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FINAL REPORT

University of Oklahoma High Energy Physics Group
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Detectors at Fermilab and Cornell (Tasks A&C); Non-Perturbative Quantum
Field Theory (Task B)"

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CONTENTS

Experimental Task A: D0 experiment at Fermilab	page 1
Theoretical Task B: OUHEP Theory	19
Experimental Task C: CLEO experiment at Cornell	50
Experimental and Theoretical Task D: Neural Nets	52
Experimental Task E: ATLAS experiment at CERN	53

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1995-1998 Close-Out Report

University of Oklahoma

Effort on the DØ Experiment

1 Introduction

During the period 1995-1998, members of the University of Oklahoma (OU) high energy physics group, contributed to the success of the DØ experiment in a number of ways. During this period of time, the members of the OU group were:

P. Gutierrez	Associate Professor	DØ Contact
M. Strauss	Assistant Professor	
G. Kalbfleisch	Professor	
G. Guglielmo	Research Associate	
E. Smith	Graduate Student	
G. Steinbrueck	Graduate Student	
T. McMahon	Adjunct Assistant Professor	Langston University
J. Snow	Adjunct Associate Professor	Langston University

During this period of time, the OU group contributed to both run I (data taking and data analysis) and run II (hardware and software). The contributions for run I data taking and processing are:

1. Maintaining level 2 trigger simulator;
2. Running the monte carlo farm for the *top* and QCD groups;

while the contributions to the physics output of the experiment are:

1. Search for charged Higgs in *top* decays;
 - (a) Contributed to indirect search;
 - (b) Lead effort on direct search;
2. Contributed to QCD with *WZ*;
 - (a) Contributed to *W* and *Z* p_T measurement;
 - (b) Contributed to measurement of ratio of *W* plus zero jets to *W* plus one jet;
 - (c) Lead the effort on the *e* angular distribution measurement as a function of p_T from *W* decay.
3. Contributed to the measurement of the anomalous *Zγ* coupling;
4. Lead on the measurement of the direct photon production cross section at $\sqrt{s} = 640$ GeV;
5. Made minor contributions to the search for scalar leptoquarks.

In addition to the work done for run I, the OU group contributed to the detector upgrade for run II. This work has included the following:

1. Silicon vertex detector hardware:
 - (a) Testing and characterizing silicon detector sensors;
 - (b) Numerous contributions to production and testing of detector elements.
2. Contributions to reconstruction software and vertex trigger;
 - (a) Contributed to global track reconstruction;
 - (b) Contributed to vertex finding;
 - (c) Contributed to secondary vertexing studies.

The following are brief reports on some of the projects listed above.

2 Physics Analysis

2.0.1 $t \rightarrow H^\pm$ Analysis

The minimal extension to the Higgs sector of the Standard Model involves the addition of another complex scalar doublet [10]. In the model considered, one doublet couples to the up-type quarks and leptons, while the other couples to down-type quarks and neutrinos. After symmetry-breaking, one is left with five physical Higgs fields: two charged scalars, two neutral scalars, and one neutral CP odd field. Flavor changing vertices once associated only the $W^\pm$ are now found in the couplings of the charged Higgs, $H^\pm$, as well. If the $H^\pm$ exists, one could find it in processes which normally occur through a weak channel.

The DØ collaboration has made a complete study of the acceptance and backgrounds to the process $t \rightarrow Wb$. If DØ were to measure a top quark production cross section which was very low compared to that predicted by theory, one might suspect that the top was decaying through a channel other than $t \rightarrow Wb$, and the products of that decay were being missed. On the other hand, if one could account for the entire theoretical cross section for top production by measuring the weak decay channel only, then a theory which predicted additional modes *might* be ruled out. This is the strategy behind the indirect search for $H^\pm$.

In the model considered (which happens to be the same as that for MSSM), one has two complex scalar doublets in the Higgs sector, with opposite weak-hypercharges:

$$H_1 = \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix}, \quad H_2 = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix} \quad (1)$$

The ratio of the vacuum expectation values of H_1 and H_2 define a parameter, $\tan(\beta)$. If one considers only this extension to the Higgs sector (ie no other MSSM couplings), the branching ratio of $t \rightarrow Hb$ and the branching ratio of the subsequent decay of the Higgs into it's final states are determined completely by $\tan(\beta)$. The top group completed their analysis using the assumption that the branching ratio of $t \rightarrow Hb$ is zero. The indirect $H^\pm$ search concentrated on the standard model mode:

$$t\bar{t} \rightarrow WWbb, \quad WW \rightarrow \text{jet} + \text{lepton} \quad (2)$$

If two scalar doublets existed, this final state would also contain, for example:

$$t\bar{t} \rightarrow HHbb, \quad HH \rightarrow c\bar{s}\tau, \quad \tau \rightarrow \text{lepton} \quad (3)$$

and the exact contribution expected from this mode could be calculated using $\tan(\beta)$. The detector acceptance, as determined by monte carlo, to these new modes given the top group's standard lepton + jets cuts would provide a prediction of the number events measured by the top group, given the existence of $H^\pm$.

To this end, a series of monte carlo experiments were run at various values of M_H and $\tan(\beta)$. $t\bar{t}$ events were generated, and decayed into W, H with a branching ratio appropriate to the value of $\tan(\beta)$ under consideration. The $H^\pm$ were also decayed according to $\tan(\beta)$. The final states were then kept as good events, or thrown out with a frequency consistent with the acceptance for each particular decay, as determined by the top group, and monte carlo studies for charged Higgs. Background events were also added using Poisson statistics, such that the average number of background events was 11, the top group's estimate for background to their $t\bar{t} \rightarrow \text{lepton} + \text{jets}$ data sample. A distribution of the predicted number of events was generated for various m_t, M_H , and where the number of events observed by the top group was beyond the 95% confidence limit as predicted by the simulated experiments, points in the M_H vs $\tan(\beta)$ plane were excluded. The first draft of a paper to be submitted to PRL, is in the process of being review by the DØ top group.

In addition to the indirect search described above, a direct search for $t \rightarrow H^+$ is being carried out. In this case, we look for the appearance of a charged Higgs signal in the background as opposed to looking for its absence in the top signal. Because of the large QCD backgrounds present in the all jets decay mode, we will concentrate on decays that have lepton + jets in the final state. This is a region of $\tan(\beta)$ in which $BR(H \rightarrow \tau) \approx 1$. One important strategy in the rejection of background, then, is the ability to identify τ decays. The efficiency of the τ id, and the fake rate from b jets has been determined for the case of:

$$t\bar{t} \rightarrow HHbb, \quad HH \rightarrow \tau\tau \rightarrow \text{jet, jet} \quad (4)$$

and determinations of $W + \text{jet}$, $Z + \text{jet}$, and $t\bar{t}$ backgrounds are forthcoming. When the study of rejection of background by τ id is complete, the focus will be on the training of a neural net in the recognition of Higgs events.

Currently we are completing a study on the τ finding efficiency, and background to this signal. A preliminary result should be available during the spring of 1999.

2.0.2 Electron Angular Distribution versus p_T^W

One of the first measurements performed on the W after its discovery was to show its left-handed coupling to fermions along with establishing that it has a spin of one [11]-[12]. This was accomplished by measuring the angular distribution of the charged lepton which follows the well know $V - A$ form $(1 - \cos \theta^*)^2$, where θ^* is measured in the rest frame of the W with respect to the direction of the hadron with the same charge as the charged lepton. All these measurements have been performed for low values of the transverse momenta where the dominate production mechanism is $\bar{q} + q' \rightarrow W$.

At higher values of the of the W transverse momenta (p_T^W), other processes contribute, such as [13]:

1. $\bar{q} + q' \rightarrow W + g$
2. $q + g \rightarrow W + q'$
3. $g + g \rightarrow W + \bar{q} + q'$

where only the first two contribute at Tevatron energies. None the less, these two processes change the angular distribution of the emitted charged lepton, which now becomes:

$$\frac{d\sigma}{dp_T^2 dy d\cos\theta^*} = K_{\theta^*} (1 + \alpha_1 \cos\theta^* + \alpha_2^2 \cos^2\theta^*) \quad (5)$$

where the various parameters depend on p_T^W and W rapidity. Further to reduce the effect of the unknown longitudinal momentum of the neutrino, the angle θ^* is measured in the Collins-Soper frame [14]; this is the rest frame of the W where the z -axis bisects the angle formed by the proton and minus the anti-proton momenta. This frame is chosen since it reduces the ambiguity on the neutrino longitudinal momentum to a sign ambiguity on $\cos\theta^*$.

We have started on the measurement of the angular distribution of electrons from W decays, as a function p_T^W . The data set used comes from run 1B, which corresponds to about 80 nb^{-1} , or to approximately 41,000 W events in the central calorimeter (we plan to eventually include the forward calorimeter into the analysis). Before looking at the data, we performed a number of monte carlo experiments to obtain an estimate of the level of statistical accuracy that we expect to obtain from the measurement. We found that for data out to $p_T^W = 200 \text{ GeV}$, we should be able to differentiate between pure V-A and V-A with next to leading order QCD at the 3 standard deviation level.

Given this result, we have been looking at the data. At present a preliminary measurement is being reviewed by the collaboration for possible presentation at the DPF conference in January, 1999. Beyond this, there is additional fine tuning of cuts, better understanding of backgrounds and efficiencies, and the inclusion of electrons in the forward region before this data is ready for publication.

2.0.3 Direct Photon Analysis

We are completing the analysis necessary to publish a measurement of the cross section for production of direct photons in $p\bar{p}$ collisions. Such a measurement probes the parton-parton interaction without the ambiguities associated with jet fragmentation and energy measurements. At large center of mass energies such as those found at the Tevatron, the dominant production mode for photons with low transverse energy is through gluon Compton scattering. Therefore, the direct photon production cross section is sensitive to the gluon distribution within the nucleon, and a measurement of this cross section can be used to test next to leading order QCD.

High energy measurements of the direct photon production cross section were initially made by the UA2 experiment at CERN in 1991 at $\sqrt{s} = 630 \text{ GeV}$ [15]. The UA2 results showed an excess of photons produced at low E_T^γ . Later measurements at $\sqrt{s} = 1.8 \text{ TeV}$, which have never been published, also show this same excess [16]. The results from UA2 and CDF are shown in Figure 1.

During Run 1, the DØ experiment at Fermilab has observed $p\bar{p}$ collisions at energies of both $\sqrt{s} = 1.8 \text{ TeV}$, and $\sqrt{s} = 630 \text{ GeV}$. The 1.8 TeV direct photon production measurement by DØ

Figure 1: Production of direct photons at 630 GeV from UA2 and CDF. The curves show the theoretical predictions using NLO calculations and NLO calculations with parton showers.

has been published, and does not show any photon production excess at low E_T^γ [17]. An updated result using the entire data set from Run 1 is nearing publication, as well. By performing this measurement at $\sqrt{s} = 630$ GeV and comparing the results with those of previous experiments and with theory it should be possible to determine if there is truly an excess of photon production, or if the cross section is explained by current understanding of parton interactions within the protons.

The experimental differential cross section is given by

$$\frac{d\sigma}{dE_T d\eta} = \frac{NP_\gamma}{L\Delta\eta\Delta E_T A\epsilon} \quad (6)$$

where N is the number of observed photons, $\Delta\eta$ is the range of pseudo-rapidity, ΔE_T is the range of transverse energy, L is the integrated luminosity, A is the geometric acceptance, P_γ is the photon purity, and ϵ is the product of various experimental efficiencies. The main source of experimental error comes from the error on the photon purity. DØ has pioneered a method of determining the purity, which uses the ratio of energy deposited in the first electromagnetic layer (E_1) of the calorimeter to the energy deposited in the total calorimeter (E_{total}). Monte Carlo distributions of $\log(E_1/E_{\text{total}})$ for both signal and background events are used to form a function which is fitted to the data distributions. Events from single photon production tend to have lower values of $\log(E_1/E_{\text{total}})$ than background events from π^0 and η decay.

Publication of the single photon production cross section at 630 GeV with the DØ detector should occur very soon. Preliminary results have been obtained, and we are currently investigating the low energy trigger efficiency using Monte Carlo events. All other aspects of the analysis have been completed including an estimation of the photon purity and the systematic errors involved, an accurate determination of the integrated luminosity at 630 GeV and the DØ tracking efficiency. An internal DØ note has been written, though not yet distributed to the collaboration.

