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**A PROPOSAL TO EXTEND
THE INTENSITY FRONTIER OF
NUCLEAR AND PARTICLE
PHYSICS TO 45 GeV
(LAMPF II)**

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LOS ALAMOS NATIONAL LABORATORY

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ACKNOWLEDGMENTS

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The following committees and individuals made important contributions to this document:

SCIENCE POLICY ADVISORY COMMITTEE

Barry C. Barish	California Institute of Technology
Val L. Fitch	Princeton University
Leonard S. Kisslinger	Carnegie-Mellon University
Alan D. Krisch	University of Michigan
Malcolm MacFarlane	Indiana University
Sydney Meshkov	National Bureau of Standards
Peter Parker	Yale University
Frederick Reines	University of California, Irvine
S. Peter Rosen	Los Alamos National Laboratory, P-DO
Lee C. Teng	Fermi National Accelerator Laboratory
Akihiko Yokosawa	Argonne National Laboratory
John D. Walecka	Stanford University

SCIENCE POLICY ADVISORY COMMITTEE - NUCLEAR SUBCOMMITTEE

Andrew Bacher	Indiana University
Donald Geesaman	Argonne National Laboratory
William Gibbs	Los Alamos National Laboratory
Charles Glashausser	Rutgers University
Jochen Heisenberg	University of New Hampshire
Leonard Kisslinger	Carnegie-Mellon University
Malcolm MacFarlane	Indiana University
Gerald Miller	University of Washington
Christopher Morris	Los Alamos National Laboratory

LAMPF II ACCELERATOR REVIEW COMMITTEE

Lee Teng, Chairman	Fermi National Accelerator Laboratory
Matthew Allen	Stanford Linear Accelerator Center
Mark Barton	Brookhaven National Laboratory
Michael Craddock	TRIUMF
Grahame Rees	Rutherford Laboratory
Robert Siemann	Cornell University
Lloyd Smith	Lawrence Berkeley Laboratory

OTHER PHYSICS COMMUNITY CONTRIBUTORS

P. D. Barnes	Carnegie-Mellon University
M. Chanowitz	Lawrence Berkeley Laboratory
J. Cronin	University of Chicago
R. Cutkosky	Carnegie-Mellon University
J. Deutsch	Université de Louvain, Belgium
Glennys Farrar	Rutgers University
A. S. Goldhaber	Stony Brook
V. Hughes	Yale University
T. Kalogeropoulos	University of Sycracuse
E. Lomon	Mass. Inst. of Technology
H. Lubatti	University of Washington
D. Marlow	Princeton University
Benjamin Nefkens	University of California, Los Angeles
W. Praeg	Argonne National Laboratory
L. Ray	University of Texas
R. Redwine	Mass. Inst. of Technology
F. Reines	University of California, Riverside
R. Shrock	Stony Brook
H. White	Brookhaven National Laboratory
M. Zeller	Yale University

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L. Agnew	T. Hiron
H. Baer	C. Hoffman
T. Bhatia	E. Hoffman
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D. Bowman	A. Jason
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G. Garvey	S. P. Rosen
W. Gibbs	V. Sandberg
B. Gibson	G. Sanders
T. Goldman	R. Silbar
D. Hagerman	G. Spalek
V. Hart	R. Stokes
A. Harvey	D. Strottman
W. Haxton	H. Thiessen
P. Herczeg	J. Warren
	M. Wilson

CHAPTER 1

EXECUTIVE SUMMARY

CHAPTER 1. EXECUTIVE SUMMARY

We propose to construct and operate LAMPF II, a high-intensity, medium-energy synchrotron addition to the Clinton P. Anderson Meson Physics Facility at the Los Alamos National Laboratory. This new facility will

- provide the nation's principal capability to address the frontiers of nuclear physics with an unmatched array of light hadronic probes;
- probe physical regimes beyond the standard model of the strong and electroweak interactions;
- provide unique opportunities to address the quark/gluon frontier in a manner complementary to other forefront facilities in nuclear and particle physics; and
- provide key facilities for basic research and education to a generation of scientists in the 1990's and well into the next century.

LAMPF II will consist of

- a 6-GeV, 170- μ A booster accelerator and
- a 45-GeV, 34- μ A, 3-Hz main synchrotron, with a 50% duty factor.

The injector accelerator will be the existing 800-MeV, 1-mA LAMPF linac. With its relatively high energy and compact emittance, the LAMPF linac is superior to any other existing potential injector. A broad array of beams will serve research programs using kaons, pions, protons, antiprotons, muons and neutrinos.

These beams will provide the capability to probe the underlying behavior of quarks and gluons in nuclear matter. This is the essential frontier of nuclear science today. It is the extension of the historical development of nuclear physics. The early description of nuclei using only bound nucleons evolved to include the second treatment of nuclear behavior with explicit meson properties. LAMPF has been a key tool in this endeavor. Nuclear science is embracing a third and even more fundamental constituent picture, emphasizing quarks and gluons. To explore the consequences of quark/gluon substructure in nuclei, the field theory of the strong interaction, quantum chromodynamics (QCD), is being applied to the nuclear many-body problem.

To be sure, all three of these perspectives are required to contribute to the principal endeavor of nuclear physics, a comprehensive understanding of the nucleus and its response to external probes. Indeed, even to attack only the newest frontier requires a complementary set of probes, consisting of high-energy electrons, relativistic heavy ions, and high-energy light hadrons.

Electrons are the best projectiles for the study of the response of the charged constituents of nuclei, described by the electromagnetic currents. In the newest picture the charged constituents are the quarks which have charges of

$-1/3$ and $+2/3$ of an electron charge. Electrons interact with nuclear matter primarily through the electromagnetic force. Thus, a very important class of nuclear processes is accessible with electrons, those processes excited by a charged, nonhadronic probe.

Nuclei, however, consist of both quarks and gluons. Typically, the quarks carry only about half of the momentum content of nuclear matter. The gluons, which bind the quarks through the color force carry the remainder. Gluons are not electrically charged. Thus, electron probes are unable to directly stimulate a significant class of nuclear responses. High energy collisions of heavy nuclei, produced by relativistic heavy ion accelerators, can create the high-density and high-temperature conditions in nuclear matter in which the many-body modes of quark and gluon behavior can be studied. Heavy ions provide tools to test the gross behavior of this hadronic matter under extraordinary conditions. Relativistic heavy ion collisions can create conditions in which a rich mixture of up and down valence quarks, virtual sea quarks, and color gluons are heated to high temperature or compressed to high density.

We propose to provide nuclear physics with a third and complementary tool, light hadrons. The principle advantage of light hadronic probes is selectivity. Like electrons, these probes are charged and therefore are sensitive to electromagnetic properties. Like heavy ions, these probes access both the quarks and gluons. By selecting the appropriate hadron beam, choosing from a menu of kaons, antiprotons, pions, and protons, the experimenter can choose the quark content of the beam and thereby select the nuclear characteristic to be investigated. Opportunities for frontier research with nuclei using a light hadron facility like LAMPF II will be unmatched in versatility and the breadth of capabilities offered.

While several of the selected programs of science we have identified may be addressed at existing facilities, there is no current prospect anywhere to advance this forefront of nuclear physics in a concerted and timely manner. Isolated programs at high-energy physics accelerators and scattered efforts at other laboratories fall short by a factor of a hundred in the rate at which this research may be advanced. The vitality of nuclear physics in the next generation will depend strongly on a facility with the range of capabilities promised by the LAMPF II project. A continuing renaissance in nuclear physics will be within the reach of LAMPF II.

LAMPF II will provide the principal facility for the next generation of profound tests of forefront questions in particle physics. The very high-energy facilities now being exploited or planned by the high-energy physics community directly access particular higher mass scales that could provide answers to key questions about the family structure of fundamental particles and the possibility of constructing a unified treatment of the forces of nature. Some or all of the answers sought may lie outside the mass ranges opened by these accelerators. The history of nuclear and subnuclear physics has taught us that sensitive experiments at relatively low energies can often uncover the critical clues of high-energy phenomena.

Striking examples exist in which current limits on the rates of many rare decays of muons, pions and kaons, in which family number is not conserved, now provide experimental evidence not possible with high-energy facilities. The current limit on one rare kaon decay can be used to bound the mass of the boson

which would mediate family-changing interactions above 10^5 GeV. This mass is already an order of magnitude higher than the mass scale to be probed by the Superconducting Super Collider! The weak bosons, which carry known parity-violating weak currents, were directly observed in 1983 by a high-energy experiment, with masses in the 80- to 100-GeV range. The bosons which may serve to restore parity symmetry at very high unification energies must be far more massive to account for the maximal violation of parity symmetry seen at low energies. The current lower limit on the right-handed W boson is 380 GeV, set not by a direct search at a high energy laboratory, but by precise studies of ordinary muon decay at a high-intensity low-energy muon beam!

The LAMPF II booster and main ring accelerators are of relatively simple design, incorporating a number of technological innovations. The high energy of the main ring beam, 45 GeV, will be unique in affording the flexibility required to address quark and gluon degrees of freedom in nuclei. These accelerators will be the foundation for a new generation of research.

The experimental facilities include two new areas, providing high-energy hadron beams and neutrinos. The existing LAMPF meson area will be upgraded, providing facilities for medium-energy hadron beams.

Construction of the LAMPF II accelerators will cost \$280 million and experimental facilities for the initial research program will cost an additional \$172 million, including the cost of upgrading existing LAMPF experimental areas. With construction commencing in FY 88, neutrino and pulsed muon beams will be available for initial use in FY 92, and the full facility will be operational in FY 95, providing an unprecedented combination of high-purity, high-intensity hadron beams.

LAMPF II will provide the physics community with a diverse spectrum of capabilities to probe fundamental questions in strong-interaction physics and in electroweak physics. In this summary, we survey the principal physics topics to be addressed by LAMPF II, the conceptual design of the accelerator and experimental facilities, the project costs and schedule, and the unique resources which Los Alamos National Laboratory will bring to bear on this project.

1.1 SCIENTIFIC AND TECHNICAL SUMMARY

1.1.1 Strong-Interaction Physics

Nuclear physics is predominantly concerned with the study of the response of nuclei to external stimuli, in an attempt to uncover the way in which nuclei are bound together and the nature of the constituents from which nuclei are assembled.

Nuclei can be viewed as consisting of nucleons bound by a potential which is described by the exchange of mesons. The individual nucleons and mesons are known to consist of quarks, bound by the raw strong interaction, due to the exchange of color gluons. During the last decade, the principal force in nuclei, which is the strong interaction, has been described by a promising gauge theory, quantum chromodynamics (QCD).

With both the fundamental constituents and the field theory apparently known, one might postulate that a detailed description of nuclear physics is merely dependent upon an extreme effort to produce sufficiently detailed calculations. This is far from the truth, for the constituents and the candidate field theory provide only the initial tools for the description of a preeminent problem in many-body physics.

As an example, knowledge of the field theory of electromagnetic physics, quantum electrodynamics (QED), did not lead to a description of superconductivity. Both experimental evidence and a picture which elucidated the appropriate collective degrees of freedom were required. The strong-interaction many-body problem in nuclei, which is much more complex than the problems addressed by QED, will only be solved by a broad, structured program of experiment using high-energy electrons, relativistic heavy ions and high-intensity light hadron probes, and by a theoretical confrontation with the resulting data using nuclear models based on traditional pictures and on QCD. Furthermore, a fundamental understanding of the theory should clarify the details of the transition between the ranges of validity of the different pictures (nucleonic, mesonic, QCD). The long-range strong interaction which binds nuclei should then be clearly seen as the collective result of the raw strong interaction between quarks, in the same way that molecular binding is the collective result of the pure Coulomb force between nuclei and orbital electrons.

The traditional view of nuclear behavior uses only nucleons, the protons and neutrons, interacting via two-body potentials. Physics often must resort to such simplifications, trading away detail (and momentarily intractable problems) for clarity and utility. This model does describe many characteristics of nuclei, but relies upon effective potentials and phenomenological models with adjustable parameters. This approach conceals flaws, and makes such models uncertain tools to use in predicting nuclear behavior in unexplored physical regimes. When discrepancies occur, it is not known whether the difficulties originate in the models, or because contact has been lost with more fundamental underlying structures.

Nuclear physics has long treated the potentials in which nucleons are bound as arising from a strong force carried by meson exchanges with the lightest mesons, pions, dominating this exchange interaction. Thus, on a

distance scale larger than the nucleon size, meson exchanges have been the basis for the theory.

In the last decade, very strong evidence has been amassed that the mesons themselves play a direct role in the description of the nucleus, even at low energy. Nuclear matter displays a combination of nucleon and meson degrees of freedom.

At even smaller lengths, a new scale (shorter wavelength and correspondingly higher probe energy) has emerged in which the advances of particle physics are applied in developing a picture of the nucleon and meson constituents. Both types of hadrons are known to consist of quarks, confined by a potential generated by the exchange of electrically neutral color gluons. In extremely energetic collisions with hadrons, these quarks appear to behave as if they were free (asymptotic freedom) but at more moderate energies the quarks appear confined. This is just the opposite of the picture arising in quantum electrodynamics, in which two charged objects are coupled most strongly at short distances, and appear free from their mutual influence at long distances. Quantum chromodynamics, which appears to describe the phenomenological behavior of quarks and gluons, is being addressed in its simplest form by high energy physicists in experiments which isolate the pure color dynamics of asymptotically free quarks; the short distance, high-energy regime of QCD physics.

This approach ignores the fundamental physics essential for the understanding of normal nuclear matter, the regime in which quarks are still confined. This is an intermediate distance scale, between the extremely short distances of asymptotic freedom and the long and medium distance scales associated with the nucleon and meson pictures. This sub-asymptotic regime is the frontier of strong-interaction nuclear physics, which can be studied by beams of electrons, heavy ions and light hadrons at appropriate energies. However, as indicated earlier, only light hadronic probes provide the capability to selectively probe both the quark and gluon content of nuclear matter. LAMPF II will permit a broad assault on this exciting new frontier.

The two main questions to be addressed by nuclear physics in the next decades are

- the understanding of the relevant degrees of freedom required to explain the nucleus, and
- the nature of quark confinement in nuclear matter.

LAMPF II will provide the nuclear physicist with the three principal tools required to address these questions. These tools will be used to do

- insertion experiments which make use of a "test charge" in the form of the strange quark to probe the nuclear volume;
- scattering experiments which make use of a gamut of traditional hadron probes (pions, protons, kaons, etc.) to sample the variety of length and time scales which characterize the hadronic response of nuclear matter; and

- deposition experiments in which the high energies and momenta available in LAMPF II beams are deposited in nuclei, creating extreme conditions.

In applying each of these three tools, a number of specific research areas deserve identification. It is always difficult to specify particular experiments to be carried out a decade hence, but having described the capabilities afforded by LAMPF II, today's perspective provides a guide to how one might begin these exciting studies.

1.1.1.1 Inserting the Strange Quark Into the Nucleus. Hypernuclear physics is the study of the nucleus with a new dimension. LAMPF II will permit a major advance beyond conventional nuclear studies, limited to neutrons and protons, or, alternately, up and down quarks. It will become possible to carry out sensitive studies of nuclei in which one or more nucleons is replaced by a hyperon (a strange nucleon), or, in the more modern constituent description, in which one or more strange quarks are embedded among the up and down quarks in the nucleus. LAMPF II kaon beams can be viewed as strange quark beams. The new nucleon created in such studies differs from naturally occurring nucleons and thus populates a new realm of nuclear states and produces excitations with distinct signatures. A new world of nuclear matter can be created and tested against meson, nucleon and quark descriptions.

Contemporary nuclear problems accessible in LAMPF II hypernuclear studies include testing for the existence of stable, doubly-strange dibaryons: states in which two strange nucleons are created, bound together, or alternately, assemblies of six quarks including two strange quarks. Such dibaryons are predicted by nuclear "bag" models, major contenders to explain quark confinement. LAMPF II beam intensities are likely prerequisites for the successful search for such a rare phenomenon.

The anomalously weak spin-orbit force felt by lambda hyperons bound in certain hypernuclei has led scientists to seek experiments which might distinguish between the traditional hadron and modern quark pictures. These alternate descriptions predict very different sigma hyperon spin-orbit forces. The flux and beam purity provided by LAMPF II will permit the high resolution spectroscopy needed to test these predictions.

An exotic form of ${}^5\text{He}$, containing a bound lambda hyperon, has been found to have a peculiarly weak binding energy. Promising but untested justification from emerging quark models has been offered, perhaps not explicable in traditional hadron pictures. Again, we will likely develop a real explanation only with the unique beams of the LAMPF II facility, testing the consequences of such explanations in other systems.

Extremely narrow states in sigma hypernuclei, hypernuclei with strangeness greater than unity, and studies of the strangeness-changing part of the nucleon-nucleon weak force are other examples of the exotic research which address the frontiers of nuclear physics, for which the intense meson beams of LAMPF II are needed.

1.1.1.2 Scattering LAMPF II Hadron Beams from Nuclei. The unfolding of the assortment of length and time scales in nuclear responses will be advanced dramatically by the set of intense, pure hadron beams which LAMPF II will

deliver. The choice of hadron probe will select the nuclear feature to be tested. For example, unmatched intense beams of antiprotons will be useful in studying the relativistic approach to ground state properties of nuclei, a treatment which has grown out of recent meson factory research with protons. The formation of pions within nuclei, resulting from antiproton annihilation, can be compared to ordinary coherent production from nuclei with pion beams, revealing whether these annihilation pions mimic "real" pion behavior and leading to the unfolding of the space-time evolution of hadronic interactions. To carry out this structured program, in which time scales of the order of 1 fm/c are investigated, pion momenta up to 1.5 GeV/c are required. The necessary antiproton momentum is 8 - 10 GeV/c. Only a facility with the high energy and high intensity of LAMPF II will be able to provide pure, high intensity beams at these momenta.

LAMPF II K^+ beams will provide a strikingly different projectile. Unlike the strong annihilation of the antiquarks in an antiproton entering a nucleus, the strange antiquark in the K^+ encounters no real (valence) strange quark in nuclear matter and thus annihilation does not take place. It remains stable and active, therefore, to probe the entire nuclear volume, and its comparatively weak scattering can be treated with precision. Even more tantalizing are the recent results which indicate that low-energy K^+ -nucleon scattering is well described by nucleon quark distributions. This unprecedented example of how a low-energy interaction tests modern constituent models implies that the K^+ can reveal the quark densities in a nucleus, not just a nucleon, directly. The duality of quarks in nucleons and quarks in nuclei can, thus, be confronted with the high purity, high flux K^+ beams to be produced at LAMPF II.

The nature of quark confinement and the quark-gluon degrees of freedom in nuclear physics will be uncovered in several crucial research programs at LAMPF II. Recent striking evidence for quark degrees of freedom in nuclei has emerged in the deep inelastic scattering experiments of the CERN European Muon Collaboration (EMC). These experiments have shown that quark distributions in nucleons do indeed depend upon the nucleus in which they are measured, demonstrating in a convincing way that the nuclear environment influences quarks. LAMPF II can be used to explore this exciting insight with a complementary study of the Drell-Yan process. In this process, a quark (or antiquark) from a beam hadron annihilates with an antiquark (or quark) from the target, leading to the production of a virtual photon which subsequently decays into a pair of leptons. It has recently been shown that a high-precision study of the Drell-Yan process performed in a manner similar to the EMC experiment (measuring the anomalous dependence of this process on the atomic number) could lead to a definitive choice among competing models of the EMC effect. The key to this sensitivity is the complementary way in which the virtual sea and real valence quark distributions of the target enter in the cross sections for deep inelastic lepton scattering and the Drell-Yan process. Proton beam energies near the LAMPF II energy of 45 GeV are excellent for these studies. In fact, the energy dependence of the appropriate range of the dilepton cross section favor this energy rather than higher energies.

Dilepton production using LAMPF II K^+ and K^- beams ($\bar{s}u$ and $s\bar{u}$) can uncover the poorly understood virtual sea quark distributions in nuclei. The K^+ probes the sea primarily, but the K^- probes both the valence and sea quarks. The intensity of LAMPF II beams will facilitate the precision comparative studies needed to do these experiments.

The existence of bound gluon states, called gluonium, which are mesons constituted without quarks, is one of the basic predictions of QCD and the search for such mesons will be the subject of a new forefront hadron spectroscopy at LAMPF II. Only with the dedicated facilities to be provided will a comprehensive study of this subject, fraught with small cross-sections and large backgrounds, advance in a meaningful and timely manner. The novel nature of gluonium makes it an essential part of the advancement of hadronic nuclear physics.

1.1.1.3 Depositing Energy in Nuclear Matter With High Energy Beams. The response of nuclear matter at extreme temperature and density has been identified in many forums as a particularly attractive way to address quark physics. Under some combinations of these conditions, a startling transition of nuclear matter to a plasma of quarks and gluons is predicted. The identity of individual nucleons vanishes. Relativistic collisions between heavy ions have been identified as a promising way to search for this postulated plasma. Antiproton annihilations in nuclei are a complementary way to achieve extreme conditions in nuclei. The release of annihilation pions energetically heats a small volume of nuclear matter. This heat spreads to the rest of the nucleus, cooling the original hot region. A region that is large enough and still hot enough to see the plasma effect is expected to appear at energies above those studied to date, that is for antiproton momenta above 2 GeV/c.

Detailed intranuclear cascade and hydrodynamical calculations strongly suggest that annihilation of antiprotons and antideuterons in this energy range would heat a portion of a target nucleus into the regime where the quark/gluon phase transition is predicted to exist. Even if this spectacular transition eludes the expectant nuclear science community, the study of nuclei at the accessible temperatures and densities are of substantial interest.

1.1.2 Electroweak-Interaction Studies of the Standard Model and Its Extensions

The study of electroweak interactions, an important frontier of particle physics, will be the second major arena for LAMPF II research. The problems traditionally probed by low and medium energy particle physics, long pursued in parallel to the activities at high energy laboratories, will expand greatly. The Superconducting Super Collider, the Tevatron and SPS, and the SLC, LEP and HERA will pursue higher mass scales in the hierarchy of particles. LAMPF II will be the stage for an unmatched pursuit of the low energy manifestations of these higher mass scales, and for significant tests of the dynamics of the standard model, a vital complementary thrust forward.

The principal advance of the last generation of particle physics is the unification of the electromagnetic and weak forces in a single theory. Comparable to the landmark derivation of Maxwell's equations in the late nineteenth century, in which the electric and magnetic fields were united, the electroweak theory is consistent with all known physical observations, especially the long-awaited observation in 1983 of the quanta of the weak interaction, the W and Z bosons, which together with the photon, mediate the entire known range of electroweak processes.

Just as the classical synthesis of Maxwell broke down only a generation later in the face of radioactive decay and other new phenomena, we can expect to

find that today's theories require extension. LAMPF II will test the limits of the electroweak theory in crucial ways.

Physicists have moved rapidly to attempt the combination of the candidate gauge theory of the strong interaction, QCD, with the electroweak theory. This theory, constructed with minimal assumptions (no assumption about nature not required by experimental evidence), is known as the standard model. In the standard model all of nature's mesons and baryons, and their strong, weak and electromagnetic interactions, can be described by combinations of nature's building blocks, the fundamental fermions, bound by nature's glue, the mediating color gauge bosons. These are the bosons that carry the strong interaction, the gluons, and the W, Z and the photon, which carry the weak and electromagnetic interactions. Our model is a minimal one, though. Outstanding questions confront us. We do not understand why there are three gauge groups (at least) of these bosons (arranged in mathematical groups called $SU(3)_C$, $SU(2)_W$, and $U(1)_B$ and why these objects have the particular masses they exhibit.

The fermion building blocks appear to fit into three known, and curiously symmetric families, as well. Pairs of quarks (up and down, strange and charm, top and bottom) and leptons (electron and electron neutrino, muon and muon neutrino, tau and tau neutrino) can be organized into three sets. Are these groupings truly meaningful? Why are some transitions among these families observed, and others never appear? Why do these objects have their observed masses or are the neutrinos truly massless?

Theories of grand unification postulate that all these separate groups collapse together at some supreme energy. Attempts are being made to treat quarks and leptons as interrelated. The question of a possible underlying relation between the bosons and fermions has been raised. Are they fundamentally different? Are these objects themselves constructed of more basic constituents? Can nature's fourth force, gravity, be included in an even more grandiose synthesis? At its most reductionist extreme, the goal of particle physics is a single, elegant, and concise theory which explains, at least in principle, all of the physical phenomena in the universe.

A frontal assault on these questions is being made by the high-energy particle physics community. High energy accelerators may directly produce new higher mass objects that may clarify the organization of the pantheon of elementary constituents. Higher energy research will likely clarify the transition region in which some of the separation between families and forces is reduced and nature's anticipated underlying symmetry is restored. Complete surprises will almost certainly emerge.

A pursuit so compelling, the very cutting edge of physics, should not be left only to these new high energy facilities. Even within their enormous energy range, the key insights may be experimentally inaccessible. Experiments at non-accelerator facilities and low- and medium-energy laboratories will be invaluable, indeed essential, in uncovering the subtle signatures of higher energy phenomena. Some studies will be unique to these laboratories.

The most immediate examples of the utility of LAMPF II in particle physics are the study of decays that violate family conservation, CP-nonconservation, and neutrino mass measurements. LAMPF II beams will provide fluxes of muons,

kaons and pions ideal for extending searches for family-nonconserving decays of these particles beyond present limits by many orders of magnitude.

The kaon beams at LAMPF II will be spectacular, providing flux two orders of magnitude more intense than previously available, with great flexibility to achieve purity and optical control. The current limit on the decay $K_L \rightarrow \mu e$ is $< 2 \times 10^{-9}$. If the process $K_L \rightarrow \mu e$ were mediated by a heavy family gauge boson, the current limit on this decay would set a lower limit of 10^5 GeV on the mass of this boson. As we have pointed out earlier, this mass is already an order of magnitude higher than the scale directly accessed by the SSC. Experiments are under way that hope to probe the 10^{-12} range in a 2000-h run. If 10^{-12} were reached, this limit (sensitive as the $1/4$ power of the branching ratio reached) would be pushed to about 7×10^5 GeV. The proposed 2000-h search could be accomplished in one day of running at LAMPF II! A one-year run could reach 10^{-14} , corresponding to a boson mass limit of 2×10^6 GeV! More realistically, detectors and data acquisition systems for such experiments will probably be developed after the first generation of LAMPF II experiments, in which the raw flux available would be used instead to tailor beam phase space and deliver truly pure kaon beams to researchers. LAMPF II will provide the capability and flexibility for all of these possibilities.

Neutrino masses stand out from the general particle mass enigma because historically neutrinos were thought to be massless, although it is now clear there was no fundamental reason for this presumption. One class of LAMPF II experiments would search for massive neutrinos in decays such as $\pi \rightarrow \mu \nu$ and $K \rightarrow \mu \nu$, in which the decay lepton energy distribution might display structure. A distortion of the π^+ spectrum in $K^+ \rightarrow \pi^+ \nu \nu$, caused by massive neutrinos, would be another appealing indicator. In each of these experiments, the extraordinary flux, purity and optical flexibility of LAMPF II beams will dramatically improve the prospects for new physical advances.

LAMPF II neutrino beams will provide an unequalled opportunity to search for the subtle oscillation of specific neutrino types into other neutrino types. This oscillation is a likely occurrence if neutrinos have mass. The experiments search for changes in the proportions of various neutrino types along the beam direction. The high flux and continuous availability of the neutrino beams at LAMPF II will be a major advantage over current neutrino beams, because the neutrino oscillation signal will likely be observable only with many detected events and the smallest possible backgrounds.

In addition to searching for entirely new physics phenomena which may force extensions of the standard model, LAMPF II beams will be a prime source of precision tests of the standard model. One dare not assume that the standard model is complete and accurate and so it must be confronted with ever more stringent challenges.

The scattering of muon neutrinos on electrons is a clean, direct measure of the leading parameter of the electroweak theory, the electroweak mixing angle θ_w . LAMPF II will be the first opportunity for real precision in this area, with the possibility of event samples with 10^4 scatterings observed. Current experiments are accumulating a few hundred events. High energy accelerators will refine our knowledge of the masses of the W and Z bosons, and θ_w at high energies. LAMPF II results will complement this work and permit the energy-dependent radiative corrections in the standard model to be tested. This

is the only known quantitative test of the field-theoretic aspects of the electroweak theory. It is comparable in significance to the study of $g-2$ in the early days of quantum electrodynamics (QED).

CP nonconservation is one of our generation's greatest mysteries. It is only accessible with precision in the kaon system (strange quark decays) for heavier quark decays are far less copious at the world's accelerator facilities. The riddle of CP nonconservation has withstood twenty years of research. The possibilities at LAMPF II, with beams yielding more than 10^9 kaons per second, may be the only way in which the origin of CP nonconservation may be unmasked.

The ultimate frontier of particle physics is the physics of the first instant after the Big Bang. During that instant the mass scales and the symmetries of nature took shape. LAMPF II will improve our knowledge of the early universe by better understanding of nature's fundamental symmetries.

1.1.3 The LAMPF II Accelerators

The LAMPF II accelerators will serve as an extraordinary kaon factory. One remarkable asset, the existing LAMPF linac, is superior to any existing or proposed injector because it provides a high injector energy (800 MeV), compact emittance and because this extraordinary linac already exists. It provides a convenient 120 macropulses per second, delivering both H^+ and H^- particles.

We propose to inject every other LAMPF macropulse into the booster accelerator, the others being reserved for the existing Proton Storage Ring. The booster will accelerate the beam to an energy of 6.0 GeV. The 60 pulses per second will orbit the 331-meter ring circumference.

Beam will be extracted in a single turn from the booster with four pulses in sequence transported to the main ring. While the main ring accelerates these pulses, the next 16 booster pulses will be sent to the neutrino target area (Area N). In this sequence, all of the 170- μ A booster beam is continuously used to do physics research, with 48 of every 60 pulses available for neutrino and pulsed muon physics.

The main ring, approximately 1324 meters in circumference, will provide 45 GeV protons, with an average current of 34 μ A at 3 Hz. The beam will have a microstructure of pulses less than 2 ns long, and separated by 16 ns, ideal for timing measurements in many experiments.

The beam from the main ring will be extracted slowly and transported to an area providing beams of kaons, pions, and antiprotons with energy less than 10 GeV (Area A), and an area in which hadrons above 10 GeV will be the projectiles (Area H).

The design of the LAMPF II accelerators has been guided principally by the physics goals of the project, but substantial attention has been paid to producing a cost-effective and reliable design. Though the design of the accelerators is relatively conservative, requiring no techniques beyond the state-of-the-art, an effort has been made to use the latest techniques available and to make a number of innovations designed to lower costs and promote reliability. Novel perpendicularly biased ferrites in the rf accelerating cavities, an advantageous magnet power supply design, and careful attention to

the details of the magnet lattice have all contributed to lower power requirements, lower costs, and reduced operational complexity.

The principal design feature in the accelerators is the use of rapid cycling to achieve high average beam intensity while avoiding high intensity in each pulse. Rapid-cycling synchrotrons have been built elsewhere and operated successfully. We have, therefore, made appropriate use of an existing technology.

The high main ring energy, 45 GeV, will make LAMPF II an exceptionally powerful and versatile high-intensity facility. In addition to the obvious utility of the energy in particle physics research, this energy will afford maximum flexibility in carrying out the program of quark/gluon nuclear physics. While beam power is the leading factor in particle production rates, our choice of higher energy at constant beam power yields a marked advantage.

LAMPF II is the only kaon factory design proposal to date employing only two accelerator rings. The simplicity of the LAMPF II design makes facility upgrades in current or duty factor dependent only upon the addition of a stretcher ring. Thus, the long-term utility of LAMPF II is assured.

1.1.4 LAMPF II Experimental Areas

LAMPF II will provide three powerful experimental areas:

- Area A will receive the main ring beam via slow extraction and will be a substantial revision of the existing LAMPF meson area. It will provide two primary beam targets and an array of low-energy secondary beams and spectrometers.
- Area H will receive the main ring beam via slow extraction and will be a new high-energy area with two primary targets.
- Area N will receive the booster beam via fast extraction, providing facilities for neutrino and pulsed muon physics.

The design of these facilities emphasizes

- high beam intensity;
- high beam brightness;
- high beam purity;
- high resolution for nuclear physics;
- high beam availability;
- multiple ports;
- complete coverage of all the beam energies and particle types available at LAMPF II, including ν , K, π , μ , $\bar{p}$, and polarized protons; and

- flexibility to meet the requirements of future physics experiments.

The areas are conservatively designed, building especially on the extensive experience at LAMPF in transporting and targeting high-intensity beams. Attention has been paid to reducing costs, through the use of such techniques as underground areas to reduce shielding costs, shared primary beam target cells, and careful choice of remote-handling technology.

The beams to be provided will include

- Area N:
neutrino source, with pion focusing horn;
pulsed muon channel.
- Area A:
0.4- to 0.8-GeV hypernuclear and stopped kaon beam;
0.7- to 1.2-GeV/c kaon beam and spectrometer;
1.0- to 2.0-GeV/c kaon beam and spectrometer;
2- to 10-GeV/c kaon and antiproton beam;
wide-band neutral kaon beam;
stopped muon channel;
test beam.
- Area H:
35-GeV separated kaon and antiproton beam;
35-GeV unseparated kaon and antiproton beam;
double-arm spectrometer area.

These beams include a capability for research with each particle type, over all accessible energies, at LAMPF II.

1.1.5 Project Schedule and Costs

Detailed cost estimates and project schedules have been prepared by Los Alamos personnel and external consultants. This proposal is supplemented by companion documents prepared by our consultants.

A consulting team from Science Applications International Corporation has prepared a complete work-breakdown and cost estimate, including EDIA, installation, project management, and 25% contingency. The costs, in January 1985 dollars, are

• Booster	\$120,000,000
• Main Ring	160,000,000
• Experimental Area N	29,000,000
• Experimental Areas A and H	143,000,000

The construction schedule contained in the SIAC Management Plan and Schedule Report, and in the Preconceptual Site Plan prepared by Parsons, Brinkerhoff, Quade, and Douglas call for construction to be completed in 1994,

given a 1988 start. Neutrino area beams will be usable in 1992, with all beams in production in 1995.

1.1.6 Unique Resources of Los Alamos National Laboratory and the User Community

LAMPF II will provide the nuclear science community with an essential capability for its attack on the frontier problems of nuclear and particle physics. To be assured that this powerful capability is developed with utmost efficiency, building on scientific, technical and managerial capabilities essential to the endeavor and already in place, the resources, people and facilities and expertise at Los Alamos, will be crucial. The scientific accomplishments of the nuclear and particle physics programs at LAMPF and the development of the scientific interests of this community are consistent with the goals of the LAMPF II program.

The Clinton P. Anderson Meson Physics Facility (LAMPF), the largest of the "meson factories," and the single largest facility in medium energy physics, has had as one of its major goals the bridging of the gap between nuclear and subnuclear physics. This goal was formulated during the early 1960's and LAMPF has carried out a broad program of ambitious research in both areas. That the frontier of nuclear physics is once again intertwined with the intellectual product of high energy physics is not an accident. That the proposed extension of the LAMPF facility described in this document is once again concerned with the overlap of these two fields is a natural extension of the scientific and technical activities of the medium energy research community. Indeed, in recent years an extraordinary dialogue between the two poles of medium-energy physics has taken place in a series of important workshops at Santa Cruz in 1983, Steamboat Springs in 1984, and at laboratories in Canada, the United States, and Europe during the last five years. LAMPF II will provide this community with the facilities and tools needed to probe the dynamics of nuclear matter with light hadronic probes and to exploit nuclei and medium energy beams and techniques to test the limits of the models provided to us by high-energy particle physics. Thus, LAMPF II is the intellectual extension of the present-day medium energy physics effort.

Furthermore, the technical challenges which must be met to construct and operate a high-energy, high-intensity proton accelerator laboratory are dominated by the difficulties presented in the transport and targeting of high-intensity beams. The operation of LAMPF II will require the technologies painstakingly developed over the years at LAMPF and other meson factories, far more than those found at the existing high-energy, low-intensity facilities, to insure the utility and reliability of the facility. The technical expertise at LAMPF in materials science, radiation-hardened components, remote handling, beam transport and shielding are unique national resources. Indeed, these abilities are well known to define the useful limits of any high-intensity accelerator laboratory. LAMPF has led the world meson-factory community in each of these areas.

In addition to the technical and operational experience acquired in operating a high-intensity laboratory for more than fifteen years, Los Alamos possesses unique assets. Embedded in one of the nation's largest and most capable multidisciplinary laboratories, the medium-energy physics program at Los Alamos has long benefited from the extraordinary collection of scientific and technical assets which have been available. The medium-energy program has

supported, critiqued and reinforced, in crucial ways, these other programs, as well. More specifically, the LAMPF accelerator surpasses any other facility as the injector for a new higher-energy, high-intensity facility. The existing LAMPF beams, and the new Proton Storage Ring (PSR) provide unmatched facilities for developing and proving prototype systems for the LAMPF II accelerator and experimental areas.

The intellectual challenges which have been pursued by the community of researchers using the Los Alamos facilities have been met with notable success. A substantial record of achievement in scientific areas bridging the nuclear and particle domain has been accumulated. The natural development of this pursuit has quite properly followed the growth of new interests and outlooks in nuclear and particle physics. The intellectual goals of LAMPF II are the goals of the medium energy physics community.

Apart from its intrinsic scientific merit, LAMPF II will continue our contribution to the nation, by providing an important part of the knowledge and people base essential to our nation's security and well-being.

CHAPTER 2.

STRONG INTERACTION PHYSICS

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2. STRONG INTERACTION PHYSICS

2.1 OVERVIEW

The nucleus has traditionally been viewed as a collection of individual nucleons interacting via two-body potentials. This picture indeed describes many of the general characteristics of nuclei. In order to achieve agreement with experiment, however, one must often resort to effective potentials and phenomenological models. While this approach is successful in correlating an impressive body of nuclear data, it often conceals flaws by subsuming all corrections into the effective parameters. Furthermore, these models provide an uncertain base from which to predict nuclear behavior in physical regimes not yet encountered. In those cases where accurate tests of our understanding of nuclear phenomena can be made, substantial discrepancies sometimes occur. As examples of these problems, the binding energy (or form factor) of the trinucleon system cannot be calculated, the Coulomb energy difference is observed to be 10% larger than calculated throughout the entire periodic table, and the longitudinal structure function for quasi-elastic electron scattering is 20% smaller than simple theory would have it. Are these persistent problems in the domain of classical nuclear physics to be laid on the doorstep of deeper structure of the system? While partial solutions to these problems may exist within the purely nucleonic framework, recent convincing and useful discussions have been in terms of alternate degrees of freedom.

It has been known for some time that the long-range part of the effective nucleon-nucleon potential is due to the exchange of mesons (primarily pions). In the last decade very strong evidence has been amassed indicating that these mesons (as well as the resonances they form with nucleons) play a direct role in the description of the nucleus. This evidence is generally considered as proof that degrees of freedom beyond nucleons are required to describe the nucleus.

During the same era in which nuclear physics was establishing the importance of meson currents in nuclei, a revolution was occurring in particle physics. There now exists a candidate gauge theory of the strong interaction (based on quarks and gluons) known as quantum chromodynamics (QCD). The recent emphasis in a large sector of hadronic physics has been to focus on high energies and highly inelastic events where this theory is expected to have its simplest form and where the fundamental principles should be most easily tested. While the logic of this approach is clear, it unfortunately bypasses the investigation of the consequences of this theory in the crucial regime appropriate to nuclei. The aim of the ultra-high-energy approach is to circumvent the complication due to the confinement of quarks in order to reach regions where the basis of the theory is clear (where the quarks appear "free"). It tends to ignore one of the most interesting aspects of the theory, the confinement mechanism itself. We note here that leptonic probes interact only with quarks. Hadronic probes must be used to directly access the gluon degrees of freedom.

It must be understood that exact calculations based directly upon the fundamental theory of QCD at energies relevant to nuclear physics are prohibitively difficult at present. As a result secondary approaches are

employed, and effective interactions with basis in more fundamental theory are used. A principal task of nuclear physics is to test these models and provide guidance for their development. It is expected that an increasingly more faithful realization of the consequences of QCD will be developed as time progresses.

There are two tasks pointed up by the developments of the last two decades:

FIRST: Delineate and investigate the degrees of freedom relevant for the description of nuclear processes. At present the three classes generally considered are (a) nucleons, (b) nucleons plus mesons and (c) quarks and gluons. The appropriateness of the different approaches and different degrees of freedom will vary according to the energy regime and the observables under consideration. For example, nuclear structure as described by the shell model involves primarily nucleonic degrees of freedom whereas deep-inelastic scattering requires the use of meson and/or quark-gluon degrees of freedom.

SECOND: Investigate the mechanism of confinement and its relationship to nuclear physics. Here LAMPF II offers a unique opportunity for the thorough development of this field. The many two-baryon reactions allow a test of basic models of QCD. The structures of meson-meson, meson-baryon and baryon-baryon states are fundamental data on which understanding of confinement will be based. The appropriate beams at the appropriate energies will be available at LAMPF II. After the basic models are formulated their deeper structure and relevance to nuclear physics can be obtained only by use in a nuclear environment.

The three principal tools available to the nuclear physicist are: the introduction of a "test charge" (the strange quark) into the nucleus; the use of many different probes (pion, K^+ , K^- , proton, antiproton, neutrino, etc.) including the electron to sample the different parts of the hadronic components of the nucleus at different length and time scales; and the use of high energy/momentum or high energy density to study the response of the nucleus under extreme conditions.

The remainder of this section on the strong interaction is classified according to these three emerging technologies of basic research in nuclear physics. Immediately following are brief discussions of these three areas. A more complete exposition is given in sections 2 through 4.

2.1.1 Hypernuclei

Investigations seeking the relevant dynamical variables with which to describe nuclear matter have been conventionally restricted to systems comprised of neutrons and protons, or alternatively of up and down quarks. However, there exist additional systems, those having strangeness. The study of nuclei in which one or more nucleons is replaced by a hyperon (a strange baryon such as a Λ , Σ , Ξ , etc.), or in which one or more strange (s) quarks are embedded in a sea of u and d quarks, extends our ability to look beyond systems normally found in nature. Utilizing a distinguishable test probe, a hyperon or an s quark which truly differs from the components of normal nuclear matter, permits one to discriminate between probe-target interactions and target component correlations. One can explore fully the consequences of the developing pictures of nuclear matter, the complementarity of the descriptions in terms of baryons and mesons and of quarks and gluons. The emerging models can be tested in a

world beyond that in which they were conceived to explain existing data. The goal is to determine which degrees of freedom provide the most economical description of nuclear phenomena, and which offer reliable predictions in new regimes.

The "H" Particle. Hypernuclear physics offers a possible direct test of the most interesting prediction of certain quark bag models--the existence of a stable, doubly-strange "H" dibaryon. The "H" particle is presumed to have the same quark composition as two Λ 's (ssuudd). Searches for such a massive six-quark object are of paramount importance. If found, the existence of such a peculiar "di-lambda" would constitute a triumph for the bag models and would serve as a severe constraint upon all quark models of the hadrons. However, the reported observation in old emulsion experiments of doubly-strange ${}^6_{\Lambda\Lambda}\text{He}$ and ${}^{10}_{\Lambda\Lambda}\text{B}$ hypernuclei raises an important concern: why shouldn't these double- Λ hypernuclei have decayed immediately into more stable di-lambda states? Is the "H" only barely bound with respect to the $\Lambda\Lambda$ threshold, in contrast to the bag model predictions? Does the "H" exist at all? Will the baryon-meson picture provide an adequate description of known phenomena? With current estimates for the cross section for the formation of this particle it is very unlikely that it will be found without a machine with kaon fluxes comparable to those of LAMPF II.

Spin-Orbit Interaction of Hyperons. An exciting feature of the (K^-, π^-) reaction data leading to p-shell hypernuclear states was the discovery that the mean-field spin-orbit force felt by the Λ is very small (< 2 MeV). Both quantum chromodynamic and one-boson-exchange potential model arguments have been proposed to explain this initially surprising result. However, these two different pictures of the small ΛN spin-orbit force give strikingly different predictions for ΣN and ΞN spin-orbit forces. Normalized to that of the NN interaction, spin-orbit force predictions are $(N:\Lambda:\Sigma:\Xi)$ 1:0:4/3:-1/3 and 1:0.3:0.4:0.15 for perturbative QCD and OBE models, respectively. Thus, hypernuclear physics offers yet another excellent opportunity to explore possible differences between the traditional hadron and the emerging quark pictures of the nuclear interaction.

Binding Energy of Hypernuclei. The anomalously small binding of the ${}^5_{\Lambda}\text{He}$ hypernucleus has been known for twenty years. Why is it small? In the baryon picture the Λ hyperon is distinguishable from the nucleons. All five baryons can coexist in 1s states to form the ground state of the hypernuclear system. (This is in contrast to the case of ${}^5\text{He}$ where only four of the five nucleons can reside in 1s states, and consequently the ${}^5\text{He}$ nucleus is unbound.) Nonetheless, simple model calculations, based upon ΛN interactions parameterized to account for the low-energy scattering data, overbind ${}^5_{\Lambda}\text{He}$ by 2-3 MeV. Explanations of this discrepancy have been sought in terms of (1) the known strong coupling between the Λ and Σ hyperons, and (2) tensor forces which bind the triton and alpha particle less, compared to the deuteron, than do central forces. Describing the system in terms of quark variables offers an intriguing alternative possibility. In such a picture only the s quark of the Λ (a uds composite) is distinguishable and can therefore coexist with the 12 other u and d quarks in a 1s state. The accompanying u and d quarks of the Λ are Pauli-blocked from the quark 1s state. Thus, one could speculate that the binding of ${}^5_{\Lambda}\text{He}$ should be smaller than that estimated from ΛN scattering data and the binding energies of ${}^3_{\Lambda}\text{H}$ and ${}^4_{\Lambda}\text{H}$ (where the u and d of the Λ are not blocked). The baryon and quark explanations, then, represent very different physics. But

are the two pictures really antithetical? Is the propensity of quarks to cluster in colorless triplets so strong (that is, is confinement so complete) that a baryon Pauli principle yields results almost identical to a quark Pauli principle? Detailed study of many such symmetric hypernuclei are needed before these questions can be answered.

Very Strange Hypernuclei. Finally, the creation of nuclei with strangeness -2 and beyond will lead to new vistas in hypernuclear physics. The (K^-, K^+) strangeness exchange reaction opens the door to the study of doubly-strange ($S = -2$) hypernuclei (Ξ , $\Lambda\Lambda$, $\Lambda\Sigma$, $\Sigma\Sigma$). Investigations of the spectroscopy of $S = -2$ hypernuclei represents a logical next step in the evolution of hypernuclear physics. Our only hope of determining the $\Lambda\Lambda$ interaction lies in the systematic study of the spectroscopy of $\Lambda\Lambda$ hypernuclei. That information is crucial to test our knowledge of the $SU(3)$ structure of the baryon-baryon force and to address the question of the properties of multi- Λ nuclear matter. The latter is speculated to exist in dense stars. Does the one-boson-exchange potential model that describes NN , YN , ΣN scattering extrapolate correctly to describe the $\Lambda\Lambda$, $\Sigma\Lambda$, $\Sigma\Sigma$, and ΞN interactions? Will multi- Λ hypernuclei be stable with respect to particle decay? Cross sections for production of discrete $S = -2$ hypernuclear states are expected to be small. Intense K^- beams will be required for serious investigations.

Summary. These are but four examples from hypernuclear physics that address the question of relevant degrees of freedom with which to describe nuclear phenomena and the closely related issues of understanding the basic strong interaction and the microscopic structure of nuclear matter. Many fundamental points will be addressed by the study of hypernuclear physics: the hyperon-nucleon force, charge symmetry breaking, few-baryon systems, microstructure of multihadron states, γ spectroscopy, high-spin states, metastable Σ hypernuclei, isospin mixing, Y^* resonances, $S = -2$ hypernuclei, charmed systems, etc.

The important results described above were obtained with machines and beams of only modest capabilities for the production of strange particles. A machine dedicated in large part to hypernuclei physics, providing kaon beams with much higher fluxes and considerably greater purity is essential to the timely development of this revolutionary area of nuclear physics.

2.1.2 Medium- and Large-Length Scale Physics

LAMPF II will be the source of intense beams of antiprotons, positive kaons, pions, and neutrinos in the 0.5 to 2 GeV range. The interactions of these particles with nuclei offer exciting new opportunities to extend present day medium-energy physics. Modest energy and momentum transfer imply that mesons and baryons, the degrees of freedom appropriate to "medium and large length scale" phenomena, will dominate. However, the composite nature of these hadrons may reflect itself even in this regime. Examples with implications for quark degrees of freedom are discussed below.

Other investigations will result in precision information about nuclear reactions and structure that are possible only with the wide variety of new intense secondary beams available from LAMPF II. Higher energy pions, positive kaons, and antiprotons have much to offer which is complementary to future studies of nuclei with electrons.

The EMC Effect--Extra Pions? Deep-inelastic scattering of 100 to 300 GeV muons at CERN has produced convincing evidence that the quark distributions of nucleons in nuclei differ in a small but significant way from those of free nucleons. Many models have been proposed to explain the EMC (European Muon Collaboration) effect. Among the most interesting possibilities proposed is that the enhanced lepton scattering from heavy nuclei may be the result of "excess nuclear pions." At the short distances, the enhanced pion field model leads to an increase in the number of quark-antiquark pairs from which high energy leptons can scatter. At length scales seen with lower energy probes, one should be able to observe the enhanced pion field itself.

LAMPF II beams of medium-energy (1-2 GeV) neutrinos offer a unique opportunity to study the consequences of this model by inclusive charged-current neutrino reactions on nuclei. It is necessary to use a weak probe to test this physics since the electromagnetic current does not couple directly to the pion field. An excess of pions will manifest itself as an anomalous A dependence (as in the EMC experiment) in the (ν_μ, μ^+) reaction near a momentum transfer of 2 fm^{-1} for nuclear excitation energies near 100 MeV. Because the cross sections for quasi-elastic scattering from nuclei are large by neutrino reaction standards, reasonable counting rates are obtainable even with the narrow-band neutrino beams required to obtain energy resolution of 100 to 150 MeV.

Pion Double-Charge Exchange and 6-Quark Bags. Pion double-charge exchange offers intriguing possibilities in the search for quark effects in nuclei. Two nucleons must be involved in this reaction in order for two units of charge to be transferred. The relevant length scale for seeing quark effects is about the nucleon radius, which means that nucleons in nuclei are often close enough for six-quark states to form. One may ask whether the double-charge exchange occurs dominantly through meson-baryon scattering, or dominantly through six-quark structures that lie outside this framework. Double-charge exchange may be used to take a "snapshot" of the nucleons inside the nucleus in an experiment in which two nucleons and a pion are measured in coincidence. The confrontation of these views with experiment will give insight into which (if either) is the correct perspective. Such experimental programs will grow to their full potential at LAMPF II where greatly enhanced fluxes of pions (factors of 10-100) will exist, along with increased duty factors and vastly improved data-taking capability.

Antiproton Reactions. The antinucleon holds a special place in the spectrum of probes. Recent proton scattering experiments have strongly suggested that a relativistic approach to nuclear theory is necessary, even at zero (or negative) energy. The implications of these results are that the virtual production of $N\bar{N}$ pairs and their propagation in nuclei are relevant to our understanding of even the ground state properties of nuclei. The study of the interaction of antiprotons with nuclei will be of fundamental importance in such approaches.

Since the antiproton annihilates to form (mostly) pions and since we know much about the interaction of pions with nuclei, the study of annihilation in nuclei will allow us to determine how pions are formed from hadronic matter. Coherent production from nuclei with high-energy pion beams has provided crude information about the time it takes for the quarks to form a physical hadron--the hadronization time. Antiproton annihilation therefore leads us to the fundamental issue--the space-time evolution of the hadronic interaction.

Investigations of interactions over these time scales (~ 1 fm/c) are made possible by both creating and scattering an antiproton within a single nucleus. This requires final pion energies in the range 0-1.5 GeV, implying antiproton energies of the order of 8-10 GeV.

K^+ -Nucleus Scattering--Neutrons or Quarks? In contrast to the antinucleon, whose antiquarks can easily annihilate when it encounters a nucleon, is the K^+ , whose strange antiquark cannot readily annihilate in a nucleus. For this reason the K^+N interaction strength is relatively weak and this probe can be treated more accurately in scattering theories. These properties also give it a unique position among hadronic probes. Certainly it can be used to measure neutron densities in a manner superior to any other probe.

There is speculation, however, that K^+ interactions may allow interpretation beyond the nucleon picture and bear directly on our view of quarks in nuclei. Recent studies suggest that low-energy K^+ -nucleon scattering is well described directly in terms of the quark distribution of the nucleon. If this is indeed true then the K^+ can be used as a direct probe of the quark density of the nucleus in a manner somewhat analogous to the way in which the electron probes the nuclear charge distribution. In a purely nucleonic picture the assumption is made that the quarks reside in nucleons and that, even in the nucleus, those nucleons retain the properties they exhibit in free space. This assumption is testable and provides a direct comparison between the nucleons-only point of view, and one in which quark degrees of freedom are essential. The experiments require beams of positive kaons of greater purity and much higher intensity than those presently available.

Summary. As an intense source of novel hadronic and electroweak probes, LAMPF II will open an entirely new field of nuclear research. Each probe has a special sensitivity to a particular characteristic of the nuclear medium, in much the same way as electrons are sensitive to only electromagnetic charge and current, and hence to only one property of quarks. The addition of kaons, antinucleons, neutrinos and high-energy pions to our arsenal will significantly enhance the ability of physicists to explore different views of the nucleus. Because the number of degrees of freedom needed to describe the nucleus has increased it is clear that more probes are needed to obtain an unambiguous delineation of the various components.

2.1.3 Quark/Gluon Physics

As QCD is a theory that is most readily implemented at high momentum transfer, the high-energy capability of this facility is extremely important. LAMPF II will provide beams of sufficient intensity and purity to study constituent (quark and gluon) scattering under conditions difficult to achieve at low-intensity facilities of equal or higher energy range.

Quark Distribution in Nuclei--The Drell-Yan Process. The EMC effect has shown us that nuclei are something beyond A interacting "free" nucleons. Many ideas have been directed toward understanding the anomalous A dependence of deep-inelastic lepton scattering. Deeper understanding of this phenomenon will require experiments which yield information complementary to deep-inelastic scattering.

One very promising area readily accessible with a high-intensity machine is the Drell-Yan process. Here a quark (or antiquark) from a beam hadron annihilates with an antiquark (or quark) from the target, leading to the production of a virtual photon which subsequently decays into a dilepton pair. It has recently been shown that a high-precision study of the Drell-Yan process performed in a manner similar to the EMC experiment (detection of an anomalous A dependence) could lead to a definitive choice among competing models of the EMC effect. The key to this sensitivity is the complementary way in which the sea and valence quark distributions of the target enter in the cross sections for deep-inelastic lepton scattering and the Drell-Yan process. The requirements for such an experiment are a high-intensity proton beam in the energy range 30 to 50 GeV. This range, rather than higher energies, is dictated by the $1/E_{\text{beam}}$ dependence of the dilepton cross section in the appropriate kinematic region. Further applications of this process in nuclei abound. For example, dilepton production from a K^+ beam ($\bar{s}u$) arises in first order only from the sea of antiquarks of the target. Reactions induced by a K^- ($s\bar{u}$) beam receive contributions from both valence and sea quarks. A comparative study of the two processes offers a direct method to study the sea-quark distribution in nuclei.

The New Generation Spectroscopy. An enormous amount of new spectroscopy awaits the intense secondary beams of pions, kaons and antinucleons provided by LAMPF II. Among the major goals are the searches for quarkless mesons--gluonium, and for narrow structures associated with the dibaryon and the baryon-antibaryon systems. This new generation of hadron spectroscopy almost certainly will deal with small cross sections in the presence of large background. The rewards are commensurate with the difficulty of the program, however. Resonances or bound states in the quarkless or multiquark sectors would represent a new form of matter. The ability of LAMPF II to provide for the entire gamut of experiments necessary for these studies makes it a complete nuclear physics facility.

Very High Momentum Transfer Physics. At asymptotically large momentum transfer perturbative QCD provides a simple description of hadronic scattering processes in terms of a small number of calculable Feynman diagrams. This simplicity translates into simple relations among experimental quantities. Observables in polarized scattering experiments, for example, become almost trivially simple because helicity is conserved in single-gluon exchange. In exclusive pseudoscalar meson-induced reactions on nucleons the number of independent reaction amplitudes is very limited, implying cross section relationships among the numerous possible reaction channels.

There is currently considerable controversy over the applicability of perturbative QCD to exclusive processes in any energy/momentum transfer region accessible to accelerators (including the SSC). LAMPF II will provide unique opportunities for experimental guidance in this debate. High-intensity polarized beams up to 45 GeV will be the source of polarized quarks necessary to test the limits of simple or more complex theories inspired by QCD. Reactions induced by high-energy pions and kaons at high momentum transfer will be feasible for the first time. These experiments will tell us how close we are to achieving the simplicity of the asymptotic regime.

Heating Nuclei with Antimatter. One of the most promising methods proposed to probe quark effects in nuclei is to subject nuclear matter to conditions of extreme temperature or pressure. Many calculations have predicted

that under such conditions normal nuclear matter will undergo a transition to a quark-gluon plasma in which individual nucleons lose their identity. If the transition for normal matter to quark-gluon matter is sharp, then the line representing the boundary of the phase transition (drawn in a temperature-density space) is expected to have a zero-density intercept around 180 MeV. The zero-temperature intercept is much less certain, with predictions ranging from ~ 3 times normal nuclear-matter density to 10-20 times normal density.

Colliding heavy-ion beams have been suggested as a method for achieving the conditions necessary for the formation of the postulated plasma. A complementary method of achieving extreme conditions is the annihilation of antiprotons in nuclei. The release of energy in the form of a number of pions rapidly heats a small volume of nuclear matter. This heat gradually spreads to the rest of the nucleus, cooling the hot region originally formed.

Studies to date on the annihilation of antiprotons in nuclei have been done only with low incident energies. In this case approximately 2 GeV is available from annihilation and only moderate energy deposition is possible. As the incident momentum increases above 2 GeV/c a combination of circumstances causes the energy deposition to become much more efficient. The relatively small volume in which this energy is absorbed leads to a high temperature and large energy density.

Detailed intranuclear cascade and hydrodynamical calculations strongly suggest that annihilation of antiprotons and antideuterons in the 8 to 10 GeV energy range would heat a portion of a target nucleus into the regime where the quark/gluon phase is predicted to exist. Such measurements are obviously of fundamental interest. All considerations of the quark-gluon plasma aside, measurement of the behavior of nuclear matter at moderate density and high temperature is important in its own right. The calculations indicate that antimatter annihilation is an excellent method for investigating the thermal degrees of freedom of the equation of state.

Summary. We close this overview with a few remarks on the necessity for constructing a facility in the energy and intensity range of LAMPF II. The major goal will be to explore a new and exciting frontier of nuclear physics. To achieve this we need not only a variety of medium-energy secondary beams but intense high-energy beams as well. In a single facility, then, one will have the capability to cross the boundary between meson/baryon and quark/gluon dominated phenomena. We believe that such a program would spark a renaissance in nuclear physics and that the extensive physics program outlined in this proposal will lead to a new level of understanding of the structure of the nucleus.

2.2 STRANGENESS IN NUCLEI

2.2.1 Introduction

Strangeness brings a new perspective to the study of nuclei. Pioneering hypernuclear programs have convincingly demonstrated the exciting physics content of this field even with the adverse experimental conditions of low beam quality and poor energy resolution. In this section we discuss the impact of LAMPF II, a true "kaon factory," on present and future directions in the study of strangeness in nuclei.

Subsection 2.2.2 examines what is known about the fundamental interaction of hypernuclear physics--the hyperon-nucleon force. The potential for new physics here is enormous, ranging from the production of the "H" particle, which tests our understanding of the binding of quark constituents, to the determination of the ΣN spin-orbit interaction from tagged-polarized Σ^- tertiary beams.

Subsection 2.2.3 reviews what is presently known about Λ -hypernuclei and what lies ahead. The traditional production reactions, (K^-, π^-) and (π^+, K^+) , can be supplemented with untested new ones such as (K^-, π^0) to produce lambdas in nuclei with 100 times present rates and a factor of ten better resolution. Only under these circumstances will one be able to judge whether nuclei with lambdas are lambda-hypernuclei or strange-quark nuclei.

In subsection 2.2.4, Beyond the Λ , a number of extensions of hypernuclear physics are examined. Whereas in some of the examples one might be able to obtain anecdotal experimental data using contemporary, low-flux machines, a serious program of hypernuclear physics can be undertaken only at a high-intensity proton accelerator capable of producing intense secondary beams. Without such a facility, none of the physics discussed here can be addressed in more than just a cursory fashion.

2.2.2 The Hyperon-Nucleon Interaction

2.2.2.1 The ΛN Force. Experimental data on ΛN (and ΣN) scattering are scarce (Po 82, Pr 80). A full phase shift analysis of the existing ΛN scattering data is not a practical possibility. Neither meson-exchange models nor quark-gluon models of the strong interaction have been sufficiently developed to permit accurate model predictions of the ΛN force. Thus, the ΛN (and ΣN) data have been grouped in a combined analysis with the more plentiful NN scattering data to constrain one-boson-exchange potential models of the $S = 0$ and $S = -1$ baryon-baryon force (Pr 80). Although this modeling procedure is subject to considerable uncertainty (there remain a number of unknown parameters in the description of the exchange of the vector meson nonet, and the scalar singlet and/or scalar octet or pairs of pions), these theoretical potentials provide our best representation of the limited data available. More extensive data are clearly needed. LAMPF II intensities will make feasible for the first time experiments with tagged and polarized Λ (and $\Sigma^\pm$) beams produced by the interaction of kaons with a polarized hydrogen target (Fig. 2.-1).

To constrain either meson-exchange models or quark-gluon models of the hyperon-nucleon interaction, improved two-body scattering data are essential. The question of relating these two approaches still remains open. The question

Fig. 2.-1. Schematic representation of tagged-hyperon experiments with polarized and unpolarized production targets.

of why the physical meson masses are required in the one-boson-exchange potentials remains a challenge for quantum chromodynamics to answer. Furthermore, if NN short-range repulsion is due to Pauli effects at the quark level, should not that repulsion be less in the ΛN (or ΣN) system?

2.2.2.2 Charge Symmetry Breaking. Neutron targets are not very realistic, and Λ targets are even less so. Yet, information on Λn scattering compared to Λp scattering is of fundamental interest. The one-boson-exchange models predict a sizable (compared to nn versus pp) charge-symmetry-breaking difference between the low-energy Λp and Λn strong forces (Pr 80). This charge symmetry breaking results from both electromagnetic Λ - Σ^0 mixing of these $S = -1$ baryons and the mass difference in the $\Sigma^{\pm,0}$ triplet ($m_{\Sigma^+} - m_{\Sigma^-} \approx 10$ MeV), which leads to an asymmetric coupling between $\Lambda p \rightarrow \Sigma^+ n$ and $\Lambda n \rightarrow \Sigma^- p$. Quark models predict different charge-symmetry breaking effects. The $K^- d \rightarrow \Lambda n \gamma$ reaction provides an ideal means of investigating the illusive Λn interaction (Pa 73). The high-energy γ -ray provides a unique signal for tagging this transition. The fact that only two strongly interacting particles exist in the final state ensures that final-state interaction questions, which would plague such studies

as $K^-d \rightarrow \Lambda n \pi^0$, are absent. Enhanced K^- fluxes from LAMPF II will be required to make an accurate experiment of this type feasible.

2.2.2.3 ΣN and ΞN Interactions. The ΣN scattering data are our primary source of information on the strength of $\Sigma \rightarrow \Lambda$ conversion in the strong interaction (ΣN - ΛN coupling). This feature of the ΛN and ΣN interaction plays an important role in the Σ -hypernuclei width problem (Ga 83) and in the anomalous binding of ${}^5_{\Lambda}\text{He}$ (Mi 84). The ΣN spin-orbit interaction is clearly of utmost interest because of the differing predictions for the strength of the term in the asymptotic QCD model (one gluon exchange) and the one-boson-exchange potential model (4/3 versus 0.4 relative to the NN strength) (Do 82B, Po 82). Charged Σ -proton scattering studies are a real possibility using tagged Σ beams from the $p(K^-, \pi)\Sigma$ reaction (Fig. 2.-1).

The ΞN interaction investigation is even more in need of the enhanced kaon fluxes of LAMPF II. The production cross section, from the (K^-, K^+) reaction on a proton target, is much smaller than that for Λ 's or Σ 's. Yet a tagged Ξ^- production experiment in which one studies the rescattering of the Ξ^- in the production target would yield measurements of the spin-averaged and spin-orbit cross sections for $p\Xi^-$ scattering. (The Ξ particles are produced with a high degree of polarization.) Again such experiments are feasible with LAMPF II beams (Ph 83). The one-boson-exchange and asymptotic QCD predictions for the strengths of the spin-orbit term relative to NN scattering are significantly different (0.15 and $-1/3$, respectively).

2.2.2.4 Few-Baryon Systems. Few-body systems offer the possibility of testing the proposed models of the baryon-baryon interaction and of providing direct information about those interactions. The fact that ${}^3_{\Lambda}\text{H}$ ($J = 1/2$) is barely bound and ${}^3_{\Lambda}\text{H}^*$ ($J = 3/2$) is not bound implies that the ΛN interaction in the spin-singlet state is stronger than that in the spin-triplet state. The fact that ${}^4_{\Lambda}\text{He}$ is 0.2 MeV more bound than ${}^4_{\Lambda}\text{H}$ (the additional binding is opposite to what one might have expected from the Coulomb repulsion in ${}^4_{\Lambda}\text{He}$) implies that there exists a significant charge symmetry breaking in the ΛN interaction (Pr 80). The 1 MeV excitation energy of the $A = 4$ spin-flip states provides quantitative information about ΛN - ΣN coupling. (ΛN potentials fitted only to low-energy ΛN scattering data would imply little difference in the binding energies of the $J^{\pi} = 0^+$ and $J^{\pi} = 1^+$ $A = 4$ Λ -hypernuclei.) The anomalously small binding of ${}^5_{\Lambda}\text{He}$ is a significant clue in our understanding of either ΛN - ΣN conversion or in our understanding of quark confinement by gluons and the related question of a Λ - N Pauli exclusion principle (Mi 84): Are quarks so confined in baryons that only identical baryons suffer Pauli effects?

The fore/aft asymmetry measurement in Λ - ${}^4\text{He}$ scattering would provide a direct measure of the Λ -nucleus spin-orbit force. Recall that the strong spin-orbit force in the NN case leads to a resonant $J = 3/2$ state while the $J = 1/2$ state has a much weaker interