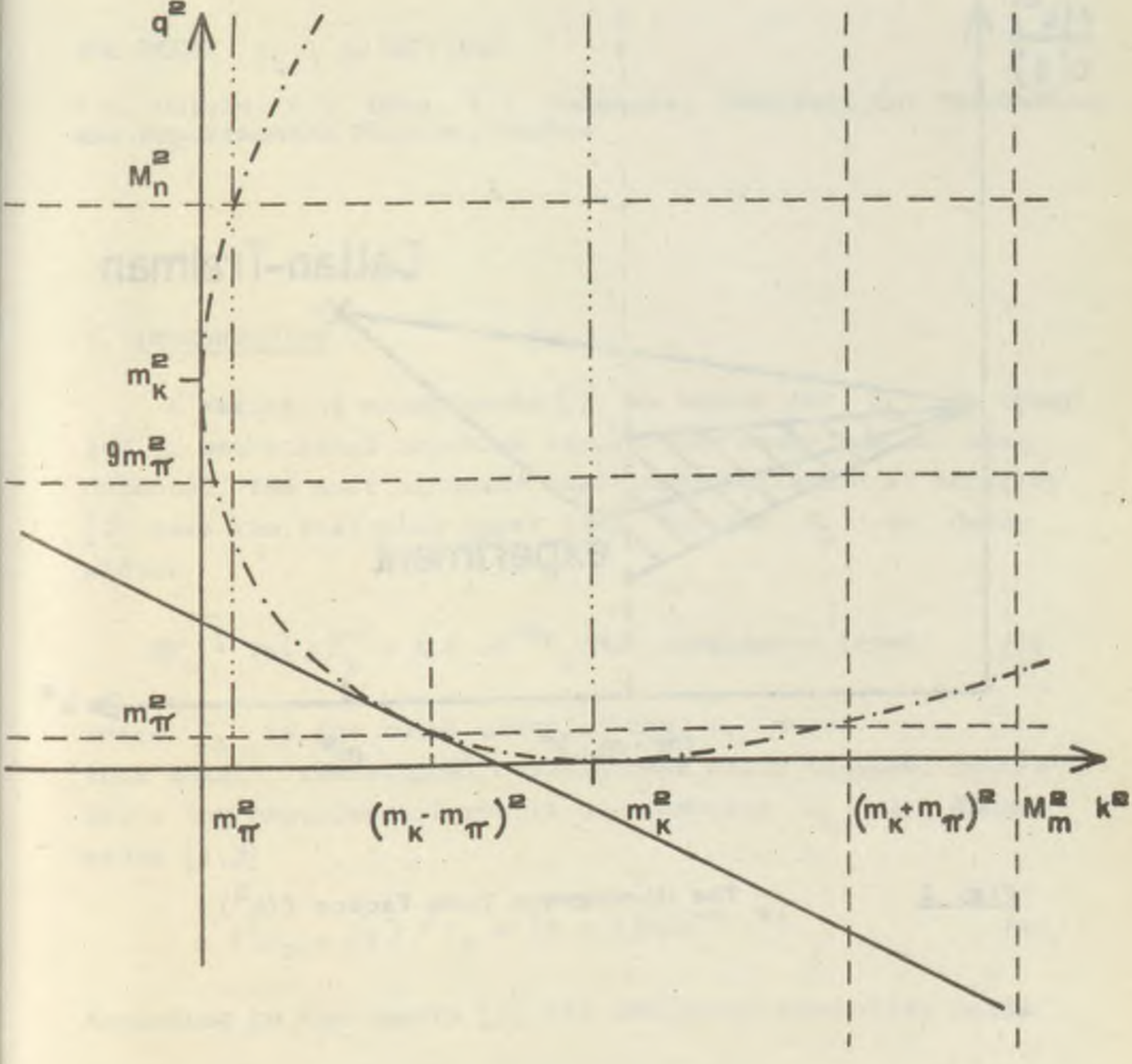

Fig. 1 — · — · — equal-time parabola

——— light-cone path Eq. (16)

– – – – lines of singularities $q^2 = M_n^2$, $k^2 = M_m^2$.

– · · – · · – lines of Ref. (10)

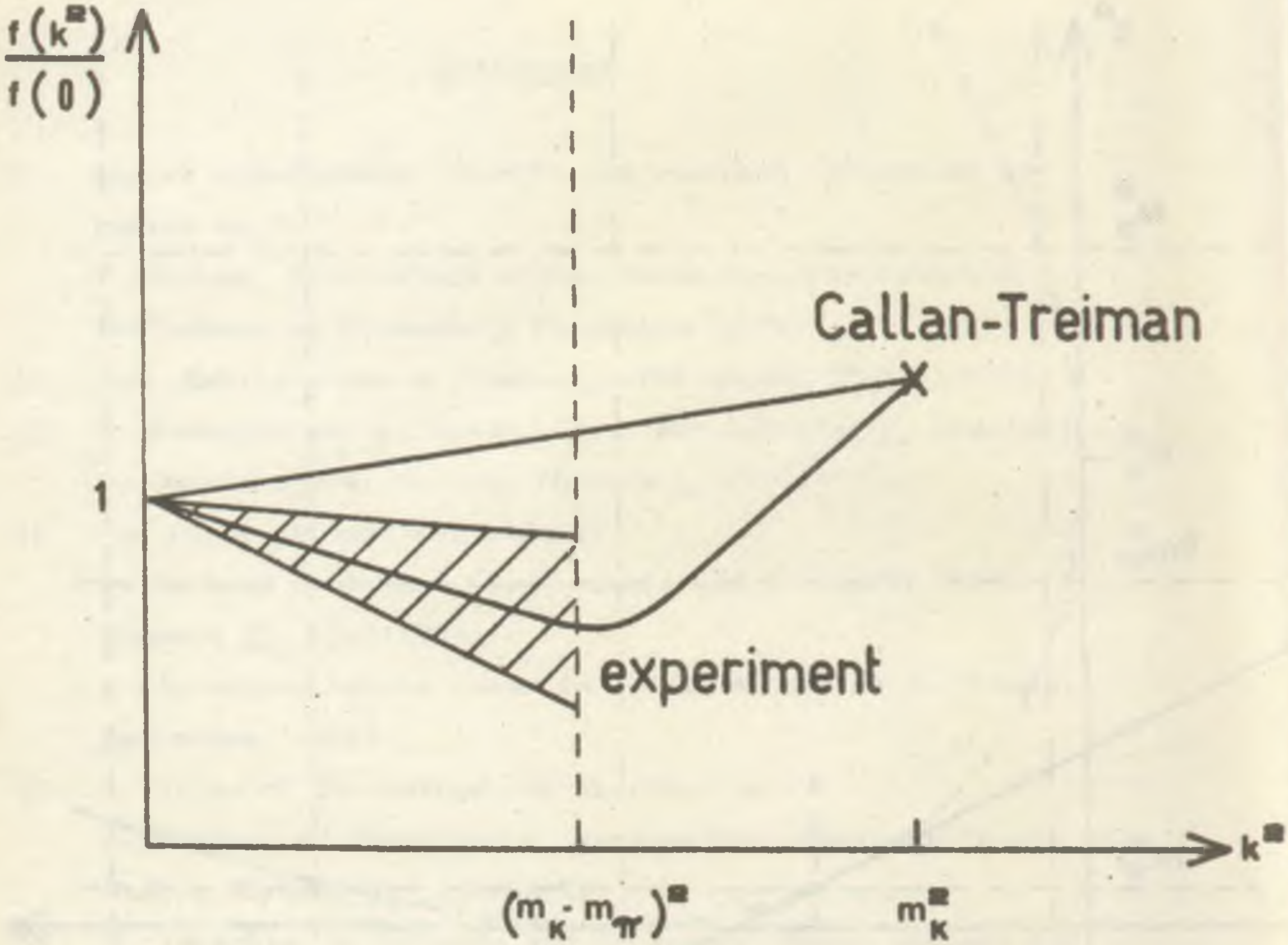

Fig. 2 The divergence form factor $f(k^2)$.

THE DECAY $K_L \to 2\mu$ (REVIEW)

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I. INTRODUCTION

A series of experiments [1] on search for $K_L \to 2\mu$ decay led to sensational negative result: the decay has not been detected. The most accurate experiment performed at Berkeley [2] gave the following upper limit for the $K_L \to 2\mu$ decay width:

$$\Gamma(K_L \to 2\mu) \leq \bar{\Gamma}_L^{\mu} = 1.8\cdot 10^{-9}\Gamma_L \text{ /90\% confidence level/} \qquad /1/$$

where Γ_L is the total width of the K_L meson.
This result, theoretical calculations being trusted, contradicts the experimental result on measuring $K_L \to 2\gamma$ decay width [1,3]

$$\Gamma(K_L \to 2\gamma)/\Gamma_L = (5 \pm 1)\cdot 10^{-4} \;{}^{*/} \qquad /2/$$

According to the theory [4] the following inequality holds

$$\Gamma(K_L \to 2\mu)/\Gamma(K_L \to 2\gamma) \geqslant 1.2\cdot 10^{-5} \qquad /3/$$

and from eq. /2/ we have the lower bound

*/Tables [1] give the ratio $(5.6\pm 0.5)\cdot 10^{-4}$. According to the paper [3e] the world average value is $(5+0.5)\cdot 10^{-4}$. The result by the ITEP group [3d], $\Gamma(K_L \to 2\gamma)/\Gamma_L = (4.6\pm 0.9)\cdot 10^{-4}$ was not taken into account in Tables [1] and in paper [3e].

$$\Gamma\left(K_L \to 2\mu\right)/\Gamma_L \geqslant \left(6 \pm 1.2\right)\cdot 10^{-9} \qquad /4/$$

which is in apparent contradiction with /1/.

1.1. Are the Experiments Reliable?

Experts assert the experiment [2] to be reliable. In many respects it is a unique one. Thus, for instance, as a background for $K_L \to 2\mu$ decay in this experiment served a rare decay $K_L \to 2\pi$, about million events of the latter decay having been detected. Nevertheless one cannot consider the experimental result to be a final one basing on a single experiment, even a very good one. The history of elementary particle physics during last 15 years warns against this. It is sufficient to recall the statements on tensor coupling in β decay of He^6, on the absence up to 10^{-5} of $\pi \to e\nu$ decay, on $\Delta Q = \Delta S$ rule violation in K^0 meson decays, on the spectrum of electrons with $\rho = 0$ in the muon decay, on the charge asymmetry in $\eta \to \pi^+\pi^-\pi^0$ decay, on the violation of $\eta_{00} = \eta_{+-}$ equality and some other examples when experiments which seemed to be reliable turned out to be incorrect. Thus new experiments on search for the $K_L \to 2\mu$ decay seem necessary.

As for the result /2/, it is an average of several experiments [3] performed with different technique and being in agreement with each other. Nevertheless, it would be desirable to determine the $K_L \to 2\gamma$ decay width with much higher accuracy.

Bearing in mind that the experiments under discussion /and especially, the search for $K_L \to 2\mu$ decay/ require good skill, great efforts and long time it is quite necessary to analyze theoretical aspects of the problem. /Note that if both the theory [4] and the experiment on $K_L \to 2\gamma$ decay are correct the probability to obtain experimentally the result /1/ for the decay $K_L \to 2\mu$ is $\left(2\div3\right)\cdot 10^{-3}$/.

1.2. Is the Theory Reliable?

This question has been discussed recently both in number of original papers /we shall consider them in detail below/ and in some reviews [5].

The derivation of the theoretical lower bound /3/ on the $K_L \to 2\mu$ decay width seems to be quite strustworthy. Just for this reason the experimental result [2] is so exciting.

In the next section we will discuss in detail the derivation of relation /3/. Here let us point out only that the main idea is to calculate the imaginary part of the amplitude. The account of the real part may only increase the value of the decay probability.

The imaginary /absorptive/ part emerges from transition on the mass shell $K_L \to n \to 2\mu$ where (n) is some real intermediate state. The main contribution into imaginary part is given by the two-photon intermediate state /see fig.1/. This contribution can be calculated unambiguously if the $K_L \to 2\gamma$ decay amplitude is known experimentally and the transition $2\gamma \to 2\mu$ is described by quantum electrodynamics. The phase space of other states /e.g. 3π, $2\pi\gamma$/ is much less. For this reason one may think that the contribution of the twophoton state dominates the imaginary part of the amplitude. Then inequality /3/ is valid.

Theoretical foundation of relation /3/ is, if one may say so, three-fold. Firstly, the general principles of physics such as unitarity of the S-matrix and CPT theorem, are assumed to be valid. Basing only on these principles, one can express imaginary part of the amplitude of $K_L \to 2\mu$ decay through the amplitudes of the transitions on the mass shell.

Secondly, the hypothesis is made that there is no "leakage" of K-mesons into unknown decay channels, i.e. there are no new particles with the mass less than /or approximately equal to/ the K-meson mass. The contribution of new decay channels if they exist would change theoretical bound on $K_L \to 2\mu$ decay rate.

And finally it is assumed that the known particles produced in K_L decay possess only usual interactions and that the estimates of imaginary parts arising from various intermediate states are not wrong by several orders of magnitude. If for example π-mesons interact strongly with muons the neglect of 3π intermediate state would be unjustifiable.

1.3. Organization of the Paper

In Section 2 we give a detailed derivation of relation /3/.

From discussion in the preceding Section one infers that if experimental and theoretical value of the $K_L \to 2\mu$ decay probability are in real contradiction to each other our concepts on elementary particles and their interactions should be modified in some essential points. In accordance with three layers of theoretical foundation of inequality /3/ these possible modifications may be grouped in the following way:

1/ introduction of new interactions of known particles /see Sec. 3/,

2/ introduction of new particles /Sec. 4/,

3/ violation of fundamental priciples of physics /Sec. 5/.

It is obvious that the existence of new particles or new interactions would manifest itself not only in $K_L \to 2\mu$ decay but in some other phenomena. The discussion of these manifestations is one of the main purposes of this short review.

II. CONSERVATIVE THEORY OF $K_L \to \mu^+\mu^-$ DECAY

2.1. Unitarity Condition

The lower bound /3/ on the $K_L \to \mu^+\mu^-$ decay width follows from the unitarity condition which allows one to express the imaginary part of the amplitude in terms of the matrix elements of transitions $K_L \to n$ and $n \to 2\mu$ where n is some real intermediate state. It is worth emphasizing that in particular case of $K_L \to 2\mu$ decay CPT invariance alone is sufficient to calculate the imaginary part while usually more restrictive condition of T-invariance is needed. Let us review the argument.

If we introduce in a standard way S and T matrices

$$S = 1 + iT \qquad /5/$$

then the unitarity condition

$$SS^+ = 1 \qquad /6/$$

can be written as

$$T_i^f - \left(T_f^i\right)^* = i \sum_n \int d\tau_n T_i^n \left(T_f^n\right)^* \qquad /7/$$

where T_i^f is the transition amplitude from state i into state f and the sum extends over all possible real intermediate states n, τ_n being the corresponding phase space. The asterisk in eq. /7/ denotes the complex conjugation of the amplitude /for normalization of the amplitudes see Appendix I eq./1.5/ and /1.6//.

In the case of $K_L \to 2\mu$ the initial state is

$$K_L = K_2 + \varepsilon K_1 \qquad /8/$$

where K_2 and K_1 have definite CP parity: -1 and +1, respectively. Due to smallness of ε $\left(|\varepsilon| = 2.10^{-3}\right)$ we neglect the term εK_i in this Section and discuss it later.

As a result of $K_L \to 2\mu$ decay the muonic pair with the total angular momentum equal to zero is produced. There are two possible states of this pair ^1S_o and ^3P_o which have CP parity equal to -1 and +1 and which are denoted in the following by indices μ_- and μ_+, respectively. If muon polarization is not fixed, these states do not interfere with each other.

Let us consider the decay of K_2 meson into one of these states. It follows from CPT invariance that

$$T_f^i = T_{\tilde{i}}^{\tilde{f}} \qquad /9/$$

where state $\tilde{i}$ is obtained /up to a phase factor/ from state i by CP conjugation. Since the initial and final states are eigenstates of CP conjugation operator, the CPT transformation in this case is reduced /up to a phase factor/ to T-transformation and we have

$$T_f^i = \pm T_i^f \qquad /10/$$

The amplitudes T_f^i can always be defined so that the unitarity relation takes the same form irrespectively of the sign in the right-hand part of eq. /10/

$$\mathrm{Im}\, T_2^{\mu\pm} = \frac{1}{2} \int d\tau_n\, T_2^n \left(T_{\mu_\pm}^n\right)^* \qquad /11/$$

where index 2 means the state K_2^o.

2.2. Two-Photon Imaginary Part

As was already mentioned, the unitarity sum /11/ is dominated by two-photon intermediate state /see Fig. 1/. Corresponding contribution into imaginary part is calculated in

Appendix I. If we keep only this contribution and assume that CP is conserved in $K_2 \to 2\gamma$ decay, then

$$\frac{\Gamma(K_2 \to 2\mu)}{\Gamma(K_2 \to 2\gamma)} = \frac{\alpha^2}{2v}\left(\frac{m_\mu}{m_K}\right)^2 \ell n^2\ \frac{1+v}{1-v} = 1.2\cdot 10^{-5} \qquad /12/$$

where v is the velocity of muon in the rest frame of K-meson $(v \approx 0.9)$.

If the amplitude of $K_2 \to 2\gamma$ decay is totally CP-non-conserving, the lower bound for the ratio of the rates of the decays $K_2 \to 2\mu$ and $K_2 \to 2\gamma$ will be the same as for the K_1 decays with CP conservation [6].

$$\frac{\Gamma(K_1 \to 2\mu)}{\Gamma(K_1 \to 2\gamma)} = \frac{\alpha^2 v}{2}\left(\frac{m_\mu}{m_K}\right)^2 \ell n^2\ \frac{1+v}{1-v} = 1\cdot 10^{-5} \qquad /13/$$

2.3. <u>Other Imaginary Parts</u>

The states $2\pi\gamma$, 3π, $3\pi\gamma$, give the contribution of the same oreder in weak and electromagnetic coupling constants as the two-photon state. However their phase space is negligibly small as compared with that of the two photons and one may think that the account of these states changes the result inessentially. Some quantitative estimates are given in Table I. As to these estimates some remarks are worth making:

<u>State $2\pi\gamma$ /Fig. 2/</u>. The calculation here [6] is a bit less definite than in the case of the two-photon absorptive part. If one neglects the form factor of π and form factors in $K_2 \to 2\pi\gamma$ decay the result is expressed through the $K_2 \to 2\gamma$ decay width. The account of form factors can change the result approximately by a factor of two [7, 8]. At present only the upper limit for the $K_L \to 2\pi\gamma$ decay width is known experimentally. The corresponding upper limit for the $2\pi\gamma$ contribution which is given in Table I is taken from paper [8] and is approximately by an order of magnitude less than that obtained in paper [7].

State 3π /Fig. 3/. The calculation here is essentially model dependent. The amplitude of transition $3\pi - 2\gamma$ can be found in framework of the current algebra [9,10]. The value 10^{-5} for the ratio of 3π and 2γ contributions is obtained [10] on the basis of these calculations and dispersion relations. Simple order of the magnitude estimate /accounting for the ratio of phase space of 3π and 2γ states/ gives 10^{-4} for this ratio.

2.4. Terms of Order ε.

Up to now we have neglected the term εK_1 in formula /8/ assuming that $K_L = K_2$. Let us now estimate the contribution from this term. Unfortunately, this estimate will be extremely unreliable because the decays $K_S \to 2\mu$ have not been observed experimentally and thus the amplitudes of transition $K_1 \to 2\mu$ and $K_1 \to 2\gamma$ are not known. Let us recall that if CPT is conserved then

$$K_S = K_1 + \varepsilon K_2 \qquad /14/$$

The experimental upper bounds are

$$\Gamma\left(K_S \to 2\mu\right) \leqslant \overline{\Gamma}^{\mu}_S \overset{[1]}{=} 7\cdot 10^{-6}\ \Gamma_S \qquad /15/$$

/see also below formula /25// and

$$\Gamma\left(K_S \to 2\gamma\right) \leqslant \overline{\Gamma}^{\gamma}_S \overset{[5]}{=} 12\cdot 10^{-3}\ \Gamma_S \qquad /16/$$

If we accept conservative hypotheses

$$\Gamma\left(K_S \to 2\mu\right) \sim \Gamma_{\text{theor.}}\left(K_L \to 2\mu\right) \sim 10^{-8}\Gamma_L \sim 10^{-11}\Gamma_S \qquad /17/$$

$$\Gamma\left(K_S \to 2\gamma\right) \sim \Gamma\left(K_L \to 2\gamma\right) \sim 10^{-6}\Gamma_S \qquad /18/$$

then the terms proportional to ε will contribute $\sim 10^{-3}$ to the amplitude $K_L \to 2\mu$. Note that the phase of this contribution is not only due to the phase of ε /experimentally

$\varepsilon \approx 2.10^3 e^{i\pi/4}$/ but also due to the absorptive part which arises from the contribution of the real two-pion intermediate state.

2.5. Terms of Order G^2

At first the interest in the $K_L \to 2\mu$ decay was arisen by the search for weak neutral currents. Even if the Lagrangian of weak interaction does not contain neutral currents they may and generally speaking must appear in the second order in the weak interaction constant /see Fig. 5/. The amplitude of Fig. 5 quadratically diverges and according to paper [11] the result can be written as

$$\frac{\Gamma(K_L \quad 2\mu)}{\Gamma(K^+ \to \mu\nu)} = G^2 \left(\frac{\Lambda}{4\pi}\right)^4 x \begin{cases} 2 \text{ in the theory with intermediate W-boson} \\ 32 \text{ in the four-fermion theory} \end{cases}$$

Here $G = 10^{-5}/\, m_p^2$ and Λ is the cut off.

It should be emphasized that if general principles such as unitarity of S-matrix and CPT invariance are valid, the amplitude corresponding to Fig. 5 is purely real and thus cannot cancel the two-photon absorptive part /Fig. 1/. The absorptive part in order G^2 appears due to the real intermediate $\pi\mu\nu$ state /see the graph of Fig. 6/. However, the contribution from this state is at least five orders of magnitude less than the contribution of 2γ state.

2.6. Breaking of Conservative Theory

We have seen that as compared to the contribution from the state 2γ the contribution from all other intermediate states is small. However the above estimates use theoretical values of the coupling constants of K and π mesons, photons, and muons. At present the experiment allows in a number of cases essentially larger values of these coupling constants. Therefore it cannot be ruled out that the contribution of some channels considerably exceeds the conservative estimates. We have already mentioned other possibilities for breaking of the conservative theory: "leakage" of K_L-mesons into yet unobserved decay channels and violation of fundamental principles.

In all the explanations of the $K_L \to 2\mu$ puzzle proposed so far an "accidental" cancellation of the two-photon imaginary part by some mechanism is introduced. Let us notice one trivial numerical fact. To remove the discrepancy between the theory and experiment it is sufficient to cancel not the whole two-photon imaginary part but 0.45 of its magnitude. If the partial width of $K_L \to 2\gamma$ decay is not 5.10^{-4} but 4.10^{-4}, this figure being in no contradiction with any particular experiment on measuring of $\Gamma(K_L \to 2\gamma)$ [3], then the fraction of cancellation term can be decreased from 0.45 to 0.35.

III. NEW INTERACTIONS?

3.1. Anomalous $K_1\bar{\mu}\mu$ Interaction

Crist and Lee [12] /see also [13]/ suggested that the observed suppression of the $K_L \to 2\mu$ decay is due to the small admixture of K_1 in the wave function of K_L meson

$$K_L = K_2 + \varepsilon K_1 \qquad /19/$$

For this to be the case two conditions must be satisfied. Firstly, to compensate the smallness of ε, the amplitude of the $K_1 \to 2\mu$ decay must be much larger than the amplitude of $K_2 \to 2\mu$ decay. Secondly, in $K^o \to 2\mu$ decays CP must be strongly violated. Otherwise final states in $K_1 \to 2\mu$ and $K_2 \to 2\mu$ decays are different and do not interfere, so the cancellation of the two-photon contribution is impossible.

CP Violation in $K_S \to 2\mu$ Decay

Let us first suppose that there exists CP odd interaction of the form [14, 15]:

$$if K_1 \bar{\mu} \gamma_5 \mu \qquad /20/$$

and the value of constant f is such that for the amplitude $\tilde{T}_1^{\mu}$ the following inequality is fulfilled:

$$\frac{m_K v}{8\pi} \left| \tilde{T}_2^{\mu-} + \varepsilon \tilde{T}_1^{\mu-} \right|^2 < \bar{\Gamma}_L^{\mu} \qquad /21/$$

/see Appendix I, eqs. /1.3/ and/1.5//.

By comparing eqs. /21/ and /12/ one comes to inequality

$$\left(\sqrt{1.2\cdot 10^{-5}\Gamma_L^{\gamma}} + \sqrt{\bar{\Gamma}_L^{\mu}}\right)^2 \geqslant |\mathrm{Im}\varepsilon|^2 \left|\tilde{T}_1^{\mu-}\right|^2 \frac{m_k V}{8\pi} \geqslant \left(\sqrt{1.2\cdot 10^{-5}\Gamma_L^{\gamma}} - \sqrt{\bar{\Gamma}_L^{\mu}}\right)^2$$

This relation implies in turn that the $K_S \to 2\mu$ decay rate is large and lies within the limits

$$1.2\cdot 10^{-5} \geqslant \Gamma(K_S \to 2\mu)/\Gamma_S \geqslant 1\cdot 10^{-6} \qquad /22/$$

It is worth noting that the width of hypothetical CP violating decay $K_S \to 2\mu$ turns out to be of the same order of magnitude as that of the known CP violating decay $K_L \to 2\pi$. Thus the interaction /20/ has a natural order of magnitude. The CP odd interaction of neutral currents of such strength was discussed several years ago by Lipmanov [16].

Let us now consider some more exotic possibilities.

CP Violation in $K_L \to 2\mu$ Decay

Let CP be conserved in $K_S \to 2\mu$ decay /the amplitude of which as in the preceding section is considered to be much larger than the amplitude $T(K_L \to 2\mu)$/. Let us further assume that the amplitude of $K_2 \to 2\gamma$ decay is CP-odd. Then the two-photon imaginary part is smaller /see sec. 2.2./ and the amplitude must satisfy weaker inequalitites

$$\left(\sqrt{1\cdot 10^{-5}\Gamma_L^{\gamma}}+\sqrt{\Gamma_L^{\mu}}\right)^2 \geq |\mathrm{Im}\varepsilon|^2 |\tilde{T}_1^{\mu+}|^2 \frac{m_K V}{8\pi} \geq \left(\sqrt{1\cdot 10^{-5}\Gamma_L^{\gamma}}-\sqrt{\Gamma_L^{\mu}}\right)^2 \qquad /23/$$

The probability of $K_S \to 2\mu$ decay in this case lies within the limits

$$10\cdot 10^{-7} \geq \Gamma(K_S \to 2\mu)/\Gamma_S \geq 6\cdot 10^{-7} \qquad /24/$$

Professor Kleinknecht has informed one of the authors on the preliminary result of the CERN group

$$\Gamma(K_S \to 2\mu) \leq \Gamma_S^{\mu} = 1.5\cdot 10^{-6} \ /90\%\ \text{c.l.}/ \qquad /25/$$

If this bound is improved by the factor of 3 it will be possible to rule out the possibilities /22/ and /24/. When comparing the theory with the experiment one should keep in mind however the uncertainty in theoretical predictions arising from experimental error in $K_L \to 2\gamma$ decay width. If $\Gamma(K_L \to 2\gamma)/\Gamma_L = 4.10^{-4}$ then the lower bounds in relations /22/ and /24/ decrease to 6.10^{-7} and 3.10^{-7}, respectively.

3.3. Anomalous $2\pi - 2\mu$ Interaction [17,18,19]

In the preceding section we assumed that $K_1 \to 2\mu$ decay is due to direct coupling of K_1 mesons and muons. Then the phase of the amplitude is fixed by hermiticity of the effective Hamiltonian /or, more strictly, by the unitarity of S-matrix and CPT invariance/. It is clear, however that the amplitude of $K_2 \to 2\mu$ transition is maximally cancelled if matrix element T_2^μ and the product εT_1^μ have opposite phases. If the amplitude T_2^μ is purely imaginary then the lowest bound for $\Gamma(K_S \to 2\mu)$ is achieved when the phase of T_1^μ is equal to $\pi/4$ or $5\pi/4$. /Remind that $\varphi_\varepsilon = \pi/4$/. It is obvious that in this case the bound /22/ or /24/ is decreased by a factor of 2 and becomes

$$\Gamma(K_S \to 2\mu) \,/\, \Gamma_S \geqslant 5\cdot 10^{-7} \qquad /22'/$$

if CP is violated in $K_1 \to 2\mu$ decay;

$$\Gamma(K_S \to 2\mu) \,/\, \Gamma_S \geqslant 3\cdot 10^{-7} \qquad /24'/$$

if CP is violated in $K_2 \to 2\gamma$ decay.

The phase of the amplitude $K_S \to 2\mu$ could be large if strong $2\pi - 2\mu$ interaction would exist. If the amplitude of $K_2 \to 2\gamma$ decay is CP even then the anomalous $2\pi - 2\mu$ interaction must violate CP invariance and has the following effective Lagrangian

$$\frac{C_-}{m_K}\left(\vec{\varphi}_\pi \vec{\varphi}_\pi^{\,*}\right)\left(\bar{\mu}\gamma_5\mu\right) \qquad /26/$$

If CP is violated in $K_2 \to 2\gamma$ decay then the anomalous $2\pi - 2\mu$ interaction may be CP invariant. In this case the effective Lagrangian of this interaction has the form

$$\frac{C_+}{m_K}\left(\vec{\varphi}_\pi \vec{\varphi}_\pi^{\,*}\right)\left(\bar{\mu}\mu\right) \qquad /27/$$

The latter possibility was discussed in papers [17,18,19]. Starting from the requirement that the phase of the amplitude $T_S^{\mu\pm}$ which we denote as $\varphi_S^{\mu\pm}$ is close to $\pi/4$, it can be easily obtained with the help of the unitarity relation that

$$C_-^2 = \frac{256\pi^2 \sin^2\varphi_S^{\mu-}}{eV_\pi V_\mu}\;\frac{\Gamma_S^\mu}{\Gamma_S} \approx 3\cdot 10^{-4} \qquad /28/$$

$$C_{+}^{2} = \frac{256\pi^{2}\,\sin^{2}\varphi_{S}^{\mu+}}{3V_{\pi}V_{\mu}^{3}}\;\frac{\Gamma_{S}^{\mu}}{\Gamma_{S}} \approx 2\cdot 10^{-4} \qquad /29/$$

where for the ratio Γ_S^μ/Γ_S we used the corresponding lower bounds /22'/ and /24'/.

The obtained value of the 2π-2μ interaction constant $C_{\pm} \sim 10^{-2}$ is much larger than that arising from the two-photon mechanism in perturbation theory [6]. Nevertheless in experiments with muons performed up to now /measurement of muon anomalous magnetic moment, of muon scattering and production cross sections, of μ-mesoatom levels/ CP-even interaction /27/ would not have been detected. As for the CP non-invariant interaction /26/, with the coupling constant $C_{-} \sim 10^{-2}$ it would have given the neutron dipole moment by 2-4 orders of magnitude larger than existing upper limit for it. It is a serious argument against existance of such CP-odd interaction.

Let us notice that in the paper by Gaillard [20] the limit was obtained: $\Gamma(K_S \to 2\mu) \geqslant 1.6\cdot 10^{-7}\,\Gamma_S$. This limit can be reached if the phase of T_S^μ is equal to 45°. Thereby the paper [20] implicitly introduced the anomalous $2\pi - 2\mu$ interaction. Besides, when analysing $K_L \to 2\gamma$ decay it was assumed is paper [20] that in expression $T_L^\gamma = T_2^\gamma + \varepsilon T_1^\gamma$ the modulus of the term εT_1^γ equals to its maximally permissible by experiment value /16/ and its phase coincides with the phase of T_L^γ. This additionally requires anomalously strong $2\pi - 2\gamma$ interaction. Such a combination of four anomalies /in decays $K_L \to 2\gamma$, $K_S \to 2\gamma$ and in interactions $2\pi - 2\mu$, $2\pi - 2\gamma$/ seems to us very unplausible.

Thus if one manages to prove that Γ_L^γ is indeed 5.10^{-4} and $\Gamma_S^\mu/\Gamma_S < 3.10^{-7}$ then the hypothesis [12] on cancellation of the amplitudes T_2^μ and εT_1^μ could be rejected. The dependence of the bounds for Γ_S^μ /22'/ on the value of Γ_L^γ was already discussed above. If $\Gamma_L^\gamma = 4.10^{-4}\Gamma_\omega$ then instead of /22'/ and /24'/ we would have $\Gamma_S^\mu \geqslant 2.9\cdot 10^{-7}\Gamma_S$ and $\Gamma_S^\mu = 1.6\cdot 10^{-7}\Gamma_S$, respectively.

Let us now turn to discussion of such mechanism of cancellation of the two-photon absorptive part of the amplitude T_L^μ for which the term εT_1^μ is inessential.

3.3. Anomalous $3\pi - 2\mu$ Interaction [18,21]

Is it possible to enhance the three-pion imaginary part by 3-4 orders of magnitude as compared with the natural estimates given in the Table I? If one tries the $3\pi - 2\gamma$ anomalous interaction as a mechanism for the enhancement, the answer to this question seems to be negative.

Indeed, such an interaction would lead, generally speaking, to a large cross section of 3π mesons production in colliding lepton beams [21, 22] /Fig. 7/ in contradiction with experiment. This contradiction can be avoided if one assumes [22] that the $3\pi- 2\gamma$ interaction is strong when mass of 3π mesons is less than or equal to the K-meson mass, and at larger mass its amplitude falls shaprly. However, such a possibility seems to be very artificial.

If the three-pion contribution is large due to direct interaction of muons with hadrons, then there appears no direct interaction of muons with hadrons, the there appears no direct disagreement with experiment. In particular, the data on energy level of μ-mesoatoms, on $(g - 2)$ of muon, on elastic muon-proton scattering and inelastic muon-proton scattering with pions production /see Figs. 8,9,10, where the black bulb denotes the discussed anomalous $3\pi- 2\mu$ interaction/ do not rule out such rather strong interaction. The most effective way to search for such anomlaous interaction seems to be experiments on elastic muon-proton scattering at intermediate energies /$\sim$1 GeV/ and large angles /$\sim 180^\circ$/ and experiments on measurement of $(g-2)$. /The graph of Fig. 8

gives the value close to experimental upper bound if the cut off is equal to $\sim$ 1 GeV/.

Nevertheless, the explanation of the result /1/ by anomalously strong $3\pi - 2\mu$ interaction seems to be rather uncredible for the reason pointed out by M.Zh. Shmatikov. Indeed, the imaginary part of the amplitude of transition $K_L \to 3\pi \to 2\mu$ is relatively small because the phase space of three π mesons is proportional to Q^2 and small, $Q \approx 70$ MeV being the energy release in $K \to 3\pi$ decay. In the real part of the amplitude this suppression is, generally speaking, absent. Thus, if imaginary part of the amplitude of Fig. 11 is equal to imaginary part of the amplitude of Fig. 1, then the real part corresponding to Fig. 11 would be $(m_x/Q)^2$ times larger, m_x being some intrinsic hadronic mass, and the contribution from the real part to the probability would be $(m_x/Q)^4$ times larger than the contribution of imaginary part. If for example, $m_x = m_K$, then $(m_x/Q)^4 \approx 2.5 \cdot 10^3$. Therefore, the real part of the amplitude of transition $K \to 3\pi \to 2\mu$ must be with high accuracy canceled out by some mechanism. The necessity for such a cancellation seems to be a very serious argument against introducing $3\pi - 2\mu$ anomalous interaction.

3.4. Anomalous $2\pi\gamma - 2\mu$ Interaction [23]

Anomalous $2\pi\gamma - 2\mu$ interaction with the value of coupling constant which is necessary to cancel the two-photon absorptive part of the amplitude T_L^μ by the contribution from $2\pi\gamma$ state would lead to anomalously large cross section of muonic pair photoproduction /Fig. 13/. Basing on the experimental data the possibility of such interaction may be excluded.

3.4. Anomalous $2\gamma - 2\mu$ Interaction [24]

The two-photon absorptive part might be canceled if $2\gamma - 2\mu$ direct coupling would exist /Fig. 14/. However, in order no to come in conflict with the data on (g-2) and on muon photoproduction one should assume that the anomalous interaction strength is maximal when the mass of two muons is of the order of K-meson mass and sharply decreases with the growth of the mass of two muons and with its decrease /more rigorously, with the decrease of the photon energy/. Such behaviour requires in turn the effective size of muon to be large and qualitatively contradicts the excellent agreement between the data on $(g - 2)$ and the quantum electrodynamics calculations. One can rule out this possibility in a direct way by measuring with a per cent accuracy the cross section of muonic pair production in the processes represented in Figs. 15 and 16.

3.6. Strong Interaction of Muons

Strong interaction of muons, if exists, could of course change the value of the graph of Fig. 1. However it would change to the same extent the value of $(g - 2)$ that is impermissible taking into account the high precision to which muon quantum electrodynamics has been checked. The only possibility which seems to us to be not excluded is the existence of dimuon resonance with the mass $m(2\mu) \simeq m_K$ and so narrow that its contribution to $(g - 2)$ is small enough. Such a resonance would be in fact a new particle and we will discuss it in the next section.

IV. NEW PARTICLES?

What Are They For?

The simplest way in which some yet unknown particles could change the theoretical bound on $K_L \to 2\mu$ decay probability is to give a contribution to imaginary /absorptive/ part of the amplitude $K_2 \to 2\mu$. For this to be true these particles must be coupled both with K_L mesons and muons. Moreover, the less the K_L meson decay width into these particles is, the stronger the latter interact with muons in order to give the required absorptive part of $K_2 \to 2\mu$ amplitude.

4.1. Main Restrictions

What can be said about the properties of these new particles if they exist? First of all they seem to be neutral since otherwise K_L decay into charged particles would have been easily detected.

Furthermore they cannot interact strongly with muons since such an interaction would have lead to incorrect value of $(g - 2)$.

The branching ratio of K_L meson decay into these particles cannot be large. This statement is undoubtedly correct if these particles or their decay products are detectable. If these particle are "invisible" /stable and interact weakly or decay in invisible channels, e.g. into neutrinos/ then the situation is less certain.

If a large leakage into invisible channel would exist then the sum of partial widths of K_L meson would not be equal to the total width. However, as far as we know, there were no experiments on direct comparison of the number of produced and decayed K_L mesons. The following argument shows that the leakage into invisible channel cannot be larger than, say, ten per cent.

For K_s^o meson the absence of invisible channels can be easily checked in experiments which measure not only $K_S^o \to \pi^+\pi^-$ decays but also $K_S^o \to \pi^o\pi^o$ decays. The life time of K_s^o is short and it can be easily verified that the numbers of produced and decayed K_s^o mesons can be determined, for example, by selecting the events with Λ hyperon decay, the latter accompanying K_s^o production, or observing the charge exchange reaction $K^-p \to K^on$.

Furthermore, the phenomenon of so called "vacuum regeneration" in which the interference of $\pi^+\pi^-$ decays of K_L^o and K_S^o mesons is observed allows one to determine the parameter $\eta_{+-} = T^{\pi^+\pi^-}_{K_L} / T^{\pi^+\pi^-}_{K_S}$, and consequently, the absolute and not only relative value of $\Gamma(K_L \to \pi^+\pi^-)$. On the other hand, the ratio $\Gamma(K_L \to \pi^+\pi^-)/ \Gamma(K_L \to$ all observed channels) is known. Thus one can determine the width of all the observed channels and by comparing it with the inverse life time of K_L^o meson as extracted from exponential curve of number of decays as a function of time one can check with reasonable accuracy the coincidence of the total width with the sum of observed partial widths.

Indirectly the leakage of K_L mesons could be detected through apparent violation of $\Delta T = 1/2$ rule for leptonic and nonleptonic decays of neutral and charged kaons. Indeed, these rules are compared with experimental data under implicit assumption that the sum of partial widths is equal to the total width. At present time the $\Delta T = 1/2$ rule is known to be valid with a per cent accuracy and from this one asserts that the leakage of K_L mesons is not large if exists at all.

What experiments could shed light on existence of new neutral light particles? To answer this question let us consider some particular models discussed recently.

4.2. Model of λ Particles [25]

In this model the existence of neutral particles with spin 1/2 and mass less than 0.5 m_K is postulated, the absorptive part of the $K_L \to 2\mu$ decay amplitude receiving contribution from the graph of Fig. 17. If λ particles are stable and do not interact strongly or electromagnetically, they must be penetrating and in this respect be like neutrino. Penetrating the shield they would lead to observable effects in neutrino experiments, for instance to $\mu^+\mu^-$ pair production /see Fig. 18/ in spark chambers.

Estimates show that the number of such pairs produced by λ particles under condition of neutrino experiment performed in CERN would be of the order $0.1 \div 1$ that is by two orders of magnitude larger than the number of muonic pairs produced by neutrino through usual V-A interaction /see Fig. 19/.

It is worth emphasizing that the expected number of muonic pairs produced by λ particles is determined uniquely by the product of $K_L \to \lambda\bar{\lambda}$ decay width and cross section of reaction $\lambda Z \to \lambda\mu^{\pm}\mu^{-}Z$; but just this product is fixed by the value of imaginary part of the graph of Fig. 17. From this point of view the improvement of accuracy of the neutrino experiment by an order of magnitude would be of great interest. Of special interest would be experiments in which as a source of neutrino /and, possibly, of other penetrating particles/ serve not charged but neutral K^0, $\bar{K}^0$ mesons.

4.3. Model of χ^0 Meson [27,28]*

In this model it is postulated [27] the existence of a neutral vector meson with mass of the order 350 MeV contributing to absorptive part of the amplitude $K_L \to 2\mu$ through the graph of Fig. 20. The interaction of χ meson with muon would contribute also to the magnetic moment of muon /see Fig. 21/. From the data on $(g - 2)$ it follows the bound on the constant of this interaction. It leads in turn to the follwing inequality.

$$\Gamma\left(K_L \to \chi^0\gamma\right) / \Gamma_L > 0.6\cdot 10^{-2}\ {}^{**} \qquad /30/$$

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*We are indebted to A.G. Dolgolenko, A.G. Meshkovksy, and V.A. Shebanov for discussion of experimental bounds given in this section. The model of χ^0 was proposed independently by A.N. Moskalev.

**This number differs by a factor of 2 from that given in ref. [27].

However, the model can hardly survive comparison with the present experimental data. Indeed, $\chi^o \to e^+e^-$ decay width is bounded by

$$\Gamma(K_L \to \chi^o\gamma \to e^+e^-\gamma) / \Gamma_L < 2.7\cdot 10^{-5} \quad [28] \qquad /31/$$

The $\chi^o \to \pi^o\gamma$ decay width is bounded

$$\Gamma(K_L \to \chi^o\gamma \to \pi^o\gamma\gamma) / \Gamma_L < 1.5\cdot 10^{-4} \quad [28] \qquad /32/$$

The $\chi^o \to \mu^+\mu^-$ decay with is seemingly bounded by

$$\Gamma(K_L \to \chi^o\gamma \to \mu^+\mu^-\gamma) / \Gamma_L < 4\cdot 10^{-4} \qquad /33/$$

/This latter bound is taken from experimental work [29] where search for $K_L \to \pi^+\pi^-\gamma$ decay was made. This inequality may turn out to be wrong, if in this work the muon registration efficiency was much less than that of π meson/.

If χ^o meson is stable (for $m_\chi < 2m_\mu$) or decays in invisible channels, then for $m_\chi \leqslant 300$ MeV the $\Gamma(K_L \to \chi^o\gamma)$ width seems to be bounded by experiment [30]:

$$\Gamma(K_L \to \chi^o\gamma) / \Gamma_L < 4\cdot 10^{-4} \qquad /34/$$

If $m_{\chi^o} > 2\, m_\mu$, then $\chi^o \to 2\mu$ decay must exist. The fact that this decay was not observed might be due to that it comprises only a small fraction of χ^o decays and the main decay channel of χ^o meson is invisible (e.g. $\chi^o \to \nu\bar{\nu}$). For $m_{\chi^o} \gtrsim 300$ MeV the experiment [30] on the search for decays $K_L \to \gamma$ + neutrals does not exclude such a possibility. However, the $\chi^o \to \nu\bar{\nu}$ decay width may be bounded using the results of neutrino experiment since the events of the type of Fig. 22 were not observed.

4.4. Pseudokaons and Narrow Dimuon

Another class of models intorduces a new particle /pseudokaon/ with a mass very close to m_K so that K_L meson is a coherent superposition of two states K_L and K'_L

$$\hat{K}_L = K_L + \hat{\varepsilon} K'_L \qquad /35/$$

Such an assumption immediately initiates a number of questions: if there also exist K'_S mesons? Do K'_S and K'_L mesons possess strong interactions? What symmetry provides K' and K meson degeneration ? etc. Leaving these questions without answer, let us explain in what way the existence of K' mesons could change theoretical bound for the $K_L \to 2\mu$ decay probability. The idea is the same as of discussed above mechanism of compensation of K_2 and K_1 decays: the $K_L \to 2\mu$ decay amplitude and the $K_L \to 2\mu$ decay amplitude multiplied by $\hat{\varepsilon}$ must destructively interfere.

How can this model be checked?

Experiments with regenrator may be useful for this purpose. Indeed, if strong interactions of K_L and K'_L are different the relative phase of K_L and K'_L would be changed and destructive interference between T^μ_L and $\hat{\varepsilon}T^\mu_L$, weakened by regenerator.

As a result the number of $\hat{K}_L \to 2\mu$ decays after the regenerator would increase. The magnitude of this effect depends on the value of $\hat{\varepsilon}$ and consequently on the $K'_L \to 2\mu$ decay width which we cannot predict.

However, some bounds on the strength of $K'_L\bar{\mu}\gamma_5\mu$ interaction and consequently on the $K'_L \to 2\mu$ decay width can be obtained from existing experimental data. Thus, the width cannot be very samll since otherwise K'_L mesons if they do not possess strong interactions, would penetrate through the shield of neutrino experiment and would give 2μ decays in detectors which were not observed experimentally. The upper bound on the constant of $K_L\ \bar{\mu}\gamma_5\mu$ interaction may be deduced from the measured values of $(g - 2)$ and cross section of the reaction $\mu Z \to \mu\mu\bar{\mu}Z$.

The effect of increasing of $K_L \to 2\mu$ decay width after regenerator discussed above may be large if the $K_L' \to 2\mu$ decay width is close to its lower bound and is of the order of $K_S \to 2\pi$ decay width. The effect is very small if the decay width is close to its upper bound.

An attempt to explain data on both $K_L \to 2\mu$ and $K_L \to 2\pi$ decays in a CP invariant manner by introducing pseudokaons, was made in paper [31].

Concluding the discussion of particles with mass close to the mass of K-meson, let us mention one more possibility. As it was noted by A.L. Lyubimov, there would be no disagreement with $K_L \to 2\mu$ decay if $K_L \to 2\gamma$ decay would have considerably smaller partial width than 5.10^{-4} and in experiments [3] were observed 2γ decays not of K_L meson but of some new particle with a mass, say, by 20 MeV less than the mass of K_L. The decays of this particle into 2μ could not be detected in experiment [2] due to $K_{\mu 3}$ background in this mass region. An argument against this possibility is that the number 5.10^{-4} for the $K_L \to 2\gamma$ branching ratio appears to be the same in a number of experiments in which the energies of K^0 mesons were different.

V. NEW VIOLATIONS OF PRINCIPLES?

If all the other possibilities to solve the $K_L \to 2\mu$ problem are exhausted, then we will be led to revision of general principles which are the basis for theoretical inequality /3/. First of all the validity of the CPT theorem and of the unitarity of S-matrix will be questioned.

5.1. CPT Violation [32,33]

It is well known that there are two sources of imaginary parts in the matrix elements. Firstly it is the presence of real intermediate states giving nonzero absorptive part. Secondly, it is T-violation. So far we considered the first possibility and discussed the contribution to imaginary part of various real intermediate states: 2γ, 3π, $2\pi\gamma$, $\lambda\bar{\lambda}$, $\chi\gamma$ and so on.

Let us now consider the second possibility - the T-invariance. Suppose that the amplitude $K_2\bar{\mu}\gamma_5\mu$ has imaginary part which is due to T-noninvariant interaction with dimensionless constant of the order 10^{-12} that partially cancels out the two-photon imaginary part. Since CP parities of K_2 meson and muonic pair in ${}^1S_o$ state are the same then CP is conserved by such an interaction. Hence T-violation in this case implies CPT violation

$$CP = +1, \quad T = -1, \quad CPT = -1 \qquad /36/$$

If CPT violation satisfies conditions /36/ it is very difficult to detect it in other cases. /It may not be so if this hypothetical interaction violates also some other selection rules, such as $\Delta S < 2$, muonic charge conservation etc./. If however, there exists CPT violating interaction with selection rule

$$CP = -1, \quad T = +1, \quad CPT = -1 \qquad /37/$$

then differences between particles and antiparticles widths and masses would arise. As for the masses, then from the accuracy with which the so called Wu-Yang triangle is checked experimentally it follows that

$$\left(m_{K^o} - m_{\bar{K}^o}\right) / m_K \lesssim 10^{-18} \qquad /38/$$

This rules out the universal CPT violation with selection rules /37/ and dimensionless constant $\sim 10^{-12}$ since P-even interaction with $\Delta S = 0$ would give in this case

$$(m_{K^0} - m_{\bar{K}^0}) / m_K \sim 10^{-2} \qquad /39/$$

As for the differences in life times, they are expected at the level $10^{-4} - 10^{-5}$, that is approximately by one-two orders of magnitude less than existing upper bounds. Since there exists CP noninvariant interaction responsible for $K_L \to 2\pi$ decay

$$CP = -1, \quad T = -1, \quad CPT = +1 \qquad /40/$$

the interaction with selection rules /36/ coupled with the CP-odd interaction /40/ must lead to processes with selection rules /37/. However, dimensionless constant characterising these processes would be considerably less than 10^{-12}. The value of this constant depends upon model of CP violation. Thus, if superweak interaction with the coupling constant of the order 10^{-16} is responsible for CP violation, the dimensionless constant characterising the amplitudes with selection rules /37/ would be of the order 10^{-28}. It will be of the order 10^{-21} if CP is violated in milliweak interaction with the constant of the order 10^{-9}. In any case experiments with K^0-mesons seem most promising for discovery of CPT violation, if it exists, because of the well known enhancement mechanism caused by smallness of $K_L - K_S$ mass difference. As for the possible effects of CPT violation they need further theoretical investigation.

If we return to consideration of the amplitudes of the first order in CPT and T-odd interaction with selection rules /36/ then it is worth to emphasize that such interaction leads to additional phase factors in the amplitudes which are the same for particles and antiparticles and are not connected with the real intermediate states contrary to the usual case.

Such interaction does not lead to difference between the decay amplitudes of a particle and its antiparticle. Therefore the search for this interaction is very difficult. In the case of $K_L \to 2\mu$ decay to prove that just CPT violation is responsible for observed anomaly, it is necessary to rule out expreimentally all other possible explanations. If CPT noninvariant interaction contains pseudoscalar currents then it might manifest itself in $K_L \to 2e$ decay at the accuracy level already achieved. But if this interaction contains only axial currents then the sensitivity of experiment on search for $K_L \to 2e$ decay is to be rised by 4÷5 orders of magnitude to detect this interaction.

We have considered CPT violation within the framework of S-matrix formalism in a purely phenomenological way. In local quantum field theory CPT violation is possible only if some other fundamental principle such as causality, Lorentz invariance, positivity of energy is violated. The question whether it is possible to reconcile these principles with CPT violation in framework of the S-matrix formalism remains now open.

5.2. S-Matrix Unitarity

As is known the unitarity condition

$$S_{ik}\, S^{+}_{km} = \delta_{im} \qquad /41/$$

is a compact mathematical expression of two physical principles: probability conservation and superposition principles. Probability conservation implies that the sum of probabilities of all transitions from given initial state into all possible final ones is equal to unity. The probability conservation is described by the diagonal terms of eq. /41/

$$\sum_{k} S_{ik}\, S^{*}_{ik} = 1 \qquad /42/$$

Nondiagonal terms of eq. /41/ have the form

$$\sum_k S_{ik} S^{*}_{mk} = 0 \qquad (i \neq m) \qquad /43/$$

and express the superposition principle*.

It should be noted that a direct application of the unitarity condition to K-meson decay /derivation of eq. /11/ from eq. /43// needs some specification. The matter is that S-matrix connects stable states while K-meson is unstable. Thus relation /11/ is strictly speaking not an exact but an approximate one. However, the corrections are negligibly small as they are for sure less than Γ_{K_L}/m_K [34].

At the beginning of this review we discussed how does the unitarity of S-matrix together with CPT invariance lead to theoretical limit for $K_L \to 2\mu$ decay probability. In the preceding section we saw how this limit can be changed if one abandones CPT invariance. Unfortunately we cannot discuss the violation of the unitarity of S-matrix even on such phenomenological level as CPT violation since in this case we have to modify the very formalism of S-matrix.

The necessity to check experimentally the unitarity of S-matrix and the superposition principle in connection with the discovery of CP violation was emphasized in papers [35, 36, 37].

The best way to check the unitarity conditions seems to be a comparison with experimental data of the well known Bell-Steinberger relation [38,39] which is an explicit form of relation /43/ in case of $K_{L,S}$ mesons

$$\frac{2i(m_L - m_S) + \Gamma_L + \Gamma_S}{2\Gamma_S} \langle K_S | K_L \rangle = \sum_j \eta_j B_j \qquad /44/$$

*We are grateful to I.Yu. Kobsarev for drawing our attention to that relation /11/ considered as a relation between experimental values may be violated in theories where the superposition principle is not fulfilled.

where $m_{L,S}$ are $K_{L,S}$ meson masses, $B_j = \Gamma(K_S \to j)/\Gamma_S$, $\eta_j = T^j_{K_L} / T^j_{K_S}$.

Here $j = \pi^+\pi^-$, $\pi^o\pi^o$, $\pi e\nu$, $\pi\mu\nu$, $\pi^+\pi^-\pi^o$, $3\pi^o$, $2\pi\gamma$ etc. As it is known, all the quantities entering relation /44/ may be measured experimentally. Let us notice in this connection that it would be desirable to measure $\Gamma(K_S \to 3\pi^o)$ or to get a reasonable bound on its value.

Of interest is also the check of unitarity condition in other decays: in neutron β decay, in $K_{\mu 3}$ decay, in nonleptonic decays of hyperons etc. At present the accuracy of measurement of phases of the corresponding amplitudes is 1^o - 10^o. The absence of any contradiction with theory can be considered as an experimental proof of the validity of the unitarity condition in these processes with such accuracy.

5.3. Other Principles

Today we do not see serious reasons to call in question in connection with $K_L \to 2\mu$ decay other principles of modern physics such as Lorenz invariance, angular momentum conservation, energy and momentum conservation though the accuracy of checking these principles in high energy physics is not sufficient. In our opinion these principles deserves better experimental foundation without any reference to the problem of $K_L \to 2\mu$ decay.

However, this point of view is not shared by all the physists. Thus, according to the hypothesis by B.A. Arbusov, the momentum nonconservation is a feature inherent just to $K_L \to 2\mu$ decay, and muons from $K_L \to 2\mu$ decay were not detected because they had momenta values different from those predicted by energy-momentum conservation.

CONCLUSION

In the last twenty years K-mesons undoubtedly brought into our knowledge of microworld more than any other elementary particle. The discovery of K-mesons played an important role in introducing the quantum number of strangeness. Investigation of K-meson decays led to discovery of C,P, and CP violation.

All these years muons remained one of the most intriguing problems of elementary particle physics. It can be even said that they became even more and more mysterious. We have learned a lot about them during these years but the question why the muon is heavier than the electron has still no answer.

If experiments [2,3] are correct, they may be a kaonic key to the secret of the muon. Very likely in this case that the seventies will bring us the understanding of muon nature.

The brief review given above shows that the present experimental data are insufficient to eliminate a number of more or less plausible anomalies in the domain of physics which is usually believed well studied.

Le us list here the experiments which were discussed in connection with the $K_L \to 2\mu$ problem.

1. Measurement of $K_L \to 2\gamma$ decay probability in several independent experiments and continuing search for $K_L \to 2\mu$ decay.

2. Search for $K_S \to 2\mu$ decay modes up to 10^{-7} and measurement of $K_S \to 2\gamma$ decay probability.

3. Search for anomalous interactions of muons /muon-nucleon elastic and inelastic scattering, muonic pair creation in hadron and photon collisions, $(g - 2)$, μ - mesoatoms/.

4. Search for new light particles in K_L meson decays /study of photonic decays and the balance of produced and decayed K_L mesons, search for anomlaous events in the neutrino experiments, search for $K_L \to 2\mu$ decay under various conditions, particularly, after regenerator/.

5. Check of CPT invariance, in particular, comparison of life times and partial widths of particles and antiparticles with accuracy 10^{-5}.

6. Thorough quantitative check of the Bell-Steinberger relation, measurement of the so called T-odd correlations in $K_{\mu 3}$ decay, neutron β decay and in hyperon decays.

7. Check of the basic conservation principles /energy, momentum, angular momentum conservation, Lorenz invariance/.

Above we made an attempt to analyse how serious is the problem and to confront with experimental data the proposed solutions of this problem. We tried as far as it was possible not to trust subjective estimates and emotions but to consider bare facts. In conclusion we should like to permit ourselves to express our opinion that we came to while looking through various possibilities.

If one tries to make a prognosis then one should proceed not only from that how incomplete and even poor are the experimental data but also from that how attractive are the hypotheses proposed. Almost all the attempts to explain the $K_L \to 2\mu$ problem discussed in this review even if they do not contradict the available experimental data seem to us now quite artificial, unplausible and unattractive. In our opinion, if $K_S \to 2\gamma$ is not detected with relative probability larger than 10^{-7}, the experimental data either on $K_L \to 2\mu$ decay or on $K_L \to 2\gamma$ decay /or on both decays/ will change and come into agreement with the theory.

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Table I

Estimate of relative contributions of various intermediate state into absorptive part of the $K_L \to 2\mu$ decay amplitude

State n	2γ	$\pi^+\pi^-\gamma$	3π	2π	$\pi\mu\nu$
$\mathrm{Im}T(K_L \to n \to 2\mu)$	1	$\leqslant 10^{-2}$ x/	$\sim 10^{-5}$ xx/	$\sim 10^{-3}$ xxx/	$\sim 10^{-7}$ xxxx/

x/ The result is based on the calculation of the graph of Fig. 2 in paper [8] . The result 5.10^{-2} obtained in paper [7] in our opinion is too large.

xx/ The part $3\pi - 2\gamma$ in the amplitude $(3\pi - 2\gamma - 2\mu)$ is calculated in framework of the current algebra [9,10]. If one uses for the amplitude $3\pi - 2\mu$ the simple dimensional estimate $(3\pi | 2\mu) = \alpha\, m_K^{-2}\; \bar{\mu}\gamma_5\mu\; \varphi_\pi^3$ the result is approximately by an order of magnitude larger than that given in the Table.

xxx/ The result is based on dimensional estimate for the amplitude of $2\pi - 2\mu$ transition: $(2\pi | 2\mu) = \alpha^2 m_K^{-1}\, \bar{\mu}\mu$. The perturbation theory calculations are given in paper [6].

xxxx/ The result is based on the rough estimate of the graph of Fig. 6.

APPENDIX I

Two-photon absorptive part of $K_2 \to 2\mu$ decay amplitude

The two-photon contribution into Abs T_2^μ is determined by the graph of Fig. 1 and may be written as

$$\text{Abs } T_2^\mu = \frac{1}{2}\int d\tau \sum_{e_1, e_2} T_2^\gamma \, \left(T_\mu^\gamma\right)^* \qquad /1.1/$$

where

$$d\tau = \frac{1}{2!} \frac{d^3k_1}{(2\pi)^3 \, 2\omega_1} \frac{d^3k_2}{(2\pi)^3 \, 2\omega_2} \cdot (2\pi)^4 \delta^4\left(P-k_1-k_2\right)$$

is the phase space of two γ quanta; the sum is taken over the polarization states (e_i), P is the K-meson momentum.

The amplitude of $(2\gamma - 2\mu)$ transition is described by the electrodynamics and is equal to

$$T_\mu^\gamma = e^2 \bar{\mu}_1 \left\{ \hat{e}_1 \frac{\hat{p}_1 - \hat{k}_1 + m_\mu}{2p_1k_1} \hat{e}_2 + \hat{e}_2 \frac{\hat{p}_1 - \hat{k}_2 + m_\mu}{2p_1k_2} \hat{e}_1 \right\} \mu_2 \qquad /1.2/$$

where p_1 and p_2 are momenta of μ^- and μ^+, respectively.

The amplitudes T_2^μ and T_2^γ are theoretically unknown and we parametrize them in the most general way without assuming CP conservation.

Because of the angular momentum conservation the pair $(\mu^+\mu^-)$ produced in K_2^o-meson decay may be either in ^1S_o or in ^3P_o states. Since CP parity of the system fermion-antifermion is equal to $(-1)^{S+1}$ where S is the total spin of the pair, then for K_2 decays into ^1S_o state CP is conserved $\left(CP\,(\bar{\mu}\mu)_{^1S_o} = -1\right)$ and for K_2 decays into ^3P_o

state CP is violated $\left(CP(\bar{\mu}\mu)_{^3P_0} = +1\right)$. So the amplitude T_2^{μ} may be represented as a sum of two terms:

$$T_2^{\mu} = \varphi_{K_2}\bar{\mu}_1 \left[\tilde{T}_2^{\mu -}\gamma_5 + \frac{i}{v}\tilde{T}_2^{\mu +}\right]\mu_2 \qquad /1.3/$$

where $v = 0.9$ is the velocity of muon in the rest frame of K-meson and indices "$\pm$" indicate CP parity of the pair $(\mu^+\mu^-)$. The factor i is chosen so that in the absence of the real intermediate states /i.e., if $K_2 \to 2\gamma$ decay would have been forbidden/ the quantities $\tilde{T}_2^{\mu\pm}$ were real. Indeed, in absence of real intermediate states the decay amplitude may be considered as the matrix element of some effective Lagrangian that must be hermitian. The statement that $\tilde{T}_2^{\mu\pm}$ is real follows in this case from the fact that the quantities φ_{K_2}, $\bar{\mu}\gamma_5\mu$, $i\bar{\mu}\mu$ are antihermitian.

Analogously, the amplitude $K_2 \to 2\gamma$ can be written as

$$T_2^{\gamma} = i\varphi_{K_2}\left[\tilde{T}_2^{\gamma -}\left((k_1k_2)(e_1e_2) - (k_1e_2)(k_2e_1)\right) + \right.$$

$$\left. + \tilde{T}_2^{\gamma +}\varepsilon_{\alpha\beta\gamma\delta}e_{1\alpha}\, e_{2\beta}k_{1\gamma}k_{2\delta}\right] \qquad /1.4/$$

The first term in expression /1.4/ corresponds to CP even transition and the second one to CP odd transition.

The amplitudes are related to the decay width in the following way

$$\Gamma(K_2 \to 2\mu) = \frac{m_k V}{8\pi}\left[|\tilde{T}_2^{\mu -}|^2 + |\tilde{T}_2^{\mu +}|^2\right] \qquad /1.5/$$

$$\Gamma(K_2 \to 2\gamma) = \frac{m_k}{64\pi}\left[|\tilde{T}_2^{\gamma -}|^2 + |\tilde{T}_2^{\gamma +}|^2\right] \qquad /1.6/$$

Substituting eqs. /1.2, 1.3, 1.4/ into eq. /1.1/ we get the relation which determines imaginary parts of form factors $\tilde{T}_2^{\mu\pm}$:

$$\bar{\mu}_1\left[\mathrm{Im}\ \tilde{T}_2^{\mu -}\ \gamma_5 + \frac{i}{v}\ \mathrm{Im}\ \tilde{T}_2^{\mu +}\right]\mu_2 =$$

$$= ie^2 \int d\tau \left[\tilde{T}_2^{\gamma -}\left(g_{\alpha\beta}\ k_1 k_2 - k_{1\beta}\ k_{2\alpha}\right) + \right. \qquad /1.7/$$

$$\left. + \tilde{T}_2^{\gamma +}\ \varepsilon_{\alpha\beta\gamma\delta}\ k_{1\gamma}\ k_{2\delta}\right]\ \bar{\mu}_1\ \frac{\hat{p}_1 - \hat{k}_1 + m_\mu}{2p_1 k_1}\ \gamma_\beta \mu_2$$

Note that both terms in expression /1.2/ give equal contributions to $\mathrm{Im}\ \tilde{T}_2^{\mu\pm}$

If we use now the relations

$$\bar{\mu}_1\ \gamma_\alpha\left(p_1 - k_1 + m_\mu\right)\gamma_\beta \mu_2 = \bar{\mu}_1\left(2p_{1\alpha} - \gamma_\alpha \hat{k}_1\right)\gamma_\beta \mu_2 \qquad /1.8/$$

$$\gamma_\alpha \hat{k}_1\ \gamma_\beta\ \varepsilon_{\alpha\beta\gamma\delta}\ k_{1\gamma}\ k_{2\delta} = 2i\left(k_1 k_2\right)\gamma_5 \hat{k}_1 \qquad /1.9/$$

and take into account that the term $\sim p_{1\alpha}$ in eq. /1.8/ vanishes when integrated over $d\tau$, then eq. /1.7/ reduces to the following two equations

$$\mathrm{Im}\ T_2^{\mu -}\ \bar{\mu}_1 \gamma_5 \mu_2 = \frac{e^2}{2}\ m_k^2\ \tilde{T}_2^{\gamma -} \int \frac{d\tau}{p_1 k_1}\ \bar{\mu}_1 \gamma_5 \hat{k}_1 \mu_2 \qquad /1.10/$$

$$\mathrm{Im}\ \tilde{T}_2^{\mu +}\ \bar{\mu}_1 \mu_2 = \frac{ve^2}{2}\ \tilde{T}_2^{\gamma +} \int \frac{d\tau}{p_1 k_1}\ \bar{\mu}_1\left(m_\mu m_k^2 + 2p_1 k_1 \hat{k}_1\right)\mu_2 \qquad /1.11/$$

Integration in eqs. /1.10/, /1.11/ is trivial and finally we get

$$\frac{\mathrm{Im}\tilde{T}_2^{\mu -}}{\tilde{T}_2^{\gamma -}} = \frac{1}{v}\ \frac{\mathrm{Im}\tilde{T}_2^{\mu +}}{\tilde{T}_2^{\gamma +}} = \frac{\alpha}{4}\ \frac{m_\mu}{m_k}\ \frac{1}{v}\ \ln\ \frac{1+v}{1-v} \qquad /1.12/$$

In deriving this equation we implicitly assumed that $\mathrm{Im}\tilde{T}_2^{\gamma\pm} = 0$. The consideration of the unitarity condition for $K_2 \to 2\gamma$ decay shows that this is indeed so if there are no anomalously strong interactions $3\pi - 2\gamma$ and $2\pi - 2\gamma$.

Substituting eqs. /1.12/ into eq. /1.5/ and making use of relation /1.6/ we get the lower bound for the $K_2 \to \mu^+\mu^-$ decay width /formulae /12/ and /13//.

APPENDIX 2

Unitarity condition for K_L-meson decays

According to arguments given in sec. 2.1, the unitarity condition directly allows one to calculate imaginary part only for transitions between states with definite CP /or C/ parity. Thus we first consider amplitudes of transitions

$$T(K_1 \to 2\mu) = T_1^{\mu +} + iT_1^{\mu -} \qquad /2.1/$$

$$T(K_2 \to 2\mu) = T_2^{\mu -} + iT_2^{\mu +} \qquad /2.2/$$

Here indices "±" indicate CP parity of the $\bar{\mu}\mu$ pair. The unitarity condition for these amplitudes takes the form

$$\mathrm{Im}\, T_1^{\mu_\pm} = \frac{1}{2}\sum_n \int d\tau_n \left[T_1^{n_\pm}\left(T_{\mu_\pm}^{n_\pm}\right)^* \mp T_1^{n_\mp}\left(T_{\mu_\pm}^{n_\mp}\right)^* \right] \qquad /2.3/$$

$$\mathrm{Im}\, T_2^{\mu_\pm} = \frac{1}{2}\sum_n \int d\tau_n \left[T_2^{n_\pm}\left(T_{\mu_\pm}^{n_\pm}\right)^* \pm T_2^{n_\mp}\left(T_{\mu_\pm}^{n_\mp}\right)^* \right] \qquad /2.4/$$

Here $T_{\mu\pm}^{n\pm}$ is CP-even amplitude of $(\bar{\mu}\mu)_\pm \to n_\pm$ transition and $T_{\mu\pm}^{n_\mp}$ is CP-odd amplitude of $(\bar{\mu}\mu)_\pm \to n_\mp$ transition.

Now if ine takes into account that $K_2 = K_L - \varepsilon K_1$, then

$$T_2^{n-} = T_L^{\ -} - i\varepsilon T_1^{n-} \qquad /2.5/$$

$$T_2^{n+} = T_L^{n+} + i\varepsilon T_1^{n+} \qquad /2.6/$$

Substituting eqs. /2.5/, /2.6/ into eq. /2.4/ we get

$$\mathrm{Im}\, T_L^{\mu\pm} \pm \mathrm{Re}\left(\varepsilon T_1^{\mu\pm}\right) = \sum_n \int d\tau_n \left\{ T_L^{n\pm}\left(T_{\mu_\pm}^{n\pm}\right)^* \pm \right.$$

$$\left. \pm T_L^{n\mp}\left(T_{\mu_\pm}^{n\mp}\right)^* \pm i\varepsilon\left[T_1^{n\pm}\left(T_{\mu_\pm}^{n\pm}\right)^* \mp T_1^{n\mp}\left(T_{\mu_\pm}^{n\mp}\right)^*\right]\right. \qquad /2.7/$$

It follows from eq. /2.3/ that the last two terms in eq. /2.7/ are equal to $\pm\, i\varepsilon \mathrm{Im} T_1^{\mu\pm}$. Making use of this we get finally

$$I_m T_L^{\mu\pm} = \frac{1}{2} \sum_n \int d\tau_n \left[T_L^{n\pm}\left(T_{\mu_\pm}^{n\pm}\right)^* \pm \right.$$

$$\left. \pm T_L^{n\mp}\left(T_{\mu_\pm}^{n\mp}\right)^* \right] \mp \mathrm{Re}\varepsilon\left(T_S^{\mu\pm}\right)^* \qquad /2.8/$$

/Here we replaced εT_1 by εT_S/. This relation expresses imaginary part of the amplitude of $K_L \to 2\mu$ decay through the amplitudes of physical processes and is the basis for further theoretical analysis.

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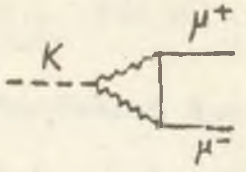

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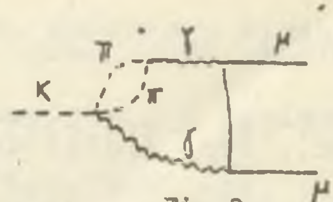

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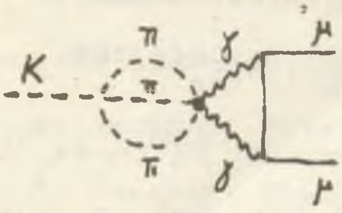

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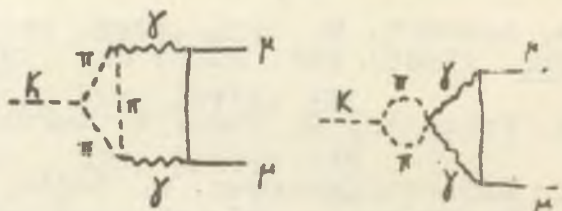

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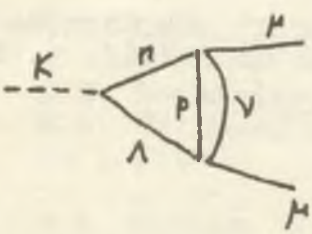

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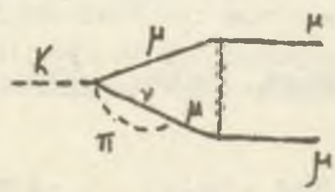

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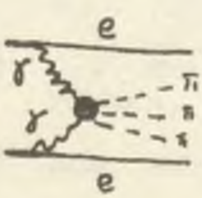

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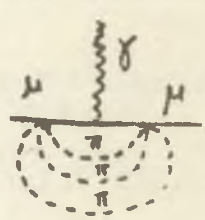

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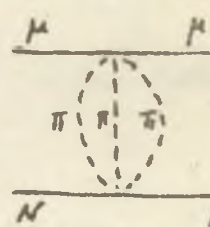

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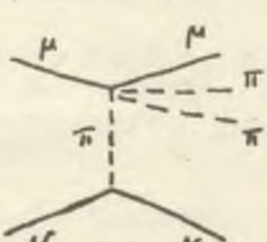

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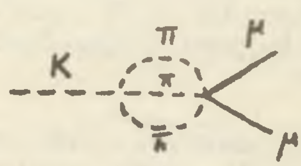

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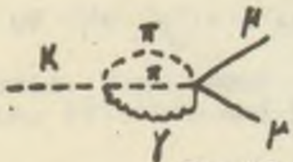

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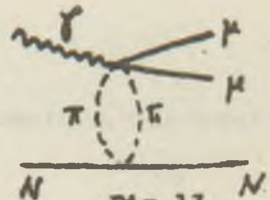
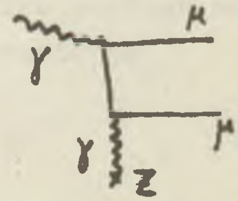


Fig.13

Fig.14



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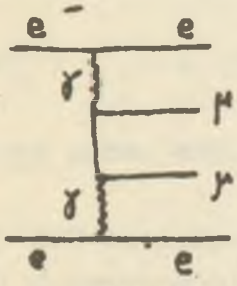

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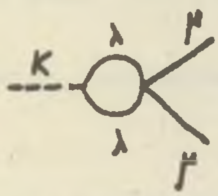
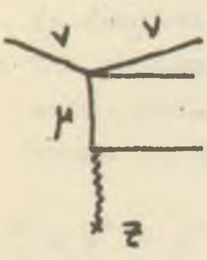


Fig.17

Fig.18



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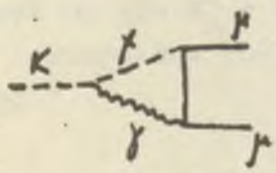

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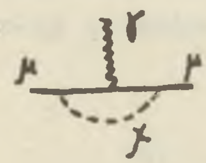

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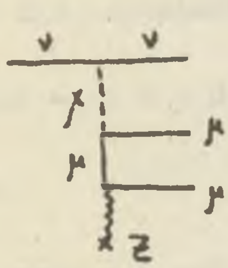

Fig.22

REVIEW OF THE $K_L^o \to \mu^+ + \mu^-$ PUZZLE*

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Among the most recent and interesting developments in weak interactions is the $K_L^0 \to \mu^+ + \mu^-$ puzzle. The present experimental upper limit on the branching ratio is [1]

$$\Gamma(K_L^0 \to \mu^+ + \mu^-)/\Gamma(K_L^0 \to \text{all}) < 1.8 \times 10^{-9} \qquad (1)$$

This experimental <u>upper</u> limit, which is the 90 % confidence level, is well below the theoretical <u>lower</u> bound of [2]

$$\Gamma(K_L^0 \to \mu^+ + \mu^-)/\Gamma(K_L^0 \to \text{all}) > 6 \times 10^{-9} \qquad (2)$$

obtained under the following assumptions:

(A) CPT invariance
(B) Conventional quantum electrodynamics
(C) CP conservation
(D) The absorptive part of the $K_L^0 \to \mu^+ + \mu^-$ amplitude is dominated by the on-shell 2γ intermediate state.
(E) The absorptive part of the $K_L^0 \to \gamma + \gamma$ amplitude is zero. With these assumptions the unitarity relation leads to the bound [2]

$$\Gamma(K_L^0 \to \mu^+ + \mu^-)/\Gamma(K_L^0 \to \gamma + \gamma) > 1.2 \times 10^{-5} \qquad (3)$$

which, when combined with the experimental branching ratio[3],

$$\Gamma(K_L^0 \to \gamma + \gamma)/\Gamma(K_L^0 \to \text{all}) = (4.9 \pm 0.4) \times 10^{-4} \qquad (4)$$

yields the lower bound (2).

Taking the discrepancy between (1) and (2) seriously, the validity of assumptions (A) - (E) has been examined in a number of recent analyses [4-10] of the $K_L^0 \to \mu^+ + \mu^-$ puzzle. Several interesting results have emerged which we shall discuss.

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Assuming CPT invariance, but not CP invariance, the decaying kaon states are

$$K_S^o = K_1^o + \varepsilon K_2^o \tag{5}$$

and

$$K_L^o = K_2^o + \varepsilon K_1^o \tag{6}$$

where K_1^o and K_2^o are the CP even and odd combinations of K^o and $\bar{K}^o$, and $\varepsilon = (2 \times 10^{-3}) \exp (i\pi/4)$. Similarly, the final $\mu^+\mu^-$ state can have a CP even $^3P_0$ component and a CP odd $^1S_0$ component. Thus, separating the CPT invariant transition interaction into a CP even piece H_+ and a CP odd piece H_-, we can define four complex amplitudes:

$$A_+ = \langle \mu^+\mu^-;\ ^3P_0 | H_+ | K_1^o \rangle \tag{7}$$

$$A_- = \langle \mu^+\mu^-;\ ^1S_0 | H_+ | K_2^o \rangle \tag{8}$$

$$B_+ = \langle \mu^+\mu^-;\ ^3P_0 | H_- | K_2^o \rangle \tag{9}$$

$$B_- = \langle \mu^+\mu^-;\ ^1S_0 | H_- | K_1^o \rangle \tag{10}$$

The amplitudes for K_L^o decay into the $^3P_0$ and $^1S_0$ $\mu^+\mu^-$ states are then, respectively

$$L_+ = B_+ + \varepsilon A_+ \tag{11}$$

and

$$L_- = A_- + \epsilon B_- \tag{12}$$

The corresponding K_S^o decay amplitudes are

$$S_+ = A_+ + \epsilon B_+ \tag{13}$$

and

$$S_- = B_- + \epsilon A_- \tag{14}$$

These amplitudes are defined so that CPT invariance implies their absorptive parts, which are what enter in the unitarity relation, are the following:

$$\text{Abs } A_\pm = \text{Im } A_\pm \tag{15}$$

$$\text{Abs } B_\pm = -i \text{ Re } B_\pm \tag{16}$$

$$\text{Abs } L_+ = -i \text{ Re } B_+ + \epsilon \text{ Im } A_+ \tag{17}$$

$$\text{Abs } L_- = \text{Im } A_- - i\, \epsilon \text{ Re } B_- \tag{18}$$

$$\text{Abs } S_+ = \text{Im } A_+ - i\, \epsilon \text{ Re } B_+ \tag{19}$$

$$\text{Abs } S_- = -i \text{ Re } B_- + \epsilon \text{ Im } A_- \tag{20}$$

In the rates the amplitudes to the $^3P_0$ and $^1S_0$ $\mu^+\mu^-$ states add incoherently, so

$$\Gamma(K_L^0 \to \mu^+ + \mu^-) = |L_-|^2 + |L_+|^2 =$$

$$= |A_- + \varepsilon B_-|^2 + |B_+ + \varepsilon A_+|^2 \tag{21}$$

and

$$\Gamma(K_S^0 \to \mu^+ + \mu^-) = |S_+|^2 + |S_-|^2 =$$
$$= |A_+ + \varepsilon B_+|^2 + |B_- + \varepsilon A_-|^2 \tag{22}$$

If CP invariance were assumed all the above equations would simplify considerably since CP invariance requires both $B_\pm = 0$ and $\varepsilon = 0$.

Next let us consider the implications and validity of assumptions (A) - (E). Since CPT [assumption (A)] and quantum electrodynamics [assumption (B)] are on relatively firm footing, we will take their validity for granted throughout. Retaining, in addition, CP invariance [assumption (C)], the validity of assumptions (D) and (E) has been studied. The effect of including the additional allowed intermediate states, $2\pi\gamma$ and 3π, in the unitarity relation has been shown not to change the bound, (2), significantly. Even allowing for CP nonconservation, estimates indicate that these contributions do not amount to more than about 20 % of the 2γ intermediate state contribution.[4-5] We shall therefore retain assumptions (A), (B), (D) and (E) and examine the consequences of relaxing assumption (C), i.e. CP invariance.[6-7]

The first important point is that CP nonconserving theories in which the CP noninvariance enters only through the mass matrix, e.g. the superweak theory,[11] can be ruled out. In the superweak theory $\varepsilon \neq 0$ but $B_\pm = 0$. Therefore, from Eqs. (18) and (21)

$$\Gamma(K_L^0 \to \mu^+ + \mu^-) > |\mathrm{Im}\, A_-|^2 = |\mathrm{Abs}\, L_-|^2 \tag{23}$$

which yields the original bound (2) in disagreement with the experimental limit, (1).

More generally Eqs. (18) and (21) lead to

$$\begin{aligned}\Gamma(K_L^0 \to \mu^+ + \mu^-) &> |\mathrm{Im}\, A_- + \mathrm{Im}(\varepsilon\, B_-)|^2 \\ &= |\mathrm{Abs}\, L_- + i(\mathrm{Re}\,\varepsilon)\, B_-^*|^2 \\ &> \left(|\mathrm{Abs}\, L_-| - \mathrm{Re}\,\varepsilon |B_-|\right)^2\end{aligned} \tag{24}$$

Taking the experimental result, (1), then implies that B_- is not at all small but is limited to the range

$$|\mathrm{Abs}\ L_-| - \Gamma^{1/2}(K_L^o \to \mu^+ + \mu^-)$$
$$< \mathrm{Re}\,\varepsilon |B_-| < \mathrm{Abs}\ L_-| + \Gamma^{1/2}(K_L^o \to \mu^+ + \mu^-) \qquad (25)$$

or

$$0.34 \times 10^3 < |B_-| / |\mathrm{Abs}\ L_-| < 1.1 \times 10^3 \qquad (26)$$

if Abs L_+ can be neglected; i.e. Abs L_- has been taken from Eq. (33) as discussed below.

This result not only rules out the possibility of explaining the $K_L^o \to \mu^+ + \mu^-$ puzzle as arising from CP nonconservation due to a superweak interaction but, in addition, there are important consequences for the $K_S^o \to \mu^+ + \mu^-$ rate which we discuss next, following largely the work of Christ and Lee [6].

More general inequalities can be obtained by retaining assumptions (A), (B), (D) and (E) but not assuming CP invariance [assumption (C)] as follows: From Eqs. (11) - (20) one finds

$$\mathrm{Abs}\ L_- - \mathrm{Im}\ L_- = -i\ B_-^* \ \mathrm{Re}\,\varepsilon \qquad (27)$$

and

$$\text{Abs } L_+ + i \text{ Re } L_+ = i A_+^* \text{ Re}\varepsilon. \qquad (28)$$

From the resulting triangle inequalities

$$|\text{Re } L_+| > \Big| \text{Re}\varepsilon |A_+| - |\text{Abs } L_+| \Big| \qquad (29)$$

and

$$|\text{Im } L_-| > \Big| \text{Re}\varepsilon |B_-| - |\text{Abs } L_-| \Big| \qquad (30)$$

it follows that

$$\Gamma(K_L^o \to \mu^+ + \mu^-) > |\text{Re } L_+|^2 + |\text{Im } L_-|^2$$

$$> \left(|\text{Abs } L| - \text{Re}\varepsilon \Gamma^{1/2}(K_S^o \to \mu^+ + \mu^-)\right)^2 \qquad (31)$$

where $|\text{Abs } L|$ is defined by

$$|\text{Abs } L|^2 = |\text{Abs } L_+|^2 + |\text{Abs } L_-|^2 \qquad (32)$$

and terms of order ε^2 have been neglected. It is clear from (31) that $\Gamma(K_S^o \to \mu^+ + \mu^-)$ is bounded both, above

and below as a function of $\Gamma(K_L^0 \to \mu^+ + \mu^-)$. However, further specific assumptions about Abs L are needed to determine the magnitude of the bounds.
Abs L can be evaluated in terms of the amplitudes for $K_L^0 \to \gamma + \gamma$ in the 0^+ CP-even and 0^- CP-odd final states. One finds,

$$\text{Abs } L_- = \lambda \langle \gamma\gamma;\ 0^- | H | K_L^0 \rangle \tag{33}$$

$$\text{Abs } L_+ = \lambda v \langle \gamma\gamma;\ 0^+ | H | K_L^0 \rangle \tag{34}$$

where

$$\lambda^2 = \frac{1}{2}\ \frac{\alpha^2}{v} \left(\frac{m_\mu}{m_K}\right)^2 \left[\log \frac{1+v}{1-v}\right]^2 \approx 1.2 \times 10^{-5} \tag{35}$$

and

$$v = \left[1 - 4\,\frac{m_\mu^2}{m_K^2}\right]^{1/2} \approx 0.9 \tag{36}$$

is the muon velocity. The $K_L^0 \to \gamma + \gamma$ total rate is then

$$\Gamma(K_L^0 \to \gamma + \gamma) = \frac{|\text{Abs } L_-|^2}{\lambda^2} + \frac{|\text{Abs } L_+|^2}{(\lambda v)^2} \tag{37}$$

and therefore (38)

$$\lambda^2 v^2 \Gamma(K_L^0 \to \gamma + \gamma) < |\text{Abs } L_-|^2 + |\text{Abs } L_+|^2 < \lambda^2 \Gamma(K_L^0 \to \gamma + \gamma)$$

Combining the inequality (31) with these bounds (38) on $|\mathrm{Abs}\, L_6|$ one readily finds the results of Christ and Lee

$$\mathrm{Re}\,\varepsilon\Gamma^{1/2}(K_S^0 \to \mu^+ + \mu^-) > \lambda v\Gamma^{1/2}(K_L^0 \to \gamma + \gamma) - \Gamma^{1/2}(K_L^0 \to \mu^+ + \mu^-) \tag{39}$$

and

$$\mathrm{Re}\,\varepsilon\Gamma^{1/2}(K_S^0 \to \mu^+ + \mu^-) < \lambda\Gamma^{1/2}(K_L^0 \to \gamma + \gamma) + \Gamma^{1/2}(K_L^0 \to \mu^+ + \mu^-) \tag{40}$$

In the CP invariant limit, $\varepsilon = 0$ and $B_\pm = 0$,(39) reduces to the original bound (2) (in disagreement with experiment) and (40) becomes trivial.

Using the experimental limit (1) on $\Gamma(K_L^0 \to \mu^+ + \mu^-)$ and the observed $K_L^0 \to \gamma + \gamma$ branching ratio (4) inequalities (39) and (40) imply

$$1 \times 10^{-5} > \frac{\Gamma(K_S^0 \to \mu^+ + \mu^-)}{\Gamma(K_S^0 \to \mathrm{all})} > 5 \times 10^{-7} \tag{41}$$

The decay $K_S^0 \to \mu^+ + \mu^-$ has not been seen and the present upper limit is [12]

$$\Gamma(K_S^0 \to \mu^+ + \mu^-)/\Gamma(K_S^0 \to \mathrm{all}) < 1.6 \times 10^{-6} \tag{42}$$

Figure 1 shows the theoretical bounds and the experimental limits. Under more restrictive assumptions the Christ - Lee bounds can be altered somewhat [13]. For example, an interesting possibility is that $K_L^o \to \mu^+ + \mu^-$ decay proceeds predominately to the CP odd ^1S_O final state. The cancellation then occurs between the amplitudes A_- and εB_- and the lower Christ-Lee bound is increased, becoming the dashed line in the figure. The lower bound (41) then becomes

$$\Gamma(K_S^o \to \mu^+ + \mu^-)/\Gamma(K_S^o \to \text{all}) > 10 \times 10^{-7} \tag{43}$$

which is not too far from the most recent experimental limit (42), offering the possibility of an early confrontation.

It is not hard to invent specific CP nonconserving interactions that are able to provide an amplitude B_- of sufficient size to cancel the absorptive amplitude coming from the two photon intermediate state. A completely ad hoc example is the interaction [13]

$$g\, G/\sqrt{2}\; A_\lambda^{(7)}\; \bar{\mu}\gamma^\lambda\gamma^5\mu \tag{44}$$

where $A_\lambda^{(7)}$ is the seventh member of the axial vector octet of currents. With the constant $g \sim 10^{-2}$, which puts this interaction in the milliweak class, the resulting CP nonconserving amplitude $\mathrm{Re}B_-$ will be of order $\varepsilon^{-1}\mathrm{Im}A_-$ making the necessary cancelation in Abs L_- possible. Clearly, there are many such CP nonconserving schemes [13 - 15].

Other possible ways out of the $K_L^0 \to \mu^+ + \mu^-$ puzzle can be thought of. For example, one might speculate on the possibility that the K_L^0 has some hitherto undetected decay modes and that the contribution from these new intermediate states in the unitarity relation cancels against the two phton contribution [16-17]. However, any such proposal must of course explain why these new states have not been detected and this is not so easy.

Finally, one might question the validity of the $\pi\pi\gamma$ and 3π contributions to the absorptive amplitude. In the case of the $\pi\pi\gamma$ state the estimate would have to be low by an order of magnitude [4,5]. For the 3π contribution to be large enough the $\gamma + \gamma \to 3\pi$ cross section would need to be about four orders of magnitude larger than expected [18,19]. These possibilities therefore seem rather unlikely.

We conclude this discussion of the $K_L^0 \to \mu^+ + \mu^-$ puzzle with a remark on the implications of the small experimental upper bound on $K_L^0 \to \mu^+ + \mu^-$ for higher order weak interactions. From the analysis by Ioffe and Shabalin[20] of the (divergent) second order weak interaction contribution to $K_L^0 \to \mu^+ + \mu^-$ one now finds the cut-off is limited to

$$\Lambda \lesssim 6 \text{ GeV} \tag{45}$$

If the cut-off were larger the most divergent part of the second order weak $K_L^0 \to \mu^+ + \mu^-$ amplitude, which is CP conserving, would exceed the experimental limit.

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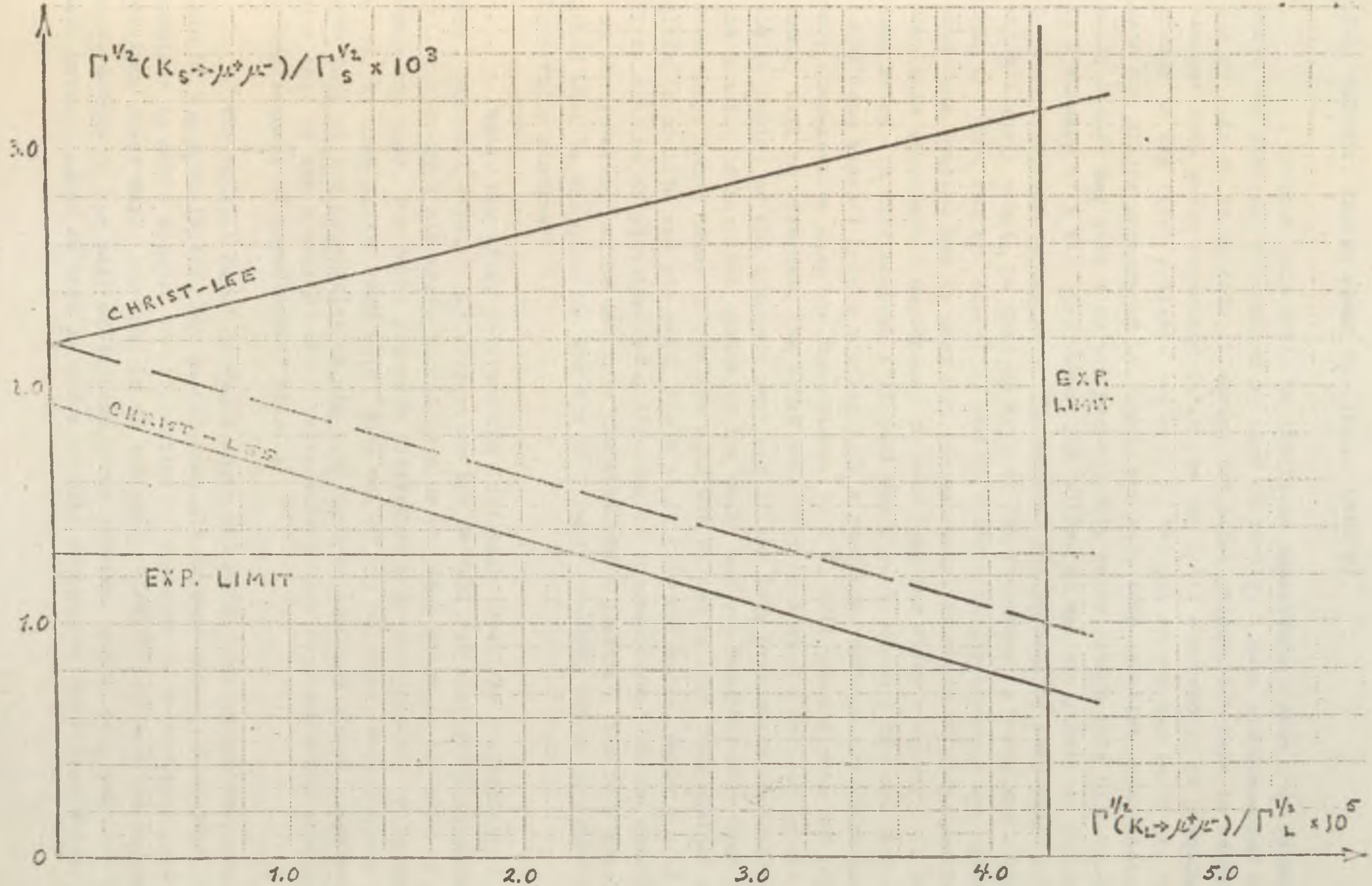

FIGURE 1. Theoretical Bounds and Experimental Limits.

MUON PHYSICS

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I thought this would be the first conference I have ever talked at where the chairman pronounced my name correctly. I have mispronounced my name all my life so as to make it easier for other people, but here in Hungary I thought they would pronounce it in the way it is supposed to be pronounced. Accent on the first syllable and open i: Telegdi so you can learn it. Gell-Mann has always pronounced it right, but he's very snobbish about these things. Well, I must say that I do not particularly enjoy giving this talk because it is a summary of a situation where you either do not see puzzle or you do not have any progress to report on the puzzle. The summary of this talk, with which I shall start, is that we know of two electrons, a heavy and a light one, the muon and the electron, and that increasingly accurate measurements have done nothing but to bear out the essential identity of these particles. Now, this statement I am suppose to make <u>last</u> for fifty minutes, as I have just heard. In this context I'll tell you a joke, which I did not dare to tell in Holland where I gave a similar talk a year ago, because the Dutch are much more conservative than the Hungarians. So I'll tell you the joke and hope nobody will be offended. The story goes that in a girl's college in the US at the beginning of the academic year one of the lady professors decided to give the benefit of her rich advice to the girls and got together with them in the dormitory. Among other things she said she wanted to warn them against the pitfalls of life and she said. "Look at it this way girls. Maybe you'll have fifty minutes of pleasure and a lifetime of consequences." At the end there was the question period and only one girl had a question. The lady-teacher asked what is the question and the girl said: "Madame, how do you make it last fifty minutes?"

Well, now after this analogy I'll talk about the properties of the muon. Before I discuss the properties let me tell you that every experimentalist who loves his trade should take an enormous pride in how much we know about the muon. The muons are made artificially by accelerators while electrons, as Dick Garwin once said, come out of any hot wire. Nevertheless we know many of the properties of the muon with an accuracy almost comparable with that of the electron, which represents a very remarkable and inspiring accomplishment of experimental physics.

The "wallet card" of Fig.1 shows the properties of the muon. The spin of the muon is 1/2. How do you know it? Because you see beautiful doublets in mesic x-rays. The statistics of the muon is Fermi-Dirac. How do you know that? Well, you know it in two ways. First of all one has accurate measurements of the pair production cross section. Better yet, you have to have several muons at once in order to make statements about statistics.

There are tridents which are the production of muon pairs by muons. I'll come back to that. And from the analysis of those one concludes that the statistics are Fermi-Dirac.

The mass of the muon is probably the best measured and least understood quantity in physics. Here it is. (There is one number I have not carried with me which is the most recent one.) How do you know that? Here you make a very precise measurement of a certain transition energy in a muonic atom; a very nice experiment, which gives the mass to a hundred parts per million. Muon physics has reached this stage, where we measure things in parts per million. Then from the ratio of the precession frequencies of muons and protons plus a knowledge of the ratio of the magnetic moments and of the magnetic moment anomaly, we know the mass to about 13 ppm. From recent experiments on muonium in conjunction again with the anomaly we know the mass at the present moment to 4 ppm.

Next, the electric dipole moment. Of course the muon has no electric dipole moment, but what limits do we have for the dipole moment? It is (5 ± 5) e cm 10^{-17}. This is obtained as a by-product of one of the Geneva (g-2) experiments.

Finally there is a remarkable quantity, the anomaly in the magnetic moment (g-2), which as I said is known to 260 ppm, (the upper number is the most recently completed CERN experiment, and underneath is the most recently completed theory). At a previous conference there was a larger discrepancy between experiment and theory, because the theory had not been carried out to that accuracy. This is a good agreement and I shall speak later about the next CERN experiment. The next CERN experiment now under construction will cut the error in the anomaly to about 20 ppm. Let me remind you that the electron anomaly is known to 3 ppm. It was published before the theory was completed and is in perfect agreement with the theory up to α^6 terms.

The life-time of the muon is a quantity which one knows better than one part in a thousand by direct measurement. The reason that this number was measured to such high accuracy by 3 different groups is because the measurement of the life-time of the muon is the Millikan experiment of weak interactions. As you measure the charge of the electron for electromegnetism you measure the lifetime of the muon to get the fundamental coupling constant of the V-A theory. Otherwise nobody would care. Of course you cannot really extract the coupling constant from the life-time without the mass, but we know the mass very accurately.

Now the only subject, where I can give you some news of possible progress (but not much - nothing of the class we have heard from prof. Reines), is on the mass of the mun - neutrino, which is in a terribly lousy condition. The best limit of a finished experiment that I am aware of is 1,5 MeV at a 90% confidence level, and is obtained by John Peoples in an unpublished thesis at Columbia University. This is to be compared to the 60 eV we attribute to the electron-neutrino. This is a completely open field of experimentation, and I hope to be able to go through some of the experiments now being actively contemplated.

I shall now talk a little bit about muon electrodynamics (See Fig.2). I shall try to do it very quickly. First, the (g-2) experiments. If g was equal to 2, the spin and the momentum would precess with the same frequency in an external field. Since g is not equal to 2 the spin precesses faster because it is larger than 2 and you measure how fast a longitudinal polarization becomes transverse. The first experiment was done in CERN at 1960, and various groups have continued ever since. Now here comes an other piece of information, which is the hyperfine structure of the muonium. I shall come back to muonium in more detail later, since it has some bearing on weak interactions. This hyperfine structure splitting is the energy difference between the two opposite spin orientations in the ground state of the μ-e atom. It is known at the present moment from a recent experiment in Chicago to 1/2 ppm. Now you may think that this is of no interest but let me point out that since the Rydberg constant, the fundamental unit in atomic physics in frequency units is known to 0.3 ppm, this number is either a source of knowledge of α or a confirmation of quantum electrodynamics. I really doubt that one should be optimistic like some people and say that the agreement of this number with theory shows that the muon has no very strange properties. Fig. 2 shows also the previous number, accurate to 2,3 ppm, obtained at Yale.

Now I go to the muon tridents. This is an experiment done by Tannenbaum at BNL. This experiment is parametrized so as to show whether we have Fermi-Dirac statistics or not. In the final result Fermi statistics correspond to parameter $\alpha = +1$, Bose statistics to $\alpha = -1$. The experiment fits the data with $\alpha = 0.98 \pm 0.23$, which gives you a rather clean confirmation of Fermi-Dirac statistics. Now the last point here is μ-π scattering. Since one has done a lot of e-π scattering, one is doing μ-π scattering, and for the most recent elastic results obtained by Lederman and co. again at Brookhaven you get a cut-off energy of 3 BeV. So there too there is nothing anomalous.

Of course people at conferences like to hear rumours and incomplete things and there is one I can tell you, the vacuum polarization. The muon in an atom is 200 times closer to the nucleus than an electron with the same quantum numbers would be. We all know that the electrostatic field of the nucleus is modified by virtual pairs over a distance of 1 Compton wave length. So in a heavy muonic atom the entire orbit is in a region where the potential is not 1/r, but something like $\frac{e^{-kr}}{r}$ because of the vacuum polarization. This effect is very, very small in ordinary atoms, but very noticable in muonic ones. I think the first person who really attracted mankind's attention to it was Pomeranchuk, but there have been many papers. E.g. for the transition of 400 Kv in lead the vacuumpolarization of order α is 2 Kv, or almost 1%. Then there are terms of order α^2 and finally of $(Z\alpha)^3$, which in the case of lead is about 50 eV.

In a recent experimental paper by Professor Anderson and his collaborators, known in the literature as Dixit et al, there were found rather serious discrepancies between theory and experiment. (These experiments of Dixit et al claim much greater accuracy than previous ones). A large fraction of that discrepancy is due to the fact that the $(Z\alpha)^3$ vacuumpolarization contribution, a term due to kroll and Wichmann, was taken over by Dixit et al from a theoretical paper where it had an incorrect sign. If it is applied with the correct sign, the discrepancy does not disappear, but becomes less flagrant. So much for this sensational piece of information.

The next point I would like to discuss is the decay of the muon. Now in this process one can measure a certain number of parameters. The electron momentum spectrum has been known for 20 years. There is a parameter ρ, that gives the shape of the spectrum, then there is a parameter for the polarized spectrum (this means that you consider the spin of the muon and the direction of the electron). That parameter is called δ. At low energies there is also a parameter η. So you have a spectrum, say,

$$S(x) = M(x;\rho\eta) + \xi B(x;\delta)\cos\theta$$

where $x = P/P_{max}$ and θ is the angle between electron momentum and muon spin. The η is important only at low x where there are practically no decays. So now we have defined parameters, and in fact you all know that the muon decays like $1 + a\cos\theta$ where this gross asymmetry a is proportional to ξ, while δ defines the shape of the energy dependence of asymmetry a. Of course we have one more parameter, which is the helicity h of the electron. So we have a total of ρ and η and δ and ξ and h, 5 parameters. And here is a point I always try to make. Every one of these parameters, which have been measured, agrees perfectly, or at least as well as one can hope, with the predictions of the V-A theory. But what is not true, is that by knowing that these parameters agree with V-A, you are authorized to conclude that the interaction is V-A. This is the usual game of necessary and sufficient. And that is a point worth noticing, because sometimes when you make a discovery somebody finds loopholes in a theory which you accepted. And this theory is the credo of everybody, including myself.

I have prepared for you Fig. 3 to show that there is interrelationship between social development, I mean large scale development of society, and physics in that sense that one can use scientific results, results of measurements, to predict, or retrospectively to establish revolutions. What you have to do is this. Take some physical quantity, let us call it capital Ξ, just for fun, which looks like no letter at all, and then you plot this quantity - let it be the velocity of light or the fine structure constant or the Michel parameter-as a function of time. If you see a sudden discontinuity in the experimental value - which obviously is invariant - then at that time there must have been a great revolution. Here on the figure you see the Michel parameter as a function of time (time is Anno Domini), and the religiously accepted value for V-A is 3/4; you see that all kinds of measurement which were done before 1957 spread and then they all converge. The sudden break in 1957 - and this is a point of a revolution - the point of discovery of V-A, where people knew what they were supposed to find, and they found it. There was a credo and so we could believe. The figure also illustrates the value of international scientific collaboration. Around 1964 two experiments were published, rather accurate ones, which disagreed with the official value. One experiment was done in the Soviet Union, the other one was done in the west, in Great Britain. And international collaboration: the mean value is 3/4. Well, you see again another small effect: this is the Chicago experiment which is almost on the red line and this is the Columbia experiment exactly on the red line. So much for that.

Well, now we go into the song and dance about how to conclude whether we have V-A or we don't have V-A. From the experimental data on the basis of the work of a lady called Jarlskog, a Persian woman married to a Swede, a former student of Kallén, there are two roads to determine V-A, (Fig 4). Road 1 is the one that we shall follow, and it is to the left. This road is full of assumptions, and later we shall go to the other side. We take the general local μ decay interaction for massless neutrinos so it is already not so general. Then we have the parameter η which governs the low energy behavior of the spectrum and that η obeys a law all by itself. It is less than $(1-h^2)$ so if the helicity h is 1, η is zero. That's a general theorem, due to a former student of mine by the name of Fryberger. Then you have ρ, δ, ξ, h, η, all ≤ 1. These are all general theorems. We now assume, for aesthetic reasons, that the neutrinos are two-component. We then have the theory which we call V-εA, i.e. there are two ingredients V and A, and ε could possibly be constant. If we make this assumption, then we find that $\delta = 3/4$ $\xi = -h$ and $\eta^2 = (1-\xi^2)/4$. A very great advantage of having made this assumption is that you can actually analyze the data, since all precision experiments of muon decay are in fact exercises in checking radiative corrections. Unfortunately, these radiative corrections can be computed only in the V and A case - in others we don't know how to do them. So it is very convenient that these things work. If you now assume T invariance, then ε is real and then you do more experiments to find that $\xi = 1$. Or, without T invariance, you have to find experimentally $\xi = 1$ and $\eta = 0$. In both cases you arrive to V-A.

Now, the road taken by Mrs. Jarlskog was to write the interaction in the charge retention form (μe) $(\nu\nu)$, so you can do the trace calculation of the neutrinos quicker. She says: Assume you have done all the experiment, and find $\rho = 3/4$, $\delta = 3/4$, $\xi = 1$, $\eta = 0$.
That means $V+\hat{A}$, where $\hat{}$ means that V and A are defined in the charge retention ordering. That automatically implies that $\eta = 0$ and you have finite radiative correction. And, now comes the impossible task. You now have to measure the spectra and polarizations of both neutrinos, and only after you have done that can you conclude thet the interection is V-A. It is very simple: in μ^- decay we have three final state particles, and only one of them is really observable; therefore you cannot find the correct theory by performing a complete set of experiments. So, that's the situation. Since several of the authors of V-A are sitting in the audience, I would like to add that the above very beautiful philosophy omits one thing: internal consistency. Consistency is that in the dynamical theory of weak interactions there is a connection between the life-time of O^{14} nuclear beta decay and muon decay. That connection just wouldn't be there without the V-A therory, and therefore this left-handed road is perhaps not as unnatural as I tried to make it appear.

Let me quickly show the present values for the μ^- decay parameters (Fig.5). These were compiled by a student of Roger Hildebrand, Derenzo and they are published. For the experimentalist I would like to say that strangely enough, the helicity of the electron is not terribly well known. Nobody is interested enough to do the hard work. Now if you want to make fits to this business to see how much you can allow for other interactions and if you normalize to one unit the vector interaction, then you find that for S and P (which are zero in the V-A theory of course by definition) the fit is ≤ 0.33. A (which should be 1) is between 1.2 and 3/4, the phase (which in V-A is π) is $180^{\circ} \pm 15^{\circ}$ and the T (that should be 0) is less than 0.28. So the situation is not so glorious. But if you look at the neutron decay, where much more work has been done, from this point of view it is not so glorious either. This is always swept under the rug. There is marvellous consistency with V-A, but never compelling evidence.

Next, we have two experiments which involve the muon in a peripheral way; the $K_L \to 2\mu$ experiment of Berkeley, which gives unbelievably small rates as compared to the theoretical expectations, and the $K_S \to 2\mu$ which should perhaps occur at a much higher rate than one thought. This problem is under active work. Let me tell you two things, that have not so much to do with muons but with K_L. The Berkeley experiment on $K_L \to 2\mu$ which shows that this rate is much lower than expected is in my personal opinion a flawless experiment, I see absolutely no source of mistake, and nor do many other competent people to whom I have talked. The $K_L \to 2\mu$ will be repeted by J.W. Cronin et al. within the next year,at a level 10 times more sensitive, than the original Berkeley experiment.

The CP violation we all know about, there is a little bit of folklore and rumor. You know that one of the effects of the CP violation is the charge asymmetry of leptons. You go very far away and count how many electrons are positive and negative and how many muons there are positive and negative. The electron experiment was done first by Steinberg et al., the muon experiment was done by Schwartz et al. There has been a somewhat embarassing tendency for the charge asymmetry of the muons to come out higher than the charge asymmetry of electrons and people who look for exciting things have not failed to notice this difference. People have to make their living and like to get excited. The experiment is now being repeted among other things by Schwartz, at Stanford, and it turns out that these muon charge asymmetries were affected by a large number of systematic effects, the most subtle of which is the range difference between the μ^+ and μ^- which Fermi once considered in 1953 but nobody thought about very seriously. And when you allow for these effects and a few others, the current status is that the muon does give the same charge asymmetry and one understands why the former value was high.

Hyperon decays. Well, whatever little one knows about, $\Lambda \to e$ $\Lambda \to \mu$ beta-decay is consistent with universality, but a high statistics in this field means 1000 events. So that is not the place to look for it.

Then we have the subject of muness, or muon quantum number. The ways of talking about lepton quantum numbers were very lucidly explained by the chairman of this conference, prof. Marx, and he mentioned the problem of additive versus multiplicative quantum numbers. We have this beautiful potential idea due to Pontecorvo I believe, of a conversion of muonium into antimuonium. I have never wanted to do this experiment because I did not see an intelligent way of doing it, but there have been ways of doing it, perhaps not too effective. The coupling constant for this transition is known to be not more that 6000 times larger than the vector coupling constant, which is hardly any information at all. It is more of a humorous nature I would say. From the experimental point of view, the problem is the following. Muons we got to burn but we have the problem of stopping them in vacuum to form muonium. And at the moment you form muonium, which you have to do by extracting an electron from an atom, you are in a gas, and the collisions with the gas atoms under almost any reasonable conditions induce more transitions than you care for, and this removes the degeneracy between the two systems. The basic idea of this experiment that I report here was a very elegant one, namely, when the muonium by its oscillation turns into the antimuonium it has two characteristics. (1) an unforeseen negative electron, instead of the high-energy positive electron from usual muon decay. (2) when you make the muonium in Argon gas, the μ^- cascades towards the nucleus and produces mesic X-ray characteristic of Argon, that you would have never gotten out of a positive muon. There is a number of schemes under way to study this thing, there has been a recent experiment in Virginia trying to make muonium by letting it exude into vacuum from gold foils. The experiment in my opinion does not give any particularly reliable answer at all, because I don't think they even proved that any muonium comes out from thet foil, and before you know what the probability of conversion of muonium is you have to prove that there is muonium there in the first place. So no particularly striking progress has been made.

Now we come to the measurement of the mass of the muon neutrino, or as I like to call it, the neutretto. The most obvious way of measuring the mass of the muon-neutrino is to look at the $\pi \to \mu$ decay. We know the mass of the muon now to 4 ppm, and we know the mass of the pion, with a question mark, to 50 - 65 ppm also from mesic X-rays. There we just measure the range, in the old fashioned way or with pulse head analysis in a germanium crystal or something of that sort. But none of this does you any good. Because if you take the current error in the pion mass, which is 58 ppm, if you believe it then that error because of a quadratic relationship that you have in a two-body decay, will give you a 700 kV bound for m_{ν_μ}. Now I have divided the experiments into two categories, class A and calss B. Class A is an experiment where the neutrino or neutretto momentum is zero, and Class B is a class of experiments where it is not zero. Now, since in the rest frame of the pion the momentum of the neutrino is manifestly not zero, you have to go to some moving frame to make it so. There have been a large number of very ingenious suggestions in the past of taking a pion in flight and choosing that pion energy at which the neutrino, now massive, emitted backward, is transformed into its rest frame. However, these experiments do not stand the acid test of statistical analysis.

Now in the class B experiments, where the neutrino has a finite momentum, there are two proposals currently investigated, or may be three, One is the reaction $\mu + Li^6$ going into 2 tritons plus a neutrino. You are supposed to look at the Dalitz plot of this two tritons, and see how far it extends. The most favored configuration (that anybody would understand) is when the two tritons go back and they share the energy. One of the problems with this reaction is that nobody has ever seen it. There is no reason that it should not be there, but nobody has ever seen it. A second problem with this reaction is that even if it exists, that particular corner of the Dalitz plot where the two tritons will come off with equal energy is a small phase space proportion and it is very sparsely populated. Quite a number of people have realized that the endpoint of a radiative decay is much better than a two-body decay. The quadratic dependence is not present in a three body system. So you study K or the π, into $\mu\ \nu\ \gamma$ and you have to measure the gamma ray energy to a high accuracy, which would mean with germanium. The energies are of the order of 40 MeV, which is not so good for germanium. No matter what you do you don't seem to have enough solid angle to do the experiment. Nobody is trying to do this experiment. The other experiment, almost equally horrible, is to study the $\pi\ \mu\ \nu$ decay in flight; with sufficient accuracy measurements of momenta in excess of 10^{-4} you can extract the mass. In particular they have a trick to send the pions and the muons through the same magnets on similar orbits. This is not a value judgement this word here: SIN. This is just the place where they plan to do this experiment: Schweizerische Institut für Nuklearforschung. If you look at the proposal, if everything works just perfectly, then they will get neutrino mass to perhaps 200 KeV. It will be a long and heroic effort which is not guaranted to succeed. The chief proponent of this experiment is Dr. Hoffe in Bern, who happens to be a personal friend of mine.

Somebody from the audience: Not everybody takes things so personally.

I told him it was a very difficult experiment. He is much younger that I am and he is much more daring, that is all. Last, but not least I would like to draw your attention to an experiment, which is connected only in a rather indirect way with neutrinos, but it is certainly connected with muons because the object of measurement are muons, and this is the Kauffel or Utah effect. Dr. Cassiday from Utah is here, and he will give a talk on this experiment. So I am really doing a commercial here for his talk. It is a very interesting experiment and I think you should listen what he has to say.

In conclusion I would like to say, I do not know any really mysterious experiment involving muons at this time, the muon neutrino mass is the object of great efforts but it is not clear when the efforts will pay off. I presume the main reason that anybody invited me anyway is not so much to give this talk, but so that I should come back to my own country. So if you didn't think my talk was very informative, please forgive me.

PROPERTIES OF THE MUON

Property	Value	How Determined
Spin	½	Mesic X-rays
Statistics	Fermi-Dirac	Pairs, Tridents
Mass	$206.76(02)\ m_e$	Energy of 3D-2P transitions in muonic phosphorus[2]
	$206.765\ m_e$	δ_μ/δ_p and $(g-2)$
Magnetic Moment	$\mu_\mu/\mu_p = 3.18333(4.8)$	from precession in H_2O (13 ppm)
	$= 3.18333(15)$	muonium (5ppm)
Electric Moment	$d_\mu < (5\pm5)ecm \times 10^{-17}$	byproduct of old g-2
Anomaly $a=(g-2)/2$	(1) $116616(31) \times 10^{-7}$ (270 ppm)	Latest CERN experiment
	(2) $1165564 \begin{pmatrix}+14\\-.7\end{pmatrix}$	Theory to order a^3 is $(a/\pi)^3\ 2.1 \begin{pmatrix}+11\\-0.5\end{pmatrix}$
	(1)-(2): $\left(61^{+31}_{-34}\right) \times 10^{-8}$ 2.8σ(7%)	
Lifetime:	$t_\mu = 2.198(1)\mu sec$	Direct measurement
Muon neutrino mass	1.5 MeV (90% CL) (!)	Endpoint of Michel Spectrum

Figure 1

2. Muon Electrodynamics

a. (g-2) Expt.

b. Muonion hyperfine structure

$\Delta\nu = 4.463305\ (10)$	Chicago double resonance
$= 4.463248\ (31)$	Yale zero field resonance

c. Muon pairs

QED verified?

d. Muon tridents 2×10^3 km mμp Tannenbaum
11 BeV/c BNL

Comment $\alpha = +1$ FD
exp. 0.98 ± 0.23
$\alpha = -1$ BE

e. μp scattering $\lambda^{-2} = 0.1 \pm 0.03\ BeV^{-2}$
$(\lambda \sim 3\ BeV)$

Figure 2

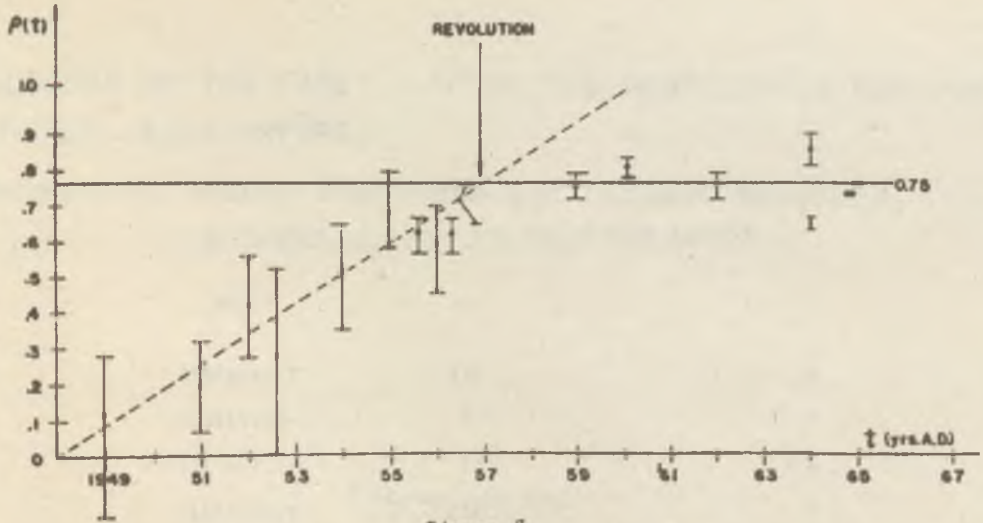

Figure 3

Experimental Value of Michel Parameter

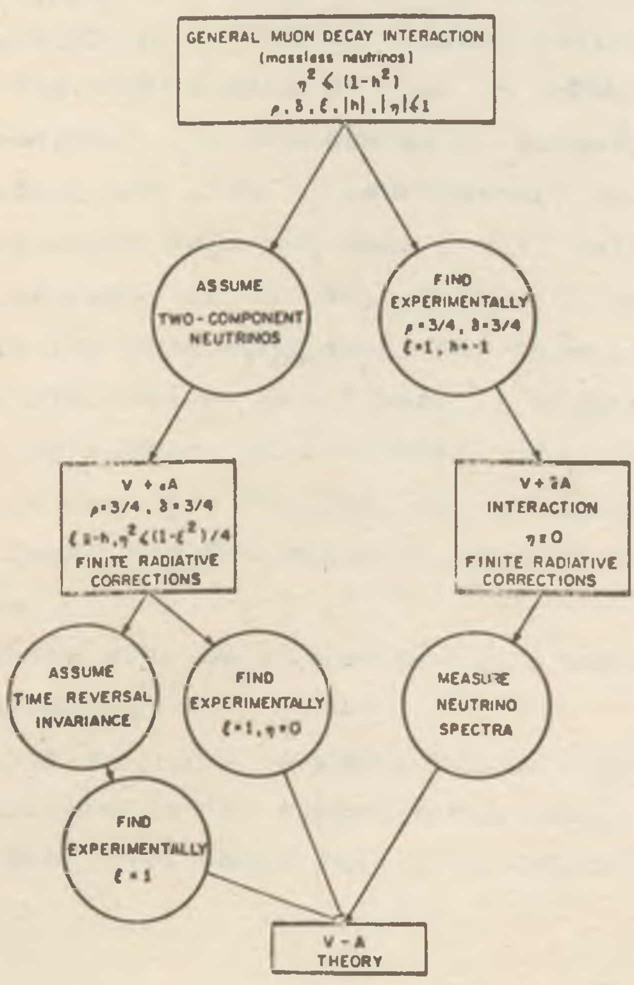

Figure 4

Logic of Experimental Test of V-A Theory

PRESENT BEST VALUES FOR MUON DECAY PARAMETERS

	V-A	EXP.
ρ	3/4	0.7518(26)
η	0	-0.12(0.31)
ξ	-1	-0.973(14)
δ	3/4	0.7540(85)
R	1	1.00(13)

COUPLING CONSTANTS (CHARGE RETENTION ORDER) IN UNITS OF G_V

	V-A	FIT
S.P.	0	≤ 0.33
A(= x)	1	$0.75 \leq x \leq 1.20$
ϕ(AV)	π	$\pi \pm 15^{\circ}$
T	0	≤ 0.28

Figure 5

CONCLUSIONS OF THE FIRST PART OF THE NEUTRINO'72 EUROPHYSICS CONFERENCE, BALATONFÜRED

B. Pontecorvo, Joint Institute for Nuclear Research, Dubna

Introduction

There have been already so many talks and discussions at our conference that it is really difficult for me to make a summary. Notice that most of the talks had a review character. This makes my task even more difficult. And then it occured to me that when I will be back in Dubna my friends will ask me: how about the conference in Hungary? I will have to answer them and give my impressions. So I decided that right now I would tell you what I will tell the people in Dubna in a few days. Of course, one has to use some subjective criteria for selecting among the papers and the comments in the discussion. Now I wish to emphasize that the omission in my summary of communications presented here does not mean in any way that they do not deserve being mentioned. The fault for the omission must be traced to the criteria I used:

First, there will be an unmistakable experimental bias in my summary.

Secondly, I am going to spend little time on problems about which there are no new experimental data, even if they are very important. This seems reasonable since such problems, as a rule, have been treated recently in review papers. Thus I neither will talk about "second class" currents, which were discussed in the interesting talk of Dr. Pietschmann, nor about muon physics, which was treated by Telegdi in a beautiful lecture rich of "first class" jokes.

Thirdly I am going to talk only about ambitious and difficult investigations in which somebody tries very hard to find and measure something, but does not see anything. You certainly have noticed that at our conference most of the experiments had such character. Search experiments usually give results which very improperly are called "negative". At the conference there were presented brilliant and brave experiments, yielding very low upper limits, the significance of the results being very great. However, the fact remains that results are presented not in terms of a measured quantity being equal to a certain value but through the $<$ sign. I must say that this is becoming more and more frequent, and that is one of the reasons why life is much harder for people doing experiments than for theoreticians.

Thus our conference, at least the first half of it, is an "inequality" meeting, where new effects were searched for at an incredible sensitivity level. On the basis of this "inequality" principle I am lead to make the following classification of the material in the summary:

1) The $K_L \to 2\mu$ puzzle (plenty of inequalities)
2) Solar neutrino's (upper limit)
3) Lepton charge conservation (upper limits)
4) "Stable" heavy leptons ("negative" results)
5) Antineutrino electron scattering, reactor + electronics (upper limit)
6) Neutral currents, accelerators + bubble chambers (upper limits)

The $K_L \to \mu^+ + \mu^-$ puzzle

The Oakes report and its discussion were terminated only a few minutes ago, so that it would be tiresome for you if I were now continuing to talk in detail about the $K_L^0 \to 2\mu$ puzzle.

When correcting the magnetophone tape, however, I decided that I had to write something on the puzzle, mainly in order to mention such points in the discussion which were not known, at least to me, before our conference.

As is well known, the puzzle consists in the following (see also the extensive review paper of Dolgov, Okun and Zakarov, ITEP, No. 924): according to experiments

(1) $$\left(\frac{\Gamma_{K_L \to 2\gamma}}{\Gamma_{K_L \to all}}\right)_{exp} = (5 \pm 1) \cdot 10^{-4}$$

(2) $$\left(\frac{\Gamma_{K_L \to 2\mu}}{\Gamma_{K_L \to all}}\right)_{exp} \leq 1.8 \cdot 10^{-9} \quad \text{(90\% confidence level)}$$

Theoretically one can obtain the following lower limit:

(3) $$\left(\frac{\Gamma_{K_L \to 2\mu}}{\Gamma_{K_L \to 2\gamma}}\right)_{theor} \geq 1.2 \cdot 10^{-5}$$

From (3) and (1) one gets

$$\left(\frac{\Gamma_{K_L \to 2\mu}}{\Gamma_{K_L \to all}}\right)_{exp.theor.} \geq (6.0 \pm 1.2) \cdot 10^{-9}$$

in contradiction with (2).

As you know, the puzzle is a serious thing. The reason is that the theoretical bound (3) is quite reliable: one calculates only the imaginary part of the amplitude, which arises from transitions on the mass shell $K_L \rightarrow$ intermediate real states $\rightarrow 2\mu$; the real part of the amplitude can only increase the probability transition and this is why you get a lower limit of the rate. The two photon intermediate state, which has been observed experimentally, should dominate the immaginary part of the amplitude, because other states have much less space phase.

In the discussion Marschak (as well as Okun and collaborators in the quoted paper) reminded us that frequently in the past puzzles arose from wrong experiments. However, as Telegdi has emphasized, the best experts think now that the experimental results (1) and (2) are correct. If the experiments are right, the puzzle must be solved. Many theoretical proposals have been made in which some kind of cancellation of the 2γ imaginary part of the amplitude is invented "ad hoc". Most of these proposals (violation of CPT invariance, violation of unitarity of the S-matrix, introduction of new particles with "necessary" properties) are neither attractive nor plausible,in the opinion of many physicists.

As you know, Christ and Lee, instead, made a proposal which is quite attractive. They note that because of the usual (small) CP violation, $K_L = K_2 + \varepsilon K_1$ (with the usual notations). To suppress the $K_L \rightarrow 2\mu$ decay these authors assume that the necessary cancellation is due to the $K_1 \rightarrow 2\mu$ decay. Since ε is small, the $K_1 \rightarrow 2\mu$ amplitude must be much larger than the $K_2 \rightarrow 2\mu$ amplitude; in addition in the $K^o \rightarrow 2\mu$ decays there must be a strong CP violation, in order that the final states in $K_1 \rightarrow 2\mu$ and

$K_2 \rightarrow 2\mu$ may interfere to cancel the two photon contribution. The analysis by Oakes of various theoretical bounds and experimental limits has shown that an experiment designed to detect the $K_S \rightarrow 2\mu$ decay if its relative probability $\frac{\Gamma K_S \rightarrow 2\mu}{\Gamma K_S \rightarrow all}$ is larger than 10^{-7} would either confirm or exclude definitely the Christ and Lee schema.

Now we heard at the conference that several experiments are being performed at present to detect the $K_S \rightarrow 2\mu$ decay at the necessary sensitivity level. According to an information of Telegdi one of these (at CERN) has already given one good $K_S \rightarrow 2\mu$ decay candidate! Similar experiments will be performed at the Argonne Laboratory and within a year there should be a definite answer. It was also very interesting to hear from Telegdi that an experiment designed to detect the $K_L \rightarrow 2\mu$ decay if its relative probability is larger than $2 \cdot 10^{-10}$ is being prepared by the Croning group. This is an order of magnitude below the previous result (2) of Clark et al.

Now two words about a comment by Marschak. Since he will give his full talk later at this conference, I will not go into the real business now, limiting myself to few remarks. The work of Marschak and collaborators on the "strong cubic intermediate vector boson" model is not new (1969), but it seems that its relevance for the $K_L \rightarrow 2\mu$ puzzle was not generally recognized and became clear only at our conference. I would say that the proposed theory is in fact a model of the phenomenological Christ and Lee proposal. In this sense even if the model does not appeal to some people on estetical grounds, it is certainly of great interest, since it is not an "ad hoc" proposal, and "required" the Christ and Lee schema even before the $K_L \rightarrow 2\mu$

puzzle exploded. Other good things of the model, in my opinion, are: first, the fact that the well known "weak" CP violation is a consequence of the theory, where basically there are strong CP violations, and secondly, the predictions of gross CP violations in the production of intermediate bosons, and in various processes, such as $K^+ \to \pi^+ e^+ e^-$, $K^+ \to \pi^+ \mu^+ \mu^-$ and, of course, $K^o \to 2\mu$.

Solar neutrino's

1) The Brookhaven experiment

Now I am going to talk on solar neutrino's and I will spend more time on this problem than on any other problem. Of course I will start with the famous Brookhaven National Laboratory experiment. This is a brave experiment of Davis et al. and I would say it is one of the few experiments which are being performed without competition. The same can be said of the Reines experiment on $\tilde{\nu}_e - e^-$ scattering with reactor antineutrino's, about which I will talk later. Suppose that Davis were feeling like going in vacation for a few years: there is no the slightest risk that somebody else would perform an analogous experiment in the meantime!

I will say few words to convay to the theoreticians the difficulty and the scale of this experiment. Deep underground there is a mass of about 600 tons of C_2Cl_4. If events of the reaction

$$\nu_e + Cl^{37} \to A^{37} + e^-$$

are originated within such "swimming pool" by solar neutrino's, you get a few A^{37} atoms. This is a radioactive noble gas and it is possible to extract these few atoms from the tremendous amount of C_2Cl_4 with a small amount of argon carrier and by He purging. The argon fraction is then separated again from the large helium one and is introduced inside a small proportional counter, in which one measures the characteristic energy emitted in the K capture decay of A^{37}. If you understand the difficulty in pushing the effective background of the counter down to about one count per month, then you will realize how acrobatic is this experiment.

My impression is that the necessary checks were done carefully and, as a person who has been working quite a lot with proportional counters, partially in order to prepare an experiment similar to the one which Davis is doing, I am very impressed by the improvement in the counter used at Brookhaven.

The improvement is due to the fact that, in order to decrease the effective background, also the pulse time of rise is measured (in addition to the pulse amplitude spectrum). This gives a substantial rejection factor for pulses originated by (background) particles, which inside the counter are less localized than the Auger electrons from K capture in A^{37}.

Here are the results:
the A^{37} production rate is 0.18 ± 0.10 events/day. In part this rate is due to the muon background (muons produce protons which produce A^{37} via the reaction Cl^{37} (p,n) A^{37}). This background, partially measured and partially calculated (see the talk of Young) is 0.12 ± 0.04 events/day. The difference is 0.06 ± 0.14 (it is safe to add the errors).

Thus solar neutrino's have not yet been detected. The capture rate of solar neutrino's in the detector is less than one in five days (70% confidence level). This corresponds to a rate $\leq 10^{-36}$ sec^{-1} (Cl^{37} atom $)^{-1}$ = 1 SNU (solar neutrino unity, according to a convenient notation of Bahcall).

There are two astrophysical conclusions which can be made with reasonable certainty. As it will be seen below, most of the expected counting rate from solar neutrino's in the Brookhaven detector is due to high energy neutrino's from B^8 (very few but very effective). So the first conclusion made by Davis is that the Sun emits much less B^8 neutrino's than expected. The other conclusion is that the C-N cycle is of little importance in the Sun, since otherwise the A^{37} counting rate should be much higher.

2) Interpretation.

The interpretation of the Brookhaven experiment was given at the conference by Bahcall.

The thermo-nuclear reactions of the hydrogen cycle in the sun are shown below together with their expected relative percentage:

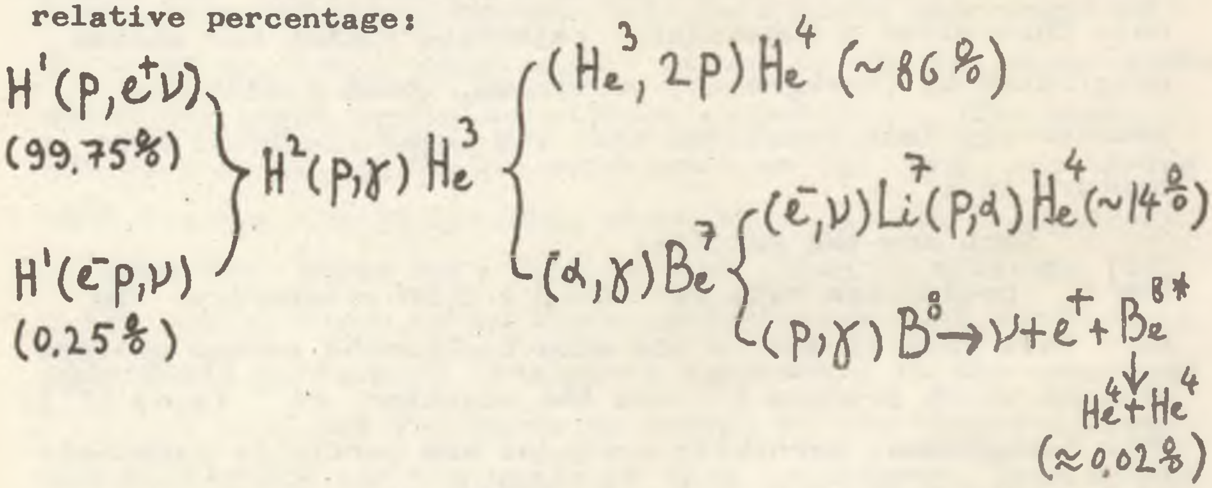

Fig. 1.

In the following table, presented by Bahcall in his interesting talk, are given the expected capture rates due to neutrino's produced by different nuclear reactions in the Sun , and expressed in SNU.

You will see the serious discrepancy already mentioned between the measured rate and the theoretical expectation, which is the summ of all the contributions in the table

(a possible small contribution from the C-N cycle is not included).

Neutrino source	Maximum neutrino energy (MeV)	Expected capture rate in the Cl-A detector (SNU)
$H^1(p,e^+\nu)H^2$	0.42	0
$H^1(e^-p,\nu)H^2$	1.44 (monoenergetic)	0.26 ± 0.03
Be^7	0.86 (90%) 0.38 (10%) (both monoenergetic)	1.0 ± 0.2
B^8	14.1	4 - 7

Is the discrepancy serious enough to force us to draw revolutionary conclusions about the Sun or about the neutrino properties? My opinion is: no. Let us look at the table. Most of the expected A^{37} rate is due to B^8 neutrino's, which represent a very small fraction of the total number (B^8 neutrino's are very energetic and consequently very effective). The reactions leading to the production of B^8 are quite unimportant from the point of view of the structure of the Sun. The A^{37} rate due to B^8 neutrino's was calculated by using currently accepted solar parameters; well, astrophysicists will have to change these parameters, and the Sun will nevertheless shine as before; seriously speaking, I mean that the Sun with new parameters will not substantially differ from what we think it is now. For the time being there is no astrophysical tragedy: the Brookhaven result is very important, since it will help to change the current solar parameters in the right direction. I think that this is not far from the opinion of Bahcall, although may be I am more conservative than he is.

After the Bahcall lecture there has been a comment by Chudakov. He said that in Moscow Kopysov and Fetisov suggested that may be the low flux of solar B^8 and Be^7 neutrino's is due to the existence of a resonance in Be^6: the existence of a resonant Be^{6*} state will decrease the concentration of He^3, and consequantly of Be^7, which arises in the $He^3(\alpha,\gamma)\,Be^7$ reaction, and of course of B^8 (see Fig. 1).

I think that this possibility is rather unlikely, since the resonance must be at the needed energy, but the proposal is very interesting and reasonable. After the comment was delivered Chudakov found out that this suggestion had already been made by Fowler and asked me to apologize for him: he just did not know about the Fowler work. Anyway the suggestion is an instructive example of "what might be true" : the Sun structure practically would not change, while the flux of B^7 and B^8 neutrino's would substantially decrease.

Now if we look at the remaining (other than B^8) contributions in the Table, you will see that their summ is not in serious disagreement with the result of Davis ($\lesssim$1 SNU, 70% confidence). Thus I repeat the conclusion: there is no reason to think that the Sun is substantially different from what we believe it to be and, even more emphatically, there are no reasons to believe that the neutrino's have very exotic properties.

3) The future.

Now let us see how the future looks like. Prof. Davis is going to get very interesting results in the near future using his improved counter and probably will either detect solar neutrino's or get to the limit his experiment permits ($\sim$ 0.5 SNU at the given depth underground). As Bahcall pointed out, if you yet a sensitivity of $\sim$0.3 SNU and you still do not see neutrino's you really have got something very exotic. This is so because the expected rate for pep neutrino's ($\sim$0.3 SNU) is known quite well. As a matter of fact the total flux of solar neutrinos is obtained directly from the Sun luminosity and from the basic fact that there are liberated $\sim$ 25 MeV energy, when 4 protons are transformed into one α particle $+ 2e^{+} + 2\nu_e$. Unfortunatcly it is not a simple matter to reach the sensitivity 0.3 SNU, especially at a depth of "only" 4300 m H_2O eq., where now is located the Brookhaven detector.

What to do in the future? Prof. Davis told us about a very important new detector of solar neutrino's which he is developing: it is a Li compound from which it is possible to extract chemically a volatile compound of Be, which can be introduced into a counter. This is the beginning of a promising development, since the reaction $\nu_e(Li^7, e^-)Be^7$ is capable of detecting pep neutrino's.

Now the more remote future of solar neutrino astronomy, in my opinion, is connected with the development of huge liquid or solid (noble gas?) electronic detectors, capable of giving some information on the energy and the direction of the detected neutrino. But I will not elaborate on that.

4) Exotics.

If neutrino's will be missing at the level expected for the $H'(p,e^{+}\nu)H^2$ or $H'(pe^{-},\nu)H^2$ reactions one has to invent something more or less extraordinary.Consequently exotics is useful since it makes us ready for the worse. The danger arises if you believe really in extraordinary things even before you are forced into exotic by hard facts. There were several exotic suggestions at our conference.

In his talk Landé suggested: the neutrino's may not be here now, but in the past they were, because the Sun may be is pulsating. He than discussed the relation between the neutrino history of the Sun and the thermal history of the earth ("neutrino archeology").

Bahcall, Cabibbo and Yahil said: if neutrino's are missing, may be they decay on their way from the Sun to the earth. It is a simple explanation, if explanations are needed. Such a speculation, as it turned out, has been already useful, inasmuch as it stimulated an experiment: if the neutrino with mass$\neq$0 decay into another particle ν' neutrino-like with zero mass + a photon, one can try to detect the photons near a working reactor. During the discussion Reines told us that he has performed such experiment. For the particular decay $\tilde{\nu}_e \rightarrow \nu' + \gamma$ it was found that the $\tilde{\nu}_e$ decay path is larger than 10^5 astronomic units. Again an inequality! One can ask: why do you need such experiment? My opinion is that any correct experimental measurement is always a very respectable thing. The exotics is useful.

Let me say now a few words about the problem of neutrino oscillations, about which there was quite a lot of discussion at our conference. Oscillations were proposed and studied in Dubna, Moscow and Leningrad because they give a very sensitive method for investigating the question about possible lepton charge violations and the neutrino mass problem. The relevance of the oscillations to the interpretation of future solar neutrino experiments was immediately recognized, but I wish to emphasize that the oscillations were not invented "a posteriori" to solve the "missing neutrino puzzle".

It is argued that lepton charge non conservation and a finite value of the neutrino mass may lead to oscillations of the type $\nu_e \rightleftarrows \nu_\mu$, similar to the $K^\circ \rightleftarrows \tilde{K}^\circ$ oscillations in kaon physics. Other types of oscillations ($\nu_e \rightleftarrows \tilde{\nu}_e$ etc) can be ruled out if in nature there exist only 4 neutrino states.

Since the problem of neutrino oscillations was discussed in detail in my report at the 1970 Kiev Conference, I will not elaborate furthermore and only state some results:

1. The presence of oscillations will decrease by a factor 2 the number of neutrino detectable in a solar experiment (since half of neutrino's are sterile). Only under special exotic conditions or in very sofisticated and remote experiments can the "decrease factor" may become $>$ 2.

2. The existence or absence of oscillations could be established in various ways, the simplest method being the comparison of the measured and expected capture rates of solar neutrino's from the pp or pep reactions. As far as the problem of the neutrino mass is concerned, the oscillations are present if the mass difference of the two Majorana neutrino's which enter the theory is larger than 10^{-6} eV. This method is several millions of times more sensitive than the ordinary one of measuring the neutrino mass (sensitive to mass values larger than $\sim$ 10 eV, in the most favourable case of the tritium decay). The physical reasons why the method is so sensitive are a) the possibility of measuring an amplitude (and not a squared amplitude) and b) the huge distanced which characterise the solar system.

Lepton charge conservation

1) Double β decay

There was a very interesting review of the subject by Fiorini who presented also a beautiful experiment, done under the Mount-Blanc, searching for the process $Ge^{76} \rightarrow e^{-} + e^{-} + Se^{76}$. A Ge(Li) crystal ($\sim$70 cm^3, $\sim$400 gr) was used both as a source and as a detector. The pulse amplitude spectrum is measured with high resolution and one looks in the region around 2.045 MeV, which is the expected sum of the energies of the two electrons in the process looked for. No peak appeared at 2.045 MeV, where the background count was only 2 $(keV)^{-1}$ in 1000 hours. In other energy regions you see lots of peaks due to very minute impurities of natural radioactive elements, so you are confident that the experimental arrangement is working properly. The experiment is absolutely convincing and gives as a result another inequality for the double β decay of Ge^{76}:

$$T_{1/2} > 4.5\cdot 10^{21} \text{ years} \quad \text{(68\% confedence level).}$$

The lepton charge violating amplitude is at most one percent of the lepton charge conserving amplitude. Similar results were obtained previously by different techniques in the search of neutrinoless double β decay of Ca^{48}, Se^{82}, Te^{130}.

2) Search for the $\mu^{+} \rightarrow e^{+} + \gamma$, $\mu^{+} \rightarrow e^{+} + e^{+} + e^{-}$ processes

Concerning the muon charge violation, I am going to report now on the work of the group of Korenchenko, who at the Dubna synchrocyclotron looked for the mentioned processes. They use a cylindrical spark chamber magnetic (9200 - 4500 Oersted) spectrometer to analyse the muon decay products (see the Table below).

As for the $\mu \rightarrow e\gamma$ process, the accuracy of the result is comparable with the one obtained previously by different methods. But it may be of interest to you that Korenchenko is planning now a search for the $\mu \rightarrow e\gamma$ process where a braching ratio of $(3 - 6)\cdot 10^{-10}$ should be measurable. As for the $\mu \to 3e$ process, the result

Searched for process	Number of π^+ stopped in target	Number of photographs	Registration Efficiency	$\Gamma_{process}/\Gamma_{\mu\to all}$ (90% confidence)
$\mu^+ \to e^+ \gamma$	$6\cdot 10^9$	$2.5\cdot 10^5$	1.35%	$< 2.9\cdot 10^{-8}$
$\mu^+ \to e^+ + e^+ + e^-$	$2.9\cdot 10^{10}$	$6.0\cdot 10^5$	3.0%	$< 3.2\cdot 10^{-9}$

is ~ 40 times better than the best previous one. It is seen that the muon-charge violating amplitude can hardly be greater than 1% of the normal amplitude. The results are comparable in accuracy withthose obtained in the double β -decay investigations, but of course not only the processes but also the lepton charges which are investigated are different.

3) Multiplicative lepton charge?

I would like to mention something new I heard at our conference on the question as to wether one of the lepton charges is a multiplicative number. As you know, several proposals for experiments on this point were made long ago and I will not mention them here.

Now the IHEP-ITEP collaboration (Arbuzov et al) proposed recently for the NAL program to search for the reaction $\tilde{\nu}_\mu + e^- \to \mu^- + \tilde{\nu}_e$, which, if observed in a large H bubble chamber, would directly prove the existence of a multiplicative lepton number. At Batavia the $\tilde{\nu}_\mu$ energy is more than sufficient and the proposal, in my opinion, is the best for the solution of the multiplicative lepton charge problem.

Talking of the multiplicative lepton charge, I would like to mention also a comment made by Filippov for the benefit of experimentalists. He will talk later about his model of a "four dimensional symmetry with multiplicative lepton charge" (whatever that means), but has already presented in the discussion the predictions of his model, which are:

1) $\sigma(\tilde{\nu}_\mu + e^- \to \tilde{\nu}_e + \mu^-) = \frac{1}{2}\sigma_{V-A}(\tilde{\nu}_e + e^- \to \tilde{\nu}_e + e^-)$

2) $\sigma(\nu_\mu + e^- \to \nu_e + \mu^-) = \frac{1}{2}\sigma_{V-A}(\nu_\mu + e^- \to \nu_e + \mu^-)$

3) $\sigma(\nu_e + e^- \to \nu_e + e^-) = \sigma(\nu_\mu + e^- \to \nu_\mu + e^-) = \frac{1}{2}\sigma_{V-A}(\nu_e + e^- \to \nu_e + e^-)$

4) Theory

Prof. Marx gave an interesting review of the subject and I am unable now to go into it. I would like to mention, however, that in his talk he proposed a decrease with time of the strenght of the CP violating interaction. Irrespective of the arguments given by Marx, the idea seems to me interesting; the possible change in time of constants was discussed before, but the CP violating constant seems to me an attractive candidate for the following reasons: if nature worked that way, which is of course highly improbable a priori, the big bang approach would easily give the asymmetry between matter and antimatter even if the barion number in the Universe is equal to zero (and without the need of inventing new particles, as it was done in the papers of Zakharov and Kuzmin).

Stable heavy leptons ?

In the report of Gershtein on work done at Serpukhov by the Landsberg group there was discussed an experiment designed to detect heavy, stable $(\tau > 10^{-9} \mathrm{sec})$, charged leptons. These objects, with charge $\pm$ are supposed to be very similar to muons; they are not interacting, so they can be detected as muons usually are. The cross section for their pair production in collision of protons with nuclei is supposed to be due only to their electric charge, and can be calculated on the basis of the Lederman experiment on production of muon pairs by protons. Such charged heavy leptons were looked for but not seen at Serpukhov. A comparison with the theoretical expectation, normalized to the muon pair production data, permits to draw the following definite conclusion: there are no "stable" heavy leptons with masses in the interval 1-3.5 GeV. The reason why you get a definite statement is that a production rate of heavy leptons equal to the expected one (normalized to muon pair production) could have been measured easily if such object existed.

Neutrino scattering

Now I will turn to the subject of neutrino scattering, that is neutrino lepton and neutrino nucleon scattering. I must say that the Weinberg's theory has a very progressive influence on the work. We are now seeing a sort of renaissance of the weak interaction phyiscs and this, to a definite degree, is due to his theory. I heard that at the Tashkent's Conference Pais qualified Weinberg's work as "strategy". I like this definition. The old problem of neutral currents is now investigated experimentally on a very wide scale with the help of reactor and accelerator facilities, and most experiments are being interpreted in terms of Weinberg's theory. Certain experiments, for example the search for $\nu_\mu - e^-$ scattering (at least with electronics methods) would probably not be considered without the new theoretical encoutagement. Now such experiment is one of the first on the list at pion factories and is being performed now in large bubble chambers.

1) Antineutrino-electron scattering (reactor + electronics)

The experiment, being conducted by Reines and collaborators with the help of a large reactor, occupies a central place in our conference. Reines has been working on the problem for more than ten years. The investigation is very difficult and, as I said before, the experimental arrangement is a monopoly of Reines. Why the experiment is so difficult? An elastic collision between $\tilde{\nu}_e$ and e^- at reactor energy is an event without a very characteristic signature, and is imitated easily by background, for example, by a Compton electron generated by γ's from radioactive impurities etc. So the fight against background is the main problem and is made by clever and complicated methods, which I cannot describe now. The detector itself is a ~ 8 Kg plastic scintillator, which sounds very simple but it is not. In order that you may be appreciate the difficulties, I remind you that the counting rate for the events of interest, that is for events which could be electron recoils with energy $>$ 3.5 MeV in the reaction $\tilde{\nu}_e + e^- \rightarrow \tilde{\nu}_e + e^-$, is about one a day! This is the rate with the reactor on as well as with the reactor off, that is here again it was not possible to see what one was looking for.

The upper limit of the cross section σ_{exp} for the process $\tilde{\nu}_e + e^- \rightarrow \tilde{\nu}_e + e^-$ with fission $\tilde{\nu}_e$ is found to be

$$\sigma_{exp} < 1.7\ \sigma_{V-A} \qquad (70\%\ \text{conf.level})$$

Here σ_{V-A} is the expected cross section according to the V-A prediction (only charged currents) and the V-A spectrum of electron recoils was assumed when calculating the fraction of electrons with energy $>$ 3.5 MeV. Now in the Weinberg's theory the presence of neutral currents, to an extent depending upon the parameter e^2/g^2, changes the V-A prediction; the results of Reines give already some constraints on the Weinberg's parameter. This was also discussed in the report of Baltay.

We heard from Reines that he has been recently improving the experimental arrangement, doubling the mass of plastic scintillator without any increase of the background. He feels confident that within a year ha could "see" the scattering events if $\sigma_{exp} > 1/3\ \sigma_{V-A}$. Let us wish him success.

2) Neutrino Electron and Neutrino Nucleon Scattering. Neutral currents (Accelerator + bubble chamber)

We heard yesterday in the talk of Pullia (the data were obtained at CERN in the "heavy" bubble chamber Gargamelle exposed to ν_μ and $\tilde{\nu}_\mu$) and in the talk of Baltay (a critical and instructive review of all the data) about the state of the search for neutral currents in $\nu_\mu - e^-$ and $\nu_\mu - N$ scattering and also in other processes. Work is being done on a very wide front; sufficient to say that only at CERN the number of neutrino events obtained with Gargamelle is at least an order of magnitude greater than all the world pregargamelle statistics; the analysis of the data, however, for the time being is very preliminary and partial. The Weinberg strategy dominates. Baltay emphasized that the energy spectra of the final state electron or hadrons in processes due to neutral currents strongly depend upon the parameter of Weinberg theory e^2/g^2. This means that the detection efficiences in the experiments (and consequently the upper limits obtained) depend on e^2/g^2. The results of such analysis of Baltay are illustrated in self-explanatories figures in his report, and I will only say a few words of summary on the neutral current question (I will call symmetrical the neutral currents of the type $e\bar{e}$, $p\bar{p}$, $\nu\bar{\nu}$ ... etc. and asymmetrical the neutral currents of the type $e\bar{\mu}$, $n\bar{\Lambda}$ ... etc.).

I. There is strong evidence against asymmetrical neutral lepton currents (for example the Dubna work on the absence of processes like $\mu \to e\gamma$, $\mu \to 3e$ etc.).

2. There is strong evidence against asymmetrical neutral hadron currents (of course this comes out from the absence of certain kaon decays such as $K_L^o \to 2\mu$, $K^+ \to \pi^+ + e^+ + e^-$ etc).

3. There is no experimental evidence in favour of symmetrical neutral lepton currents (e.g. $\nu_\mu + e^- \to \nu_\mu + e^-$) nor in favour of neutral symmetrical hadron - lepton currents (e.g. $\nu_\mu + p \to \nu_\mu + p$, $\nu_\mu + p \to \nu_\mu + n + \pi^+$ etc; such processes can hardly have a cross section larger than 1/10 of the cross section corresponding to charged currents).

The predictions of the Weinberg model (requiring neutral currents, at least symmetrical ones), however, are very close to be tested. The game is only starting, but within less than a year we should have an answer. I should notice, however, that even if the Weinberg model with neutral currents will be excluded, the Weinberg's strategy will stay: I am told that B.Lee has shown that this is so in a model with heavy leptons.

In conclusion I wish to express my warmest gratitude to the Hungarian Academy of Science, to the Hungarian Physical Society and to Prof. Marx for the wonderful hospitality.