This publication should help resolve the discrepancies between the current published data and the theoretical predictions, as well as between the various experimental results. In addition, our measurement will be the first measurement of the single photon production cross section in the rapidity region greater than 1.5 at $\sqrt{s} = 630$ GeV.

2.0.4 $Z \rightarrow \mu^+ \mu^- \gamma$ Analysis

A search for anomalous $ZZ\gamma$ and $Z\gamma\gamma$ couplings, using run 1b data, is currently being concluded. This search provides an important test of the standard model, since no tree level couplings are predicted [18]. Never the less, the process $p\bar{p} \rightarrow Z\gamma X$ does occur in the standard model through initial state radiation, off one of the incident quarks, in Z production. This analysis is carried out using the reactions $p\bar{p} \rightarrow e^- e^+ \gamma X$ and $p\bar{p} \rightarrow \mu^- \mu^+ \gamma X$. This provides for two independent measurements of the same parameters and in the end, a factor of about two increase in statistics. The data sample, as stated earlier, comes from run 1b data, which corresponds to $\approx 97 \text{ pb}^{-1}$ —a previous limit from run 1a [21, 22] data ($\approx 14 \text{ pb}^{-1}$ has been published.)

The data is analyzed in terms of the most general ($ZV\gamma$) couplings that are consistent with Lorentz and Gauge invariance [19, 20]. This leads to four independent trilinear couplings, with two of these couplings leading to CP even amplitudes ($h_{i=1,2}^V$) and two giving CP odd amplitudes ($h_{i=3,4}^V$). Further, to ensure partial wave unitarity at high energies, the couplings are given in the form $h_i(\hat{s}) = h_i/(1 - \hat{s}/\Lambda^2)^{n_i}$, with $n_1 = n_3 = 3$ and $n_2 = n_4 = 4$. To extract the four parameters, the theory calculation is compared to the observed photon p_T spectra. This provides the most sensitive test, since it is expected that events with non-zero $ZV\gamma$ couplings will produce an excess of events at large p_T .

As stated above, the Run 1b analysis is nearing completion. A draft of a PRL paper is currently being circulated among members of the collaboration, with an open comment period ending October 5 1997. Unless there are some unforeseen questions raised during this period, the paper is expected to be submitted to PRL shortly after the expiration of the comment period. No indication of an anomalous signal was observed and thus upper limits on the coupling strengths have been set (see fig 2). These new results have been combined with previous results from this experiment for $ZZ\gamma$ and $Z\gamma\gamma$ and represent the tightest limits currently available. This analysis has been published in [23].

3 Run II

3.1 Silicon Vertex Detector

The primary responsibility of the OU group in the DØ detector upgrade, is in the area of the silicon vertex detector [24]. In addition to being involved in all the discussions on the detector design, the group has been primarily involved in testing and characterizing individual silicon strip detectors. This includes production testing to insure that the detectors meet the DØ specifications and source testing of prototype detectors. We have also participated in beam tests of these detectors.

Members of the OU group have also been involved in the supervision of production testing that is being carried out at Fermilab. We have also participated in the process of setting up a facility at Fermilab to test fully constructed detector ladders; a ladder consists of two detectors wire bonded together with the support structure and readout chips mounted. The test facility is composed of

Figure 2: The figure shows the limits on the CP -even couplings for $ZV\gamma$. The data includes both run 1a and 1b plus the run 1a $Z \rightarrow \nu\nu$ channel.

a laser that sits over a detector on a movable table. The readout used is the final version of the SVX-II readout.

Over the last year and a half, we have been participating in the development of tracking software. This effort has and will mainly concentrate on the inclusion of silicon strip detectors into the main DØ tracking software.

In addition to the other efforts on the upgrade, we have also become involved in an effort to study and design a tracking trigger based on the silicon strip detector. This trigger will enhance the physics capabilities of the experiment, by allowing to trigger on particles with short lifetimes ($\approx 10^{-12}$ s). This gives the DØ detector the capability to do b -quark physics and improves the tagging for t -quark events. In addition, the silicon detector trigger will help sharpen up the p_T trigger threshold. This work is still at an early stage.

3.1.1 Production Detector Testing

To carry out the production testing, plus other tests on silicon strip detectors, the OU HEP group constructed a clean-room using funds from DØ. The clean room was set up and became operational in December of 1995. Detectors received from Fermilab are stored in the clean room in air-tight boxes with dry nitrogen flowing through them. The detectors under test are placed on a probe station that is surrounded by a light-tight Faraday cage; this keeps noise to a minimum while measuring the detector characteristics. Finally all data taking and control functions are performed with a Power Macintosh 7100/80 running the Labview software.

The focus of our efforts, is the testing of single and double sided silicon micro-strip detectors, which will be used in the detector upgrade. The detectors are tested against design criteria which specify full depletion voltage, active area current, coupling capacitor integrity, polysilicon resistor

Figure 3: ADC distribution of pulse height from incident α 's on a silicon strip detector versus reverse bias voltage. The maximum pulse height defines where to set the operating voltage.

values, and interstrip resistance.

The common method for determining the full depletion voltage of a silicon micro-strip detectors, is to measure the gain and phase shift between the ohmic and the diode sides as a function of reverse bias voltage. The single sided detectors supplied by Micron for the DØ upgrade, however, contain a MOS structure below the poly-silicon resistors which dominates the body capacitance up to voltages of about twenty volts, which is close to the full depletion voltage and in many cases obscures it.

To get around this problem, we developed a technique where an α -source is placed on the ohmic side of the detector and the pulse height on the diode side as a function of reverse bias voltage is measured. Because of the very short range ($\approx 20\mu\text{m}$) of alpha particles in silicon, the charge generated by the alpha is not seen on the diode side until the detector is almost fully depleted. Ideally, the charge collected goes from zero to the full alpha energy loss over a few volts. The diffusion spot on the diode side is a function of voltage even after the detector is fully biased, and as a result, the charge collected on a single strip will continue to change for voltages beyond full depletion.

Charge deposited on a capacitor from a $0.1\mu\text{C } ^{210}\text{Po}$ alpha source is used to trigger a discriminator. The discriminator output is made $2\mu\text{sec}$ wide and was used as an integration window on an integrating ADC (Camac 2249A). It is worth noting that the distributions are qualitatively the same when a peak sensing ADC is used instead. Fig. 3 shows the distribution of charge collected as a function of reverse bias voltage.

Figure 4: Maximum charge collected as a function of detector bias voltage.

This detector was found to deplete at ≈ 37 volts using the α -source. It is clear from this distributions that a change in discriminator threshold would affect a counting rate experiment. If, however, one only considers the maximum charge collected at any given voltage, the distribution would become insensitive to changes in discriminator setting or amplifier gain. Fig 4 shows this distribution.

Full depletion is taken to be the point where the maximum ADC curve turns over most quickly, which is identified by the minimum in the second derivative. The data points are too few to give a good numerical answer, so the second derivative must be calculated from a fit to the data. The function chosen for the fit is one with only one second derivative minimum over the range of the fit, and is:

$$f(v) = a + b \frac{e^{-c(d-v)}}{1 + e^{-c(d-v)}} + e \ln(v + f) \quad (7)$$

where v is the bias voltage, and $a-f$ are fit parameters. This function agrees well with all data taken. The fit to the data and the corresponding second derivative minimum are calculated in Labview, and all relevant information about the fit is written to disk for future reference.

Double sided detectors have their full depletion voltage determined using more conventional methods. The capacitance C across the detector is measured as a function of voltage. At the full depletion voltage the capacitance changes from nearly a quadratic dependence on the voltage to one nearly independent of voltage. This change in dependence is clearly evident by a change in slope in a plot of $1/C^2$ as a function of voltage. The method described above for determining the depletion voltage of single sided detectors is not as effective in determining the voltage for double sided detectors. We are currently investigating the reason for this.

Design specs for the single sided detectors require the measurement of the p-side bias current, and the p-side guard-ring current. This is accomplished with a Keithley 237 source-measure unit,

Figure 5: Setup of the I-V measurement, showing the connections of the Keithley 237, and 617.

and a Keithley 617 electro-meter. The K237 is used to bias, and measure the current through, the entire detector. The bias line current is then extracted as: $I_b = I_{237} - I_{617}$. Fig. 5 illustrates this setup. All data is written to disk for future reference. For double sided detectors, the guard ring current is not independently measured. Instead, a measurement of the entire current across the detector is made, including bias current and guard ring current.

To test all of the AC-coupling capacitors on the micro-strip detectors, the probe station movement has been automated. A VXI interface is used to control the motion of the detector in the plane of the detector, and the stage holding the probes is moved by VXI controlled motors perpendicular to the plane of the detector. Once the Labview program, which controls the motors is given the (x, y, z) positions of the first and last strips on the detector, the entire data taking process proceeds automatically. A set of fifty current measurements at a voltage of 10 Volts is made on each of the 384 strips, resulting in a measurement time of ≈ 3 hrs/detector. All current measurements for all strips are written to disk for future reference.

The coupling capacitors on the two sides of the double sided detector are tested differently from each other. The n side is tested exactly as done for the single sided detectors described above. Every channel is tested at 10 volts to determine if the coupling capacitor is shorted. However, on the p side, this test is not adequate. When the detectors are fully depleted under operating conditions, the n side of the detector will be held at ground, and the entire bias voltage will be held across the p side coupling capacitor. Consequently, the p side capacitor must remain functional up to a potential difference of 100 V or more. Each coupling capacitor on the p side is tested by raising the voltage first to 10 V, and taking 50 measurements to check if the capacitor is functional. If there is minimal current across the capacitor at 10 V, the voltage is raised by 20 Volt increments until it reaches 100 V. The current is again measured to determine if the capacitor can adequately hold 100 V. This entire process is automated. The program makes decisions based on the current being drawn, and ramps up the voltage in steps. After each step, the voltage is raised only if the current is below some threshold value. This extensive testing of the p side coupling capacitors requires about 6 hours per detector. However, because it is fully automated it can be run with no personnel present once the initial set up has been accomplished.

Measuring the interstrip resistance is a very difficult task because the resistance between the channels is much larger than the polysilicon resistors connecting each channel to the bias line.

We have tested a number of different methods for extracting this large resistance, on the order of a gigaohm, from the polysilicon resistance, on the order of a few megaohms. We first make an accurate measurement of the polysilicon resistor by holding the bias line at ground potential, and applying the same voltage (between -5.0 and +5.0 volts) on three consecutive channels. No current flows between channels (i.e. across the interstrip resistance) because adjacent channels are held at the same voltage as the channel being measured. Therefore, an accurate measurement is made of the single polysilicon resistor in the center channel of the three channels held at the same voltage. Next, the bias line and the two outside channels are held at ground, while the inside channel is raised to a different voltage. Current now flows across the polysilicon resistor on the inside channel, and across the interstrip resistance between the two adjacent channels. This current, which is the sum of the current through the polysilicon resistor and two interstrip resistors, is measured and compared to the resistance of the polysilicon resistor alone. Thus the interstrip resistance is inferred from these two measurements.

All the data collected is written to disk, so that any piece of information can be retrieved for any detector, without having to take data again. Labview programs have been written, which can extract the data that is currently considered important by Fermilab, and translate this data into a format which is readable by Microsoft Excel. The compilation of data into the spreadsheet format is fully automated.

Testing of single-sided detectors has been completed. We tested 58 single-sided detectors at OU and 62% passed all tests. The yield for the double-sided 2° stereo has not been as good. We have tested 40 double-sided detectors and only 12 have passed the original specifications. Consequently, we have relaxed some of the requirements for a good double sided detector. The number of allowed bad channels has been raised from 1% to 2%, and the alignment required between the aluminum strips and the substrate has been altered. With these modifications, 29 of the 40 double-sided detectors pass these new specifications. Testing on double-sided 2° stereo detectors, and possibly double-sided 90° stereo detectors and double-sided wedge detectors will continue here at the University of Oklahoma.

3.1.2 Test Beam

During the summer of 1997, the DØ silicon group carried out a beam test with with the following goals in mind:

- Check the performance of the silicon tracker. This included individual silicon strip detectors, with a complete readout chain. The detectors tested were single-sided and double-sided 2° stereo detectors plus a single-sided ladder. In addition, one single-sided detector that had been irradiated to 1 Mrad, was also studied. Finally, a ladder was tested in a 2 Tesla magnetic field;
- In addition to testing the performance of the silicon detectors, the performance of the readout electronics under realistic running conditions, was given as a goal. This allowed for modifications to future versions.
- The third goal was to test the data acquisition software in a real environment.

The beam test was carried out in the NW-beamline at Fermilab, using a 125 GeV pion beam. The setup consisted of a set of trigger counters to provided a beam trigger, a silicon strip detector

beam telescope (2 detector upstream and 2 downstream of the test device with both horizontal and vertical strips) to provide the beam position and a 2 Tesla magnet where the test detector was placed—the test detector could also be mounted outside the magnet. The beam telescope was built, installed and tested by the OU group, with the telescope stands being built by the OU machine shop. In addition, the OU group carried out simple monte carlo calculations to ensure that the positioning of the beam telescope was optimized within the the beam line constraints. Other DØ detector elements were placed downstream of the silicon detectors.

The data analysis is currently winding down, with the hope of having the results written up, and published within the next six months to one year. Even before the data was fully analyzed, much was learned about operating the SVX 2 and the rest of the readout. When the system was initially switched on, the amount of noise present overwhelmed the signal, but by the end of the exercise clean data was being taken. In addition, several questions were raised about the High Density Interconnect (HDI) design; this issue is being carefully reexamined in particular for the double-sided detectors. As far as the detectors themselves (results that had previously been reported):

- The magnetic field does not degrade the performance of the detectors;
- On the double-sided detectors, the n and p-side signals were shown to be correlated. This implies that the noise was low enough to not cause any problems;
- An irradiated detector was able to be operated (depletion voltage went from about 20 V to 120 V).
- The resolution was found to be $8.8 \mu\text{m}$, which is well within what DØ requires.

Recently, the OU group has been analyzing the test beam data taken at different values of the magnetic field. The goal is to measure the cluster shape and the cluster offset from the nominal zero field case. This measurement was performed at 5 magnetic field values, ranging from 0. T to 2. T in steps of 0.5 T, with the beam incident normal to the surface of the detector.

The cluster shape was calculated using the method of reference [25]. To summarize the method briefly, the strip with the largest pulse height was identified; the is referred to as the seed. Then the pulse height ratio was formed with the neighboring strip that has the largest pulse height: the distribution:

$$\eta = \frac{Q(R)}{Q(L) + Q(R)} \quad (8)$$

where Q is the pulse height, and R (right) corresponds the strip with the largest higher strip number (see fig. 6). To convert the η distribution into a distribution given as lengths, the $dN/d\eta$ distribution is integrated:

$$x = \frac{\Delta}{N_0} \int_0^\eta \frac{dN}{d\eta} d\eta - X_0 \quad (9)$$

where N_0 is the total number of events, Δ is the strip pitch $50 \mu\text{m}$, and X_0 is half the strip pitch. Given that $d\eta/dx$ is the cluster shape, the integral distribution is used to convert the original

distribution into one that involves the spatial coordinates. After the change in variables, the cluster shape is given by:

$$\frac{dN}{d\eta} = \frac{dN}{dx} \frac{dx}{d\eta} = \frac{N_0}{\Delta} \frac{1}{f(-x)} \quad (10)$$

where $f(-x)$ is the cluster shape (see fig. 7).

Figure 6: The η distribution, as defined in the text for the the various magnetic field values. The last figure shows the integrated η distributions.

To verify that we understand the experimental distributions, and for use later in modeling the effects of the magnetic field during data taking with the DØ detector, a monte carlo program was written to model the diffusion process. The first step in the model is to calculate the electric field, since the amount of diffusion depends on the drift time which depends on the field. When the silicon is fully depleted (all charge carriers removed) the bulk material has a net charge density. Assuming this charge density to be uniform, the Poisson equation is solved giving the following electric field:

$$E(y) = \left[\frac{V - V_D}{d} + \frac{2V_D y}{d^2} \right] \quad (11)$$

where V is the applied voltage, V_D the voltage at full depletion d the thickness of the detector, and y the distance the charge travels. The drift time is obtained by integrating out the equation of motion, which gives:

$$\frac{dy}{dt} = \mu E(y) \Rightarrow t(y) = -\frac{d^2}{2\mu V_D} \ln \left[1 - \frac{2V_D(d-y)}{(V-V_D)d} \right] \quad (12)$$

Figure 7: Comparison of data and Monte Carlo distributions.

Given the drift time as a function of the applied voltage, full depletion voltage, and drift distance, the amount of diffusion can be calculated. One additional effect has to be taken into account, the crossed electric and magnetic fields. The drift due to the magnetic field is parameterized in terms of the Lorentz angle as:

$$\tan \theta_L = \mu B \quad (13)$$

Solving the diffusion equation, and integrating out the coordinate along the strip, the diffusion distribution is given by a Gaussian as follows:

$$f(x) = \frac{1}{\sqrt{4\pi D_h d}} \int_0^y \frac{dy}{\sqrt{t+t_0}} \exp \left[\frac{(x - y \tan \theta_L)^2}{4D_h(t+t_0)} \right] \quad (14)$$

where the parameter t_0 accounts for the ionized charge being produced in a cylinder of finite size about the incident track.

In the monte carlo, the charge produced in the ionization process is placed uniformly across the thickness of the detector. The diffusion due to each charge is added up giving the cluster shape for

a given set of parameters. These shapes are then integrated over a set of detector strips with an added 2 counts RMS of ADC noise to compare with the experimental distributions. Figure 7 shows the experimental and monte carlo distributions, which show reasonably good agreement [26].

3.1.3 Silicon Tracking Trigger

Over the last year, the DØ collaboration has been carrying out a study of the physics gains that a trigger based on the silicon strip detector can provide. Besides the obvious gain of an impact parameter trigger, the silicon trigger can also provide a sharper trigger turn-on, for triggers using tracking information, since track p_T is improved with the inclusion of silicon detector information. As an example of possible gains, with appropriate choices of triggers, the inclusion of an impact parameter trigger can reduce the trigger rate by almost a factor of 3, while maintaining almost 99% efficiency for triggering on $t\bar{t}$ events—this study is based on a full GEANT simulation of the DØ run 2 detector. (Other studies are in progress, we quote the $t\bar{t}$ study since it is closest to being completed.)

Currently, the level 1 trigger is envisioned to use the calorimeter, fiber tracker, muon chambers and preshower detector to generate triggers. These triggers are based on single detector elements and are required to be ready in $4.2\mu\text{s}$ —the level 1 bandwidth is limited to an accept rate of 5–10 KHz. At this point the fiber trigger consists of finding tracks in predefined p_T roads. If a trigger occurs, information from the highest p_T tracks (currently 20 tracks) is passed to the level 2 trigger.

The level 2 trigger is split into two pieces, the preprocessors and the global trigger. The preprocessors use any level 1 information available and combine it with specific detector components. This is then passed on to the global trigger, which puts together different detector components to form a trigger. Each portion of level 2 has $50\mu\text{s}$ to complete its task—a bandwidth of 1 KHz is imposed on the level 2 trigger.

The final trigger level, level 3, does a partial reconstruction of the event to impose tighter trigger requirements. The event reconstruction is done in one of 50 nodes of a processor farm. At this stage, information from level 2 will be available, in addition to the raw detector data. The level 3 accept rate is set at 20 Hz

The hardware for the silicon trigger, goes in as a level 2 preprocessor. As currently envisioned, the silicon preprocessor will use the highest p_T fiber tracks as seeds to reconstruct a full track (a track with both silicon and fiber hits). The fiber track is projected into the silicon tracker. The projection is done using an approximate form of the helix equation, which is valid for p_T values greater than a few hundred MeV in a 2 Tesla field:

$$\phi(r) \approx \frac{b}{r} + \kappa r + \phi_0 \quad (15)$$

where b is the impact parameter, κ the curvature (proportional to $1/p_T$), ϕ_0 the value of ϕ at infinity and r the radial distance from the nominal beamline. The combined track is formed from those silicon hits that are within 1 mm of the projected fiber track (see fig 8). The track parameters are calculated using a linearize fit:

$$p_j = p_j^0 + \sum_{i=1}^n \left(\frac{\partial p_j}{\partial \phi_i} \right)_0 (\phi_i - \phi_i^0) \quad (16)$$

Figure 8: Distance of silicon detector hits from project fiber tracks.

where p_j corresponds to the j^{th} track parameter and the variables and parameters superscripted with a 0 correspond to a known reference track through the silicon—note that this equation is exact for a linear least squares fit. The advantage of the linearized fit is that this form of the equation is easily handled by a Digital Signal Processor (DSP), while the known reference track parameters can be stored in memory.

The current hardware design is still not settled, but will most likely consist of the following components:

1. A cluster finder;
2. A track finder;
3. A track fitter.

A preliminary design for each of these components exists.

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Theory Task B

Nonperturbative Quantum Field Theory

Final Report 1995–8

Substantial progress was made on the topics proposed in the proposal written in October 1994, and new initiatives started before the end of the grant in March 1998. These are enumerated as follows:

1. Finite-Element Method in Quantum Field Theory
 - (a) Non-Abelian Gauge Theory
 - (b) Quasilocal Formulation
 - (c) Anomalies in Quantum Electrodynamics
 - (d) Lattice Hamiltonian and Applications to QCD
2. The δ Expansion and Variational Perturbation Theory
 - (a) The Linear δ Expansion; Variational Perturbation Theory and QCD
3. The Casimir Effect
 - (a) Identity of van der Waals and Casimir Energies
 - (b) Sonoluminescence and the Dynamical Casimir Effect
4. Magnetic Monopoles
 - (a) Binding and Production
 - (b) Theoretical Aspects of the Experiment
5. Analytic Perturbation Theory and QCD
 - (a) The Analytic Running Coupling
 - (b) Inclusive τ Decay
 - (c) Deep Inelastic Scattering Sum Rules
6. PT Symmetric Theories
 - (a) Perturbative Parity Symmetry Breaking
 - (b) Supersymmetry
 - (c) PT Symmetric Quantum Electrodynamics
 - (d) Solving the Schwinger-Dyson Equations

I. FINITE-ELEMENT METHOD IN QUANTUM FIELD THEORY

Here the major progress was extending this analytic approach to lattice field theory to a non-Abelian gauge theory.

A. Non-Abelian Gauge Theory

Our major progress in 1995 was the formulation of non-Abelian gauge theory on a finite element lattice in four dimensions. (The two-dimensional formulation was given, in part, in [90].) A beginning of that extension was sketched at the Singapore ICHEP meeting in 1990 [93], and at the Quarks 92 Seminar in Moscow in May 1992. While I was on sabbatical at Imperial College, London, in 1995, I showed how the gauge-invariant field strength could be constructed in terms of the link operator,

$$(L_\mu)_{ijkl} = e^{-igh(A_\mu)_{m_1+i, m_2+j, m_3+k, m_4+l}}, \quad (1)$$

where h is the lattice constant, and the vector potential is an appropriately averaged one:

$$A_{m_\mu, m_\perp}^\mu = \frac{1}{2} \left(A_{m_\mu, m_\perp}^\mu + A_{m_\mu-1, m_\perp}^\mu \right). \quad (2)$$

The notation here is that $m_\perp$ refers to space-time lattice coordinates other than the one singled out, m_μ . Each index $ijkl$ in (1) takes on two values, 0 or 1. The result of a detailed, constructive calculation [109] is the following simple transformation law for the link operator

$$\delta(L_\mu)_{1000} = ig[\delta\omega_{0000}(L_\mu)_{1000} - (L_\mu)_{1000}\delta\omega_{1000}], \quad (3)$$

where we have assumed that the first coordinate index refers to the μ direction. Then, it is easy to see that the “transversely-local” field strength $\mathcal{F}_{\mu\nu}$ is given by the following path-ordered product of link operators around the μ, ν “plaquette”:

$$e^{-igh^2(\mathcal{F}_{\mu\nu})_m} = P e^{-ig \oint A \cdot dl} = (L_\mu)_{1000} (L_\nu)_{1100} (L_\mu^\dagger)_{1100} (L_\nu^\dagger)_{0100}, \quad (4)$$

where the first index is the μ coordinate and the second index the ν coordinate. To construct the full field strength, which is forward-averaged over a hypercube with lower left-hand corner at m we use the gauge covariant average operators constructed from the links:

$$(\tilde{D}_\lambda \mathcal{F}^{\mu\nu})_{0000} = \frac{1}{2} \left[(L_\lambda)_1 (\mathcal{F}^{\mu\nu})_1 (L_\lambda^\dagger)_1 + (\mathcal{F}^{\mu\nu})_0 \right], \quad (5)$$

where on the right side we have only displayed the λ coordinate (the rest are 0). That is, the field strength is given by*

$$(F_{\mu\nu})_{\overline{m}} = \left(\prod_{\lambda \neq \mu, \nu} \tilde{D}_\lambda \mathcal{F}_{\mu\nu} \right)_m, \quad (6)$$

where symmetrical averaging is to be understood. It is immediately obvious that $F_{\mu\nu}$ given by (6) transforms covariantly,

$$(\delta F^{\mu\nu})_{\overline{m}} = ig[\delta\omega_m, (F^{\mu\nu})_{\overline{m}}]. \quad (7)$$

Given this field strength, a covariant Yang-Mills equation can be immediately written down in analogy with the Dirac equation (multiplication being replaced by commutation). See [109].

*The averaging over the finite element on the left side of Eq. (6), required by the finite-element prescription, is necessary to ensure unitarity. If $(F_{\mu\nu})_{\overline{m}}$ were replaced by $(F_{\mu\nu})_m$, even the free theory would not be unitary.

B. Quasilocal Formulation

We have further examined the possibility of constructing a completely local set of field equations, as suggested some years ago by Matsuyama¹ [111]. The idea was to use the link variables, in terms of which the field strength was constructed, to express the Dirac and Yang-Mills equations in local form, rather than the form involving the entire prior lattice given previously. A locally gauge-invariant construction can indeed be done, one which is inequivalent to the earlier formulation. However, the new scheme is less advantageous than it first seemed, and ultimately fails to be consistent:

- The field strength which appears in the Yang-Mills equation is not locally constructed in terms of local link operators. It appears impossible to have a local formulation consistent with local non-Abelian gauge symmetry.
- Moreover, detailed calculations, even in the Abelian theory where the equations are local, turn out to be no simpler, and perhaps more complicated than those in the original formulation.
- Disastrously, unitarity is violated. Explicitly, we have seen in the Abelian case that the transfer matrix is not unitary, even in a temporal gauge. This is traced to the fact that the covariant derivative operators are not skew Hermitian. In contrast, the original formulation is manifestly unitary (canonical). It should be recalled that preservation of the canonical commutation relations at the lattice sites was the original motivation for adopting the finite-element prescription for field theory on a Minkowski lattice.

We can interpret these results in a positive light. It was always apparent that, although some arbitrary choices had to be made to implement local gauge invariance, the requirement of that invariance, that is, that the transformation equations could be "integrated," was rather rigid. The findings presented here strengthen that conclusion. The fact that we have now shown that a rather natural and attractive alternative gauging process is ultimately inconsistent makes the pursuit of extracting physical information from the original, consistent, approach that much more compelling.

C. Anomalies in Quantum Electrodynamics

In the last few years, we have significantly clarified the issue of species doubling and chiral symmetry breaking in the finite-element scheme [108,111]. In particular we have presented new calculations of the current anomalies in (1+1) dimensions for rectangular lattices, that is, when $h \neq \Delta$. In general, both vector and axial-vector currents are anomalous, but when $h = \Delta$ we recover the familiar finite-element result

$$\langle \partial_\mu j_5^\mu \rangle = \frac{e^2}{M \sin \pi/M} \begin{Bmatrix} \cos^2 \pi/2M \\ 1 \end{Bmatrix} E, \quad \text{for } M = \begin{cases} \text{odd} \\ \text{even} \end{cases}, \quad (8)$$

where the quotation marks signify the finite-element divergence, and E is the corresponding electric field. This then raises the question of how chiral symmetry is broken, since it would appear that the Dirac equation respects that symmetry. The answer lies in the fact that there is no action from which the lattice equations of motion can be derived. It is only possible to define an action in Euclidean space, where the equations refer equally to the past and the future, but not in Minkowski space, where the Dirac equation involves fields in the entire *prior* lattice, but not fields in the future. From an action, one can indeed derive a

symmetry current, which possesses no anomaly. That current cannot be used in Minkowski space, because again it involves fields defined on the entire lattice, making it impossible to solve the lattice difference equations by time-stepping through the lattice. If, in fact, one uses this unacceptable current to compute the vacuum polarization ($\langle A_\mu A_\nu \rangle$), one sees the explicit appearance of species doublers and a vanishing anomaly.

We expect to continue exploring quantum electrodynamics on the finite-element lattice. Some preliminary work on the anomalous magnetic moment of the electron is very encouraging, and we hope to obtain some nonperturbative results for that quantity, at least for small lattices.

D. Lattice Hamiltonian and Application to QCD

In applying the finite-element equations to non-Abelian theories, techniques based on a suitable definition of a lattice Hamiltonian (which exists in the sense of a unitary time evolution operator) will prove most useful. To give the reader a hint of the power of that time-evolution operator, we note that it gives for the ground-state energy of an anharmonic oscillator defined by the continuum Hamiltonian $H = p^2/2 + \lambda q^4/4$ the following value, in a two-state approximation:

$$\omega_0 = 0.42054\lambda^{1/3}, \quad (9)$$

accurate to 0.06% [105], while a similar approximation for the q^8 oscillator is good to 0.2% [107]. Wavefunctions are equally impressive.

These results were based on a closed-form expression for the matrix elements of the unitary time evolution operator, which does not require solving the equations of motion. Although this expression has only been derived for one dimension (quantum mechanics), similar techniques should be possible in higher dimensions. We hope to derive such expressions in the coming year for Abelian and non-Abelian theories, and ultimately use them to study chiral symmetry breaking issues, glueball masses, etc. A brief status report of the current status of the finite-element gauge theory program was given in my contribution to the Lattice 96 conference [115]. Although this program has been in hiatus for a couple of years, I hope to vigorously revive it in the coming year in collaboration with a new graduate student, and I hope to obtain new results for non-Abelian theories before the new millenium dawns in 2001.

II. δ EXPANSION; VARIATIONAL PERTURBATION THEORY

While I was on sabbatical leave at Imperial College, London, I worked with Hugh Jones on a variation of the above scheme, the so-called linear δ expansion, in which he is expert.²⁻⁸ There, one interpolates between a soluble action S_0 and the true action S for the theory under consideration by writing

$$S_\delta = (1 - \delta)S_0 + \delta S. \quad (10)$$

The Green's functions may then be obtained as power series in δ , which parameter is set equal to one at the end of the calculation. In practice, only a finite number of terms in the δ series are evaluated, so Green's functions depend on some unphysical parameter [analogous to M in the nonlinear delta expansion—see the 1994 proposal]. The value of this parameter is typically determined by use of the “principle of minimum sensitivity.”⁹ Both formal issues and practical calculations using this variant of the δ expansion were

explored. We concentrated on studying $1/N$ corrections to the fermion mass in the Gross-Neveu model, a useful laboratory for examining issues of renormalization. Although we made considerable progress in clarifying fundamental issues, it is clear that much work remains to be done to understand renormalization in this framework.¹⁰

Subsequently, I have turned my attention to another, but closely related, direction. I am about to embark on a series of calculations in QCD using the “variational perturbation theory” of Igor Solovtsov¹¹. It is possible to split up the QCD Lagrangian in such a way that expansion can be made in terms of a new parameter that is always less than unity, at the cost of introducing a parameter that must be determined by some variational procedure. In particular, as a first step we propose to use this technique to study inclusive τ decay, based on improved renormalization group techniques, in part harking back to [18]. This is not unrelated to the work described in §5.

The application of this method to quantum chromodynamics has been started in the papers of Solovtsov *et al.*¹¹. In particular, we write the QCD action as

$$S(A, q, \phi) = S_2(A) + S_2(q) + S_2(\phi) + gS_3(A, q, \phi) + g^2S_4(A), \quad (11)$$

where $S_2(A)$, $S_2(q)$ and $S_2(\phi)$ are the free actions of the gluon, quark, and ghost fields, S_3 contains the trilinear vertices, and S_4 the four-point gluon vertex. By introducing an auxiliary field, and breaking up the action into “free” and “interaction” parts with parameters ξ and ζ , we easily obtain the following expression for a general Green’s function corresponding to some product of operators O :

$$G(O) = \int D_{\text{QCD}} O V(A, q, \phi) e^{iS_0}, \quad (12)$$

where S_0 is the free action of QCD and

$$V(A, q, \phi) = \sum_n \sum_{k=0}^n \frac{1}{(n-k)!} \left(-\frac{\partial}{\partial \kappa} \right)^{n-k} \frac{i^k}{k!} \frac{1}{[1 + \kappa(\zeta^{-1} - 1)]^{\nu/2}} \\ \times [g_3 S_3(a, q, \phi)]^k \exp\{ig_4^2 S_4(A)\}, \quad (13)$$

where

$$g_3 = \frac{g}{[1 + \kappa(\zeta^{-1} - 1)]^{3/2}}, \quad g_4 = \frac{g}{[1 + \kappa(\xi^{-1} - 1)]^{1/2}}. \quad (14)$$

Here κ is set equal to 1 at the end of the calculation. Gauge invariance requires the two parameters here be related by $\xi = \zeta^3$. Analysis of the structure of this variational perturbation series shows that it can be organized in powers of the new small parameter $a = 1 - \zeta$ if the standard coupling constant g is related to a by

$$\frac{g^2}{(4\pi)^2} = \frac{1}{C} \frac{a^2}{(1-a)^3}, \quad (15)$$

where C is a positive constant and $0 \leq a < 1$. It is C that must be determined by some variational condition, such as requiring appropriate asymptotic behavior of the β function.

I have arranged, through extensive University funding, for Igor Solovtsov and Olga Solovtsova to come for two extended visits here: they are at OU for the entire Spring 2000 semester. In the first phase of our collaboration, concluding with their six-month visit to Oklahoma in 1997–98, we perfected the APT approach described in Section 5 below; now, to include mass and renormalon contributions, we are planning to apply the VPT method as well, and we have now completed a paper on e^+e^- annihilation which will soon be published in *Eur. Phys. J* [137].

FIG. 1. A plot of the TE (scalar) and the TM Casimir stresses F for $-2 < D < 2$ on a spherical shell. For $D < 0$ ($D < 2$) the stress F^{TE} (F^{TM}) is complex and we have plotted $\text{Re } F$.

III. THE CASIMIR EFFECT

When a quantum field is confined to a finite volume, or subject to nontrivial boundary conditions, the energy-momentum tensor develops a nonzero vacuum expectation value, and an observable force or energy results. Although this phenomenon has been known for 50 years,¹² and has been recently directly confirmed,¹³ it remains poorly understood except in the simplest geometries, which is unfortunate since it is a fundamental quantum phenomenon with important applications from hadronic to cosmological scales, and plays essential roles in atomic and condensed matter systems. However, recently there has been serious progress. First, we have begun to understand the profound nature of the dimensional dependence of the Casimir effect. Carl Bender and I examined the Casimir stress of a scalar field satisfying Dirichlet boundary conditions on a spherical shell in D dimensions [106]. The numerical results are shown in Fig. 1. We find that the Casimir stress vanishes as $D \rightarrow \infty$ (largely because the surface area of a D -dimensional sphere tends to zero as $D \rightarrow \infty$), and also vanishes when D is a negative even integer. Remarkably, the stress has simple poles at positive even integer values of D . Does this mean that we cannot make sense of the Casimir effect in two space dimensions, an arena of tremendous interest in condensed matter physics [99]? In addition to these poles, branch point singularities occur for D even and ≤ 0 .

In 1996, I completed the complementary calculation for a vector field (TM modes) bounded by a sphere [116]. Similar results are found, as also shown in Fig. 1, although the scale, qualitatively, is shrunk by a factor of 2 for $D < 2$. Branch points now occur for integer $D \leq 2$.

A. The Identity of van der Waals Interactions and Casimir Energies

Although it has been known almost from the beginning of the subject that van der Waals, or Casimir-Polder forces, were related to Casimir forces, the precise connection was clouded by divergences in the calculations. The explicit demonstration of the identity of these two forces was given in the case of dilute parallel dielectric slabs, where the Lifshitz formula for

the Casimir energy may be easily seen to be equal, if dispersion be neglected, to the sum of pairwise long-range van der Waals energies¹⁴:

$$V = -\frac{23\alpha_1\alpha_2}{4\pi r^7}, \quad (16)$$

the Casimir-Polder retarded dispersion potential.¹⁵ (See also [33].) Here the connection between dielectric constant ϵ and polarizability α is $\epsilon = 1 + 4\pi N\alpha$, N being the number density of molecules. The simple geometry makes this calculation easy and unambiguous.

The corresponding calculation for a spherical geometry is fraught with more difficulty. The sum of van der Waals interactions (16) for a spherical ball (the case for a spherical bubble is identical) has been given in [119]. A sensible procedure for carrying out the calculation is dimensional continuation, which has been advocated in [106,116]. That is, we evaluate the integral

$$E_{\text{vdW}} = -\frac{23}{8\pi}\alpha^2 N^2 \int d^D r d^D r' (r^2 + r'^2 - 2rr' \cos \theta)^{-\gamma/2}, \quad (17)$$

by first regarding $D > \gamma$ so that the integral exists. The integral may be done exactly in terms of gamma functions, which when evaluated at $D = 3$, $\gamma = 7$ yields [119]

$$E_{\text{vdW}} = \frac{23}{1536\pi a}(\epsilon - 1)^2. \quad (18)$$

(Note that the expression (17) is formally negative or attractive, while the continued result (18) is positive or repulsive.)

Of course, the above expression (17) in three dimensions is divergent. These divergences are of two kinds: "volume" and "surface." The volume divergence is a self-energy effect that would be present if the medium filled all space, and makes no reference to the interface, and therefore is quite unobservable. If the divergences are regulated by inserting a point-splitting cutoff, and the divergent terms are simply omitted, the same finite result (18) is again obtained.

Now, we have shown [128] by an explicit, exact calculation that the same result (18) is obtained from a field theory calculation of the Casimir effect for this geometry. We consider a dilute dielectric ball, which was already considered by me in [38]. (That is, we consider $\mu = 1$ everywhere, and $\epsilon = 1$ outside of the ball.) The formula, which still admits of dispersion, becomes in that case

$$E_C \approx -\frac{1}{8\pi a} \sum_{l=1}^{\infty} (2l+1) \frac{1}{2} \int_{-\infty}^{\infty} dy e^{iy\delta} (\epsilon(y) - 1)^2 x \frac{d}{dx} F_l(x), \quad (19)$$

where

$$F_l(x) = x^2 \left(1 + \frac{l(l+1)}{x^2} \right) - \frac{1}{4} \left(\frac{d}{dx} e_l s_l \right)^2 - x^2 \left[2 \left(1 + \frac{l(l+1)}{x^2} \right) e_l s_l - \frac{1}{2} \frac{d^2}{dx^2} e_l s_l \right]^2, \quad (20)$$

and $s_l(x)$, $e_l(x)$ are Ricatti-Bessel functions. The exact numerical integration of this result is

$$E_C = (\epsilon - 1)^2 \frac{0.004767}{a}. \quad (21)$$

(An approximate calculation was given earlier by Ref.¹⁶.) This precisely agrees with the van der Waals result (18).[†]

I have also, in collaboration with Vladimir Nesterenko and his son, returned to the cylindrical geometry as well. Using a simple contour integral method of summing over eigenvalues, we have reconfirmed my old result with DeRaad, for a perfectly conducting cylindrical shell [51] (which was simultaneously reconfirmed, to greater accuracy, by Gosdzinsky and Romeo¹⁸), and found a new result for a dilute cylinder: For either a dielectric cylinder, or a for a dielectric-diamagnetic cylinder satisfying the condition that $\epsilon\mu$ be the same both on the inside and the outside of the cylinder (that is, the speed of light is uniform), we find the Casimir energy vanishes! [129] We do not understand this remarkable result. I am collaborating with August Romeo on a calculation which we hope will shed light on the origin of these zeroes.

B. Sonoluminescence and the Dynamical Casimir Effect

In related direction, I have been examining the possibility that the “dynamical Casimir effect” may be relevant to the puzzling phenomenon of sonoluminescence.¹⁹ This idea was explored, in a preliminary way, by Julian Schwinger in the last few months before his death in 1994.^{20–23} As noted above, the calculation of the Casimir energy for a spherical cavity in a dielectric medium, which significantly generalizes the calculation I presented years ago [38], indicates unambiguously that the magnitude of the Casimir energy is too small by many orders of magnitude to be relevant to the huge burst of photons seen in sonoluminescence. This subject has become somewhat controversial, in view of the notoriety given to the work of Eberlein²⁴. Her conclusions are in error, as I have shown in [113] and [119].

Three years ago, the question of the relevance of the “volume contribution” to the Casimir energy was raised by Carlson *et al.*²⁵. I have argued that this is a spurious effect, merely being a renormalization of the energy density of matter [119]. Thus, the contributions of vacuum fluctuations to the sonoluminescence regime seems much too small to be significant.

In light of recent developments, the issue seems even more clear-cut. Now that we clearly see that the Casimir energy may be identified with van der Waals interactions, it seems perfectly plain that the volume effect that Carlson *et al.*²⁵ consider, proportional to $\epsilon - 1$, simply cannot be present, because such cannot arise from pairwise interactions. (This point was already made in [38].) Our interpretation stands vindicated: An effect proportional to the volume represents a contribution to the mass density of the material, and cannot give rise to observable effects.

More subtle is the role of surface divergences [113]. The zeta function regularization calculation we presented above simply discards such terms; but they appear in more physical regularization schemes. For example, if the divergences are regulated by introducing a time splitting parameter δ , we get from the leading asymptotic expansion

[†]As our paper was being written, we received a manuscript from G. Barton¹⁷, in which exactly the same identity between the van der Waals and Casimir forces was noted. He uses an elementary method of summing zero-point energies directly in powers of $(\epsilon - 1)$, using ordinary perturbation theory. Of course, our approach is in principle more general, in that it allows for arbitrary ϵ and permits inclusion of dispersion.

$$E_{\text{div}} = -\frac{(\epsilon - 1)^2}{4a} \frac{1}{\delta^3}; \quad (22)$$

and if a simple model for dispersion is used, with characteristic frequency ω_0 , the same result is obtained with $1/\delta \rightarrow \omega_0 a/4$ [113]. (A very similar result is given in Ref.¹⁶.) We believe these terms are probably also unobservable, for they modify the surface tension of the liquid, which, like the bulk energy, is already phenomenologically described. (That surface tension has its origin in the Casimir effect was proposed in [32].) In any case, this surface energy is probably too small, and definitely of the wrong sign, to be relevant to sonoluminescence. (The flash of light is emitted at the minimum radius.)

We note that Barton in his recent work¹⁷ seems to concur with our assessment: The terms "proportional to V [the volume] and to S [the surface] would be combined with other contributions to the bulk and to the surface energies of the material, and play no further role if one uses the measured values."

It is truly remarkable that however the (true) divergences in the theory are regulated, and subsequently discarded, the finite result is unchanged. That is, in the van der Waals energy, we can simply omit the point-split divergences, or proceed through dimensional continuation, where no divergences are explicit; in either case, the same result (18) is obtained. Likewise, the same result is obtained for the Casimir energy using either a temporal point-splitting, or an exponential wavenumber cutoff,¹⁷ and omitting the divergent terms; or through the formal trick of zeta-function regularization. It is worth re-emphasizing that we are not claiming that the Casimir effect for a dielectric ball is finite, unlike the classic case of a spherical conducting shell^{26,27} [33]. It is merely that those divergent terms serve to renormalize phenomenological parameters in the condensed matter system.

So, finally, we are left with the finite term, which in the dilute approximation is given by (18) or (21). For a bubble of minimum radius $\sim 10^{-4}$ cm, the corresponding Casimir energy is only $E_C \sim 10^{-3}$ eV. This is ten orders of magnitude too small to be relevant to sonoluminescence, where about one million optical photons are emitted per flash, and again the sign is wrong. (As to the relevance of a static calculation to the dynamical regime of sonoluminescence, we note that the adiabatic approximation seems favorable, since the time scale for the flash, $\sim 10^{-11}$ s is far longer than the time scale for optical photons, $\sim 10^{-15}$ s.)

Dynamically, there are new results as well. It seems plausible that the dynamical Casimir effect is closely allied with the so-called Unruh effect,²⁸ wherein an accelerated observer, with acceleration a , sees a bath of photons with temperature T ,

$$T = \frac{a}{2\pi}. \quad (23)$$

Indeed, the observed radiation in sonoluminescence is consistent with the tail of a blackbody spectrum, with temperature 20,000 K. That is, kT is about 1 eV. Let us, rather naively, apply this to the collapsing bubble, of radius R , where $a = d^2R/dt^2 \sim R/\tau^2$, where τ is some relevant time scale for the flash. We then have

$$kT \sim \frac{R}{(c\tau)^2} \hbar c, \quad (24)$$

or

$$1 \text{ eV} \sim \frac{10^{-3} \text{ cm} \cdot 2 \times 10^{-5} \text{ eV} \cdot \text{cm}}{\tau^2 (3 \times 10^{10} \text{ cm s}^{-1})^2} \sim \frac{10^{-29} \text{ eV}}{\tau^2 (\text{s}^2)}. \quad (25)$$

That is, $\tau \sim 10^{-15}$ s, which seems implausibly short; it implies a characteristic velocity $R/\tau \sim 10^{12}$ cm/s $\gg c$. It is far shorter than the upper limit to the flash duration, 10^{-11} s. Indeed, if we use the latter in the Unruh formula (23) we get a temperature below 1 milli

FIG. 2. The sudden collapse of an otherwise static bubble.

Kelvin! This conclusion seems consistent with those of Eberlein,²⁴ who indeed stressed the connection with the Unruh effect, but whose numbers required superluminal velocities.

However, we must remain open to the possibility that discontinuities, as in a shock, could allow changes on such short time scales without requiring superluminal speeds. Indeed, Liberati *et al.*,²⁹ following Schwinger's earlier suggestion,^{20,22} indeed assume an extremely short time scale, so that rather than the adiabatic approximation discussed above being valid, a sudden approximation is more appropriate. We therefore turn to an analysis of that situation.

The picture offered by Liberati *et al.*²⁹ is that of the abrupt disappearance of the bubble at $t = 0$, as shown in Fig. 2. On the face of it, this picture seems preposterous—the bubble simply disappears and water is created out of nothing. It is no surprise that a large energy release would occur in such a case. Further, the static Casimir effect calculations employed in Ref.²⁹ are invalid in this instantaneously changing model. Therefore, rather than computing Bogoliubov coefficients from the overlap of states belonging to two static configurations, let us follow the original methodology of Schwinger,^{20,22} which is essentially equivalent.

As in Schwinger's papers, let us confine our attention to the electric (TM) modes. They are governed by the time-dependent Green's function satisfying

$$(\partial_0 \epsilon(x) \partial_0 - \nabla^2) G(x, x') = \delta(x - x'). \quad (26)$$

The photon production is given by the effective two-photon source

$$\delta(JJ) = i\delta G^{-1} = i\partial_0 \delta \epsilon(x) \partial_0. \quad (27)$$

The effectiveness for producing a photon in the momentum element centered about $\mathbf{k}$ is

$$J_{\mathbf{k}} = \sqrt{\frac{(d\mathbf{k})}{(2\pi)^3} \frac{1}{2\omega}} \int (dx) e^{-i(\mathbf{k}\cdot\mathbf{r} - i\omega t)} J(x), \quad \omega = |\mathbf{k}|. \quad (28)$$

Let us follow Schwinger and consider one complete cycle of appearance and disappearance of the bubble, which we assume disappears for a time τ . The result of an approximate calculation coincides with that found by Schwinger.

$$|J_{\mathbf{k}} J_{\mathbf{k}'}|^2 \propto (\epsilon' - 1)^2 \omega^2 \sin^2 \omega \tau \delta(\mathbf{k} + \mathbf{k}') v. \quad (29)$$

The total number of photon pairs emitted is then, if dispersion is ignored,

$$N = \left(\frac{4\pi}{3}\right)^2 \left(\frac{R}{\Lambda}\right)^3 \left(\frac{\epsilon - 1}{4}\right)^2, \quad (30)$$

where the cutoff wavelength is given by $K = 2\pi/\Lambda$. Such a divergent result should be regarded as suspect. It was Eberlein's laudable goal²⁴ to put this type of argument on a

sounder footing. Nevertheless, if we put in plausible numbers, $\sqrt{\epsilon} = 4/3$, $R = 4 \times 10^{-3}$ cm, and, as in Schwinger's earlier estimate, $\Lambda = 3 \times 10^{-5}$ cm, we obtain the required $N \sim 10^6$ photons per flash.

The problem with this estimate is one of time and length scales—for the instantaneous approximation to be valid, the flash time τ_f must be much less than the period of optical photons, $\tau_o \sim 10^{-15}$ s. This is consistent with the discussion of the Unruh effect above, and is acknowledged by Liberati et al.²⁹ On the other hand, the collapse time $\tau_c \sim 10^{-5}$ s is vastly longer than τ_f , and is therefore totally irrelevant to the photon production mechanism. The flash occurs near minimum radius, and thus the appropriate value of R in Eq. (30) would seem to be at least an order of magnitude smaller, $R \sim 10^{-4}$ cm. This would lead to $N < 10^3$ photon pairs, totally insufficient.

IV. MAGNETIC MONOPOLES

Nearly five years ago, George Kalbfeisch, Mike Strauss, and I put together an experimental proposal to look for magnetic monopoles at Fermilab, which was approved in 1996. A new apparatus has been constructed, and we have just finished phase 1 of running samples from the dismantled CDF and D0 detectors at Fermilab. (See Task MM.)

A. Binding and Production

As part of this process I undertook to look at the theory of binding of monopoles to nuclei. Although there is a considerable literature, the theory is in rather unsatisfactory shape, because the $1/r^2$ potential binding the spin of the nucleus to the monopole is so singular at the origin. The binding is very sensitive to the spin of the nucleus: For example, simple angular momentum considerations show that if Be were spin 1/2 it would bind to nuclei, whereas, because of its actual spin of 3/2, it does not bind. (This is a relevant example, because the Be beam tubes are the obvious first place to look for monopoles and have been among our first samples.) Most of the calculations done so far are nonrelativistic, with hard core cutoffs,^{30–33} although Olaussen and Sollie use form factors.³⁴ There are only two relativistic calculations, for spin 1/2^{35–37} and spin 1.³⁸ (In the latter case, additional ad hoc interactions must be inserted.) The results are extremely sensitive to the form of the interaction used: see Table I. We propose to re-examine the issue of binding, particularly relativistically for general spin, for the most promising candidate experimentally now seems to be Al, which has spin 5/2. Experimentally, we should note that only a small amount of binding is necessary to trap the monopoles permanently in matter: only 1 eV binding will give a 10 yr lifetime; anything more than a kilovolt of binding will trap the monopole forever. (It is probable that even Be, which would be disrupted by the strong magnetic field of the monopole, will provide sufficient binding.) And once bound to an atom, the lattice will keep them in the bulk indefinitely. (Some details of these considerations appear in [134].)

There are also vital questions of production of monopoles. The previous estimates based on Drell-Yan mechanisms are suspect, even perturbatively, for they do not include the most obviously required velocity dependence in the coupling. Field theory calculations are difficult, because of the occurrence of the “string” in the coupling between electric and magnetic charge. (In fact, no practical, consistent field theory of magnetic charge has yet been given.) Previous estimates have been not only extremely crude, but often explicitly in violation of unitarity bounds. It is clear that the territory is ripe for civilizing. Leonard Gamberg and I have been intensively examining these issues for the past two years. We have just submitted a detailed reformulation of the quantum field theory of magnetic monopoles, which, for example, allows a consistent string-independent formulation in the

Nucleus	Spin	γ	$\hat{\gamma}$	J	E_b	Notes	Ref
n	$\frac{1}{2}$	-1.91		$\frac{1}{2}$	350 keV	NR,hc	32
^1_1H	$\frac{1}{2}$	2.79	2.79	$q - \frac{1}{2} = 0$	15.1 keV	NR,hc	31
					320 keV	NR,hc	32
					50-1000 keV	NR,FF	34
					263 keV	R	37
^2_1H	1	0.857	1.71	$q - 1 = 0$ ($n = 2$)	$\frac{130}{\lambda}$ keV	R,IM	38
^3_2He	$\frac{1}{2}$	-2.13	-3.20	$q + \frac{1}{2} = \frac{3}{2}$	13.4 keV	NR,hc	31
$^{27}_{13}\text{Al}$	$\frac{5}{2}$	3.63	7.56	$q - \frac{5}{2} = 4$	2.6 MeV	NR,FF	34
					560 keV	NR,hc	33
$^{113}_{48}\text{Cd}$	$\frac{1}{2}$	-0.62	-1.46	$q + \frac{1}{2} = \frac{49}{2}$	6.3 keV	NR,hc	31

TABLE I. Weakly bound states of nuclei to a magnetic monopole. The angular momentum quantum number of the lowest bound state is indicated. In Notes, NR means nonrelativistic and R relativistic calculations; hc indicates an additional hard core interaction is assumed, while FF signifies use of a form factor. IM = induced magnetization, the additional interaction employed for the relativistic spin-1 calculation. Here γ is the magnetic moment expressed in proton units, $\hat{\gamma}$ the magnetic moment expressed in nucleus units. J is the angular momentum of the lowest state in which binding occurs, $|eg| = n/2$, where g is the monopole strength, and $q = nZ/2$. We use $n = 1$ except for the deuteron, where $n = 2$ is required for binding.

eikonal approximation [136]. This paper was accepted enthusiastically, almost immediately, and will appear shortly in Phys. Rev. D.

We have used our emerging insights to critique the recent indirect mass limits claimed by the D0 experiment [125,134]. This work is described in more detail in our 1999 DOE proposal.

B. Theoretical aspects of the experimental situation

There are many theoretical issues that come up in the design and operation of the monopole detector. One of these is the modification of the flux by the superconducting cylinder. I have carried out new calculations of the flux seen by the superconducting loop by a monopole or pseudopole. The principal effect, valid for an infinitely long cylinder, is a reduction of the flux produced when a monopole of strength g passes through the loop by

FIG. 3. Schematic diagram of monopole detector

a geometrical factor depending on the ratio of the radius of the cylinder R to the radius of the superconducting detection loop a : The net current induced in the loop of inductance L is

$$I = \frac{4\pi g}{Lc} \left(1 - \frac{a^2}{R^2}\right). \quad (31)$$

As we have begun the actual running, we must face up to the (remarkably large) signals that arise from small magnetization of samples presumed nonmagnetic, and induced dipoles in conducting samples produced by traversing inhomogeneous fields. I have examined these questions as well. A magnetic dipole of strength μ along the axis of the detector produces a current signal at time t of

$$I(t) = -\frac{2\pi\mu}{Lc} \frac{a^2}{(a^2 + z(t)^2)^{3/2}}, \quad (32)$$

where $z(t)$ is the distance of the dipole from the loop. Inhomogeneous fields can induce a current in the sample, which in turn produces a magnetic dipole field which produces a signal in the detector. For a small sample of volume V_s with axial symmetry, an approximate formula for the signal in the detector is

$$I(t) = \frac{v}{c} \frac{\sigma(\pi a^2)V_s}{2Lc^2} \frac{\langle r^2 \rangle}{[a^2 + z(t)^2]^{3/2}} \frac{dB_z(z(t))}{dz}, \quad (33)$$

where v is the speed of the sample, $\sqrt{\langle r^2 \rangle}$ is the RMS radius of the sample, σ is the conductivity of the sample, and $B_z(z(t))$ is the magnetic field seen by the sample loop when it is at position $z(t)$ relative to the detection loop.

We have clearly been able to disentangle these effects in the experimental apparatus, and have demonstrated that a signal corresponding to a single Dirac pole strength can be readily detected. As noted elsewhere, we have finished phase 1 of our running, and will have preliminary results on cross section and mass limits by the end of the summer.

V. ANALYTIC PERTURBATION THEORY AND QCD

A. The Analytic Running Coupling

Igor Solovtsov, Olga Solovtsova, and I have been rapidly developing a new scheme for combining the renormalization group with perturbation theory. Eight papers have now appeared [117,118,121,122,124,126,127,132], and several more are in preparation. The conventional approach is inconsistent, as manifested by the Landau singularity or the ghost pole, which was first discovered in QED.³⁹ There, that unphysical spacelike singularity is merely a theoretical artifact with no phenomenological consequence. In QCD, in contrast, the effect is large and observable. We insist that a spacelike singularity (in higher orders it is not merely a pole, but a cut as well) cannot be present, because it violates general analytic requirements based on causality. It is easy to find a definition of the running coupling that respects the required analytic behavior, while incorporating low-order perturbative information.⁴⁰ On general grounds, the running couplings, defined in the spacelike (α) and timelike (α_s) regions, must satisfy the dispersion relations (here $a = \alpha/\pi$, $\alpha = g^2/4\pi$ is the dimensionless strong coupling constant, and the overbar signifies the running coupling)

$$\bar{a}_s(s) = -\frac{1}{2\pi i} \int_{s-i\epsilon}^{s+i\epsilon} \frac{dz}{z} \bar{a}(z), \quad (34)$$

$$\bar{a}(z) = -z \int_0^\infty \frac{ds}{(s-z)^2} \bar{a}_s(s), \quad (35)$$

which are valid if the couplings satisfy the following spectral representations: First, the timelike coupling can be written as

$$\bar{a}_s^{\text{APT}}(s) = \frac{1}{\pi} \int_s^\infty \frac{d\sigma}{\sigma} \varrho(\sigma), \quad (36)$$

where $\varrho(\sigma) = \text{Im} \bar{a}^{\text{PT}}(-\sigma - i\epsilon)$. Note that

$$-s \frac{d}{ds} \bar{a}_s^{\text{APT}} = \frac{1}{\pi} \varrho, \quad (37)$$

which is a relation advocated many years ago by Schwinger.⁴¹ That is, the β function is proportional to the spectral density. The spectral density ϱ plays a central role in the APT method; the spacelike running coupling constant, $\bar{a}^{\text{APT}}$, is also expressed through ϱ as follows:

$$\bar{a}^{\text{APT}}(z) = \frac{1}{\pi} \int_0^\infty \frac{d\sigma}{\sigma - z - i\epsilon} \varrho(\sigma). \quad (38)$$

The running coupling constant has a universal infrared value,

$$\bar{\alpha}_s^{\text{APT}}(s=0) = \bar{\alpha}^{\text{APT}}(Q^2=0) = \frac{4\pi}{\beta_0}, \quad (39)$$

in terms of the first β function coefficient, and is stable under change of renormalization scheme. The resulting running couplings differ significantly from the perturbative values, as Fig. 4 shows. These considerations make a major impact on the value of the running coupling extracted from QCD analysis of τ decay, and particularly on the value of Λ extracted from experiment, which may be 50% larger than in the conventional, internally inconsistent, analysis.

FIG. 4. Behavior of the 2-loop running coupling constant in the PT and APT schemes, both in terms of spacelike and timelike definitions. This figure is taken from [121].

B. Inclusive τ Decay

The inclusive semileptonic τ decay ratio for massless quarks is given⁴² in terms of the electroweak factor S_{EW} , the CKM matrix elements V_{ud} , V_{us} , and the QCD correction Δ_τ ,

$$R_\tau = 3S_{EW}(|V_{ud}|^2 + |V_{us}|^2)(1 + \Delta_\tau). \quad (40)$$

The correction Δ_τ may be written in terms of the functions $r(s)$ or $d(q^2)$, which are the QCD corrections to the imaginary part of the hadronic correlator Π : $\text{Im } \Pi \sim 1 + r$ and to the Adler D -function: $D = -q^2 d\Pi/dq^2 \sim 1 + d$. The analytic properties of the Adler function allow us to write down the relations analogous to those satisfied by the running couplings (34) and (35)

$$d(q^2) = -q^2 \int_0^\infty \frac{ds}{(s - q^2)^2} r(s), \quad (41)$$

$$r(s) = -\frac{1}{2\pi i} \int_{s-i\epsilon}^{s+i\epsilon} \frac{dz}{z} d(z). \quad (42)$$

Because of the proper analytic properties (which are violated in the usual perturbative approach) the QCD contribution to the τ ratio is given by the two equivalent forms

$$\begin{aligned} \Delta_\tau &= 2 \int_0^{M_\tau^2} \frac{ds}{M_\tau^2} \left(1 - \frac{s}{M_\tau^2}\right)^2 \left(1 + 2\frac{s}{M_\tau^2}\right) r(s) \\ &= \frac{1}{2\pi i} \oint_{|z|=M_\tau^2} \frac{dz}{z} \left(1 - \frac{z}{M_\tau^2}\right)^3 \left(1 + \frac{z}{M_\tau^2}\right) d(z). \end{aligned} \quad (43)$$

If one likes, one can think of r and d as effective running couplings in the timelike and spacelike regions, so as with the running couplings they may be expressed in terms of an effective spectral density,

$$d^{\text{APT}}(q^2) = \frac{1}{\pi} \int_0^\infty \frac{d\sigma}{\sigma - q^2} \rho^{\text{eff}}(\sigma), \quad (44)$$

$$r^{\text{APT}}(s) = \frac{1}{\pi} \int_s^\infty \frac{d\sigma}{\sigma} \rho^{\text{eff}}(\sigma), \quad (45)$$

possessing the same universal infrared limit as the running coupling.

In the ordinary perturbative approach, the function $d(q^2)$ may be expanded in terms of the running coupling and in the third order is

$$d^{\text{PT}}(q^2) = a^{\text{PT}}(-q^2) + d_1 [a^{\text{PT}}(-q^2)]^2 + d_2 [a^{\text{PT}}(-q^2)]^3, \quad (46)$$

where numerically, for three active flavors, the coefficients are $d_1^{\overline{\text{MS}}} = 1.6398$, $d_2^{\overline{\text{MS}}} = 6.3710$. Such is not the case in APT; rather, d^{APT} is constructed from Eq. (44) with a spectral density obtained as the imaginary part of d^{PT} on the physical cut:

$$\rho^{\text{eff}}(\sigma) = \rho_0(\sigma) + d_1 \rho_1(\sigma) + d_2 \rho_2(\sigma), \quad (47)$$

where $\rho_n(\sigma) = \text{Im } a^{n+1}(-\sigma - i\epsilon)$. We use the TAU'98 conference value⁴³ value[†] $R_\tau = 3.642 \pm 0.019$. Our results are shown in Table II, where for the sake of illustration we also show the PT results obtained by using the contour integral representation given in the second line of Eq. (43).

Most remarkably, APT exhibits very little renormalization scheme (RS) dependence. The τ decay coefficients d_1 and d_2 are RS dependent, as are all but the first two beta function coefficients, defined by the renormalization group equation

$$\mu^2 \frac{da}{d\mu^2} = -\frac{b}{2} a^2 (1 + c_1 a + c_2 a^2). \quad (48)$$

Because of the existence of RS invariants, one can investigate the sensitivity of the predicted value of the QCD correction to the choice of RS by varying d_1 and c_2 in a region where the degree of cancellation in the second RS invariant $\omega_2 = c_2 + d_2 - c_1 d_1 - d_1^2$ does not exceed a specified limit, taking, for example,

$$\frac{|c_2| + |d_2| + c_1 |d_1| + d_1^2}{|\omega_2|} \leq 2. \quad (49)$$

That sensitivity is shown [122] in Fig. 5, based on $\Delta_\tau^{\overline{\text{MS}}} = 0.1906$. Observe that the relative difference between the prediction of the lower corners of the domain defined by Eq. (49) is 0.8%, while in PT that difference is 6%. Note that the $\overline{\text{MS}}$ scheme lies outside this domain, as does the so-called V scheme, the prediction of which differs by only 0.2% from the $\overline{\text{MS}}$ value in APT, but by 61% in PT! To all intents and purposes, APT exhibits practically no RS dependence. (Because of the perturbative stability, the known three-loop level is quite adequate for these RS recalculations.)

[†]This is consistent with the 1998 PDG value,⁴⁴ $R_\tau = 3.642 \pm 0.024$.

FIG. 5. Contour plot of values of $\Delta\tau^{\text{APT}}$ at the three-loop order as a function of RS parameters d_1 and c_2 . The dashed line indicates the boundary of the domain, defined by Eq. (49), and the heavy points are the positions of the A , B and $\overline{\text{MS}}$ schemes.

TABLE II. QCD parameters extracted from $R_\tau = 3.642 \pm 0.019^{43}$ in the $\overline{\text{MS}}$ scheme. This table is taken from the revised version of [122]. The two perturbative schemes Br and LP are due to Braatan⁴⁵ and to Le Diberder and Pich.⁴⁶

Approximation	Method	$\Lambda_{\overline{\text{MS}}} \text{ (MeV)}$	$\alpha(M_\tau^2)$
NNLO	PT (Br)	366 ± 14	0.328 ± 0.007
	PT (LP)	391 ± 16	0.340 ± 0.008
	APT	907 ± 94	0.403 ± 0.015
NLO	PT (Br)	492 ± 17	0.371 ± 0.009
	PT (LP)	465 ± 19	0.358 ± 0.009
	APT	954 ± 90	0.404 ± 0.014

(c) *Deep-Inelastic Scattering Sum Rules*

At present, the polarized Bjorken and Gross-Llewellyn Smith deep inelastic scattering sum rules allow the possibility,⁴⁷ as with τ decay, of extracting the value of α_S from experimental data at low Q , here down to 1 GeV.

The Bjorken sum rule refers to the value of the integral of the difference between the polarized structure functions of the neutron and proton,

$$\Gamma_1^{p-n} = \int_0^1 dx [g_1^p(x, Q^2) - g_1^n(x, Q^2)] \quad (50)$$

$$= \frac{1}{6} \left| \frac{g_A}{g_V} \right| [1 - \Delta_{\text{Bj}}(Q^2)] , \quad (51)$$

where the prefactor in the second line is the parton-level description. In the conventional approach, with massless quarks, the QCD correction is given by a power series similar to Eq. (46), with coefficients for three active flavors being $d_1^{\overline{\text{MS}}} = 3.5833$, $d_2^{\overline{\text{MS}}} = 20.2153$. However, this description violates required analytic properties of the structure function moments, which, as has been argued,⁴⁸ follow from the existence of the Deser-Gilbert-Sudarshan integral representation.

Thus we adopt the analytic approach, which says instead

$$\Delta_{\text{Bj}}^{\text{APT}}(Q^2) = \delta^{(1)}(Q^2) + d_1 \delta^{(2)}(Q^2) + d_2 \delta^{(3)}(Q^2) , \quad (52)$$

where

$$\delta^{(k)}(Q^2) = \frac{1}{\pi^{k+1}} \int_0^\infty \frac{d\sigma}{\sigma + Q^2} \text{Im} [\alpha_{\text{PT}}^k(-\sigma - i\epsilon)] , \quad (53)$$

that is, Δ^{APT} is not a power series in α^{APT} .

FIG. 6. Γ_1^{p-n} with 1-, 2-, and 3-loop QCD corrections vs. the coupling constant.

Besides possessing the correct analyticity, the APT approach has two key properties in its favor. First, successive perturbative corrections are small, so that the two- and three-loop QCD results are nearly the same. This is not the case in the conventional PT approach, as is shown [124] in Fig. 6. Second, the renormalization scheme dependence is again very small, so that various schemes which have the same degree of cancellation as the $\overline{\text{MS}}$ scheme give the same predictions all the way down to $Q^2 \sim 1 \text{ GeV}^2$. In contrast, it is impossible to make any reliable prediction for conventional PT, even if improved by the Padé approximant (PA) method,⁴⁹ for Q^2 below several GeV^2 . This is illustrated [124] in Fig. 7. One can see that instead of RS unstable and rapidly changing PT functions, the APT predictions are slowly varying functions, which are practically RS independent.

A precisely similar analysis can be performed on the GLS sum rule, which refers to the integral

$$S_{\text{GLS}} = \frac{1}{2} \int_0^1 dx [F_3^{\bar{\nu}p}(x, Q^2) + F_3^{\nu p}(x, Q^2)] \quad (54)$$

$$= 3 [1 - \Delta_{\text{GLS}}(Q^2)]. \quad (55)$$

Again, the APT approach leads to perturbative stability, and to practically no renormalization scheme dependence down to very low Q^2 [127].

There are purely theoretical implications of this work as well. Central to the approach is the recognition that the definition of the running coupling is different in the timelike and spacelike regimes. Yet the two functions possess common behavior at infinity and at zero momentum, and disagree only in three loops, where the beta function possesses renormalization scheme dependence. It therefore might be thought that with a proper choice of scheme the two couplings could be brought into agreement. Solovtsov and I have presented a simple proof, however, that such an identity is impossible based on the required analyticity [132].

FIG. 7. Renormalization scheme dependence of predictions for Δ_{Bj} vs. Q^2 for the APT and PT expansions. The solid curves, which are very close to each other, correspond to the APT result in the $\overline{\text{MS}}$, A, B, and ECH schemes. The PT evolution in the $\overline{\text{MS}}$, A, and B schemes are shown by dashed curves, the ECH scheme is indicated by a dotted curve, and the PA results in the $\overline{\text{MS}}$, A, and B schemes are denoted by dash-dotted curves. (The definition of the various schemes is given in our paper [124].) The SMC data⁵⁰ is indicated by a square, the triangle is the E154 data,⁵¹ circles are E143 data,⁵² and the stars are more recent E143 data.⁵³

VI. PT SYMMETRIC THEORIES

A. Perturbative Parity Symmetry Breaking

In the summer of 1996, Carl Bender and I initiated a new variant on the δ expansion [76,77,80–89,94–97] which should allow us to investigate, perturbatively, issues of symmetry breaking [114]. For a scalar field, instead of expanding in powers of ϕ^2 , we consider, instead,

$$\mathcal{L} = (\nabla\phi)^2/2 + m^2\phi^2/2 - g(i\phi)^{2+\delta}, \quad (56)$$

which is a free theory if $\delta = 0$. We believe this theory, surprisingly, has only positive energies for $\delta > 0$. This theory, even in first order in δ , exhibits symmetry breaking in the sense that the vacuum expectation value of ϕ is nonzero. In fact, to first order in δ , we have established the following formula

$$\langle\phi\rangle = -i\delta M^2\mu^{-3}(\mu^2/4\pi)^{D/4}\sqrt{\pi\Gamma(1-D/2)/2}, \quad (57)$$

where M is an arbitrary variational mass, $\mu^2 = m^2 + M^2$, and D is the number of space-time dimensions. Note that $\langle\phi\rangle$ is real for $2 < D < 4$. This result is remarkable, because conventional perturbation theory is unable to give any information about symmetry breaking.

B. Supersymmetry

Bender and I wrote a second paper applying this new variant of the delta expansion to a supersymmetric model [120]. We develop rather elaborate Feynman rules to enable the calculations to be carried out to second order in δ . In particular, we consider a two-dimensional model of fermions and bosons interacting through a superpotential of the form

$$\mathcal{S}(\phi) = -ig(i\phi)^{1+\delta}. \quad (58)$$

The resulting quantum field theory exhibits broken parity symmetry for all $\delta > 0$, because $\langle\phi\rangle \neq 0$. However, supersymmetry remains unbroken, because the ground-state energy vanishes, and the fermion-boson mass ratio is unity, due to detailed cancellation between diagrams involving bosons and fermions. One might think that an anomaly could develop, because this cancellation is formal, involving the subtraction of infinite integrals. However, we can find a regularization scheme, based on Schwinger's proper-time method, which respects the supersymmetry, and hence exactly justifies the cancellation.

C. PT Symmetric Quantum Electrodynamics

The results⁵⁴ regarding PT-symmetric non-Hermitian quantum field theories are intriguing because they suggest that it is possible to formulate a new kind of electrodynamics. Instead of coupling the A field to a vector current, why not couple this field to an axial-vector current? Of course, this coupling breaks parity symmetry. Therefore, we also replace e by ie , thereby breaking time-reversal invariance as well! The resulting PT-symmetric, massless Lagrangian is

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} - \frac{1}{2}\psi\gamma^0\gamma^\mu\frac{1}{i}\partial_\mu\psi + e\frac{1}{2}\psi\gamma^0\gamma^5\gamma^\mu A_\mu\psi. \quad (59)$$

We conjecture on the basis of our experience with scalar theories that the spectrum of this theory is physically acceptable in that it is bounded below.

Our conventions in Eq. (59) are as follows: γ^0 is antisymmetric and pure imaginary, $\gamma^0\gamma^\mu$ is symmetric and real, $\gamma^5 = \gamma^0\gamma^1\gamma^2\gamma^3$ is antisymmetric and real, and $(\gamma^5)^2 = -1$. The fermion field ψ is expected to be complex, as is the operator ϕ in Eq. (56).

Like the conventional electrodynamics, $\mathcal{L}$ in Eq. (59) possesses gauge invariance; $\mathcal{L}$ is invariant under the replacements

$$A^\mu \rightarrow A^\mu + \partial^\mu \Lambda, \quad \psi \rightarrow e^{-ie\gamma^5 \Lambda} \psi. \quad (60)$$

Note that this gauge transformation on the fermion field is not a phase transformation when e is real; rather it changes the scale of ψ . However, the bilinear forms in the fermion field in the Lagrangian and in the energy-momentum tensor are all invariant.

This new theory allows us to re-examine the program of Johnson, Baker, and Willey⁵⁵. The objective of the program was to find a critical value of the coupling constant e which would make massless quantum electrodynamics finite. The mass shift in this theory is finite because the unrenormalized masses are zero. Thus, the only possible infinite quantities are associated with the three renormalization constants Z_1 , Z_2 , and Z_3 . However, to any order in powers of e , it is possible to find a gauge in which $Z_1 = Z_2$ is finite. Thus, the only remaining divergent quantity is $1/Z_3$. Demanding that this be finite translates into an eigenvalue condition on the fine structure constant

$$\alpha = \frac{e^2}{4\pi}. \quad (61)$$

This eigenvalue condition takes the form

$$F_1(\alpha) = 0. \quad (62)$$

The function $F_1(\alpha)$ has been calculated to three loops (four terms) in weak-coupling perturbation theory:

$$F_1(\alpha) = \frac{4}{3} \left(\frac{\alpha}{4\pi} \right) + 4 \left(\frac{\alpha}{4\pi} \right)^2 - 2 \left(\frac{\alpha}{4\pi} \right)^3 - 46 \left(\frac{\alpha}{4\pi} \right)^4 + \dots \quad (63)$$

The first two terms in this series were calculated by Jost and Luttinger⁵⁶. Unfortunately, with just two terms the only nontrivial solution of Eq. (62) is negative, which gives an unphysical imaginary value for e . In a dramatic development, Rosner⁵⁷ calculated the third term in the series. The negative sign of his result is significant because now there is a positive root to the cubic polynomial equation obtained by truncating $F_1(\alpha)$ after three terms:

$$\alpha = 13.872. \quad (64)$$

Rosner's two-loop result is surprising because it is rational. His work suggests the conjecture that all of the coefficients in the expansion of $F_1(\alpha)$ might be rational, possibly reflecting a deep symmetry of the underlying massless theory⁵⁸. This conjecture has recently gained support through the stunning calculation of the three-loop coefficient by Gorishny and coworkers⁵⁹⁻⁶¹. See also Broadhurst's recent papers.⁶² The fourth-degree equation gives one positive nontrivial root for α :

$$\alpha = 3.969. \quad (65)$$

This value differs from the result in Eq. (64) by a factor of 3.5, which suggests that this nontrivial root is unstable.

Apart from a possible overall factor of $e\gamma^5$ in some of the Green's functions, the Feynman rules for the Lagrangian (59) give precisely the same weak-coupling expansion as in the conventional massless quantum electrodynamics, except that α is now replaced by $-\alpha$. Thus, in this new and peculiar theory of quantum electrodynamics, the expansion of $F_1(\alpha)$ becomes

$$F_1(\alpha) = -\frac{4}{3} \left(\frac{\alpha}{4\pi} \right) + 4 \left(\frac{\alpha}{4\pi} \right)^2 + 2 \left(\frac{\alpha}{4\pi} \right)^3 - 46 \left(\frac{\alpha}{4\pi} \right)^4 + \dots \quad (66)$$

Now, we find a stable sequence of roots obtained by considering 2, 3, and 4 terms in the series, converted into an appropriate Padé approximant:

$$\alpha_2 = 4.189, \quad (67)$$

$$\alpha_3 = 3.590, \quad (68)$$

$$\alpha_4 = 4.110. \quad (69)$$

Note that the sequence of roots $\alpha_2, \alpha_3, \alpha_4$ is remarkably stable. It would be extremely interesting to calculate the roots of the $(3, 2), (4, 2), (4, 3), \dots$, Padés.

We also note that the Casimir model⁶³ for the electron now actually exists: That is, the Coulomb *attraction* can now be balanced by the Casimir repulsion^{26,27} [33] and give a positive value for the fine structure constant [123].

D. Solving the Schwinger-Dyson Equations

In a recent paper submitted to Physical Review, Bender, Savage, and I have set up the Schwinger-Dyson equations for such a PT -symmetric theory, solved those equations by truncation, and renormalized such theories [135]. We first examine zero-dimensional theories where everything can be done explicitly. Very accurate numerical results are found. Most interesting are the Green's functions for the massive theory, which exhibit distinctive nonperturbative behavior. One dimensional theories are also studied numerically. However, of course, the results in D dimensions are of the most interest, because renormalization is then required. There, we first considered the PT -symmetric $ig\phi^3$ theory in 6 space-time dimensions. We find a perturbative solution of the Schwinger-Dyson equations in which the n point function goes like $G_n \sim g^{n-2}$; in particular, $G_1 = im^2/g$. In this theory the beta function turns out to be positive, so unlike the usual ϕ_6^3 theory, this is not asymptotically free. Of much more interest is the ϕ_4^4 theory, because this presumably bears on the Higgs sector of the standard model. Here there are two regimes. If $m^2 > 0$ the odd Green's functions, although nonzero, are exponentially small, $G_{2n+1} \sim e^{-1/g}$, the even Green's functions decouple from the odd ones, and the sign of the beta function reverses so the theory is asymptotically free. Even more interesting is the case $m^2 < 0$. Then all the Green's functions are substantial, $G_n \sim g^{n/2-1}$, and we have perturbative parity-symmetry breaking, while the theory is probably still asymptotically free, and certainly renormalizable. Implication of these fascinating results for physics will be explored in future papers.

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Task C: CLEO

Introduction

One of the OU experimental group's major experiments is the CLEO experiment at the CESR storage ring at Cornell University. During the grant period we have had two postdoctoral RA's working on the CLEO experiment, one resident at Cornell and one resident at OU. We have also had three graduate students working on CLEO during this period.

CLEO II Analysis

A major effort of the group was maintenance of CLEO II software and data analysis. After moving to Cornell, GRA Steve Richichi took on important responsibilities to run jobs that reanalyze previous data with a new version of the Pass 2 program. This was a major effort and the resulting new compressed data set will be used in future physics analysis. This effort is now finished and all the existing CLEO data has been reprocessed.

The CLEO analysis efforts included study of inclusive B decays to final states with strangeness. Two different analyses were carried out: $B \rightarrow \Lambda X$, and $B \rightarrow K^\pm X$. Further work remains to be done on the $B \rightarrow K^\pm X$ decay mode, which is a systematics limited measurement. In the case of the $B \rightarrow \Lambda X$ analysis, the availability of vastly improved statistics and better secondary-vertex reconstruction allow a high-precision measurement with a significantly improved precision over previous experiments. This result, which was presented at the EPS97 conference in Jerusalem, Israel, is useful in understanding the overall picture of B decays.

During the period of this grant, GRA X. Fu completed his thesis on $\bar{B} \rightarrow D^* \bar{\ell} \bar{\nu}$ decay and obtained the best current CLEO value of $|V_{cb}|$ [1]. RA B. Nemati completed analysis on $\Lambda_c^+ \rightarrow \rho \phi$ which resulted in a CLEO paper [2]. GRA M. Wood completed analysis on inclusive $B \rightarrow D X$ decay and contributed to a CLEO publication [3]. RA H. Severini finished the analysis of inclusive B meson decays to the charmed strange baryons Ξ_c^0 and Ξ_c^+ , which is a project he started before joining our group. This analysis includes a detailed measurement of the branching fractions of these never before seen decay channels as well as a measurement of their momentum spectra. A CLEO publication resulted from this work [4].

Hardware Efforts

Our group contributed to construction of the silicon vertex detector for CLEO III. The silicon work for CLEO III consisted of two major parts:

- (1) Detector quality assurance tests of microstrip detectors, including measurements of detector capacitance's and leakage currents as a function of depletion voltage, were made. These measurements are carried out using equipment such as the HP 4194A phase-gain analyzer, probe station, and computer. We use a 64 channel probe card to characterize detectors before wire-bonding and assembly. We have developed the capability to automate the measurements using computer controlled hardware and LABVIEW software. This effort included development of a VLSI analogue switching circuit to scan each of the probe card channels, installation of stepping motors to computerize the probe station, and installation of a clean room in which to do the tests. The first production lot of 42 detectors were tested. The tests are described in a paper submitted (by P. Skubic) to the VERTEX '97 conference which was held in Rio de Janeiro, Brazil during September 1-5, 1997 [5].
- (2) We are developing flexible cables using High Density Interconnect (HDI) techniques to provide needed connections between silicon detectors and electronics hybrids. During this grant period, we have had 50 prototype flex circuits constructed by the General Electric Corporate Research Facility in Schenectady, NY. These flex circuits were fully tested and meet the CLEO specifications. GE was chosen as the vendor for production of the final flex circuits and funding was included in our CLEO III budget requests (FY96 and FY97) for their purchase.

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Report on Task D: the Neural Network Initiative

Theorists R. Kantowski and Caren Marzban applied Neural Networks to find primary and secondary vertices in simulated pixel detector data appropriate for SDC. Phil Gutierrez and Pat Skubic assisted in providing the simulated data by using SD-CSIM, a ISAJET GEANT based Monte Carlo program. The goal was to develop a fast/accurate trigger for the presence of secondary B and Charmed meson vertices. One publication resulted: R. Kantowski and C. Marzban, "A Neural Network for Locating the Primary Vertex in a Pixel Detector" *Nucl. Inst. and Meth. A* **355**, 582-588 (1995).

Task E: ATLAS

Three OUHEP faculty members (P. Gutierrez, P. Skubic, and M. Strauss) and two adjunct faculty members from Langston University (T. McMahon and J. Snow) have joined the ATLAS experiment, one of two general purpose experiments to be done at the Large Hadron Collider (LHC) which is now under construction at CERN. Our group is collaborating with other institutions including LBL, U. of New Mexico, and the State University at Albany on development and construction of the ATLAS pixel vertex detector. This device will occupy the region nearest to the beam line and therefore is one of the most important and technologically advanced subsystems in the ATLAS detector. The U.S. institutions have responsibility for providing the forward (disk) pixel modules. The contribution of OUHEP to the ATLAS pixel detector included the following projects:

A. Flex Hybrid for Disk Pixel Modules

We have taken on responsibility for design, production, and testing of the flex hybrid and related circuits which carry analog and digital signals to and from the disk module front-end (FE) readout chip and the module clock and control chip (MCC). We collaborated with Genoa on design of a flex hybrid for the demonstrator pixel module. Genoa designed the MCC chip and module support PC board, while we were responsible for the layout of the flex hybrid circuit. This work was done by an EE Masters student (S. Krishnama) and Senior Electronics Technician G. Boyd. This effort builds on our experience in developing the CLEO III flex circuits. We expect that prototype flex hybrids for the pixel demonstrator may be fabricated by the General Electric Corporate Research and Development center and by CERN so that two processes can be compared. The GE process, which was used for CLEO III, uses Cu on a Kapton substrate. Preliminary layouts have been done by EE MS student K. Neriyanuri who graduated. Simulations have been done of the preliminary design based GE design rules and on current designs of the front-end chips and MCC. Results of this work were presented at the PIXEL98 workshop at Fermilab in May 1998 [1].

B. Readout Electronics

The OU group has contributed to development of the pixel readout electronics by the efforts of OU Electrical Engineering student (Mr. Atul Joshi) who completed his Masters thesis during the Fall 1996 semester [2]. Another EE student (Ms. Alice Lu)

worked on the pixel readout and made significant contributions to the characterization of the Honeywell process which will be used to fabricate radiation hardened VLSI chips. This work included testing prototype chips and developing and evaluating simulation techniques to better understand how to optimize the design parameters of the process. Test structures provided by Honeywell were irradiated and characterized to allow better simulations of rad-hard designs. A summary of her contributions is included in her thesis [3].

C. Pixel Detector Development

We contributed to the development of pixel detectors for ATLAS in close collaboration with State University at Albany and the University of New Mexico. We evaluated a demonstrator device provided by pixel detector vendor CIS as part of the prototype procurement process. This is a common ATLAS order with the U.S. funds being provided by Albany. After the prototype detectors were received, we collaborated with UNM and Albany to test irradiated and non-irradiated wafers, pixel detectors, and test structures. Tests on detectors included leakage current and capacitance measurements versus bias voltage.

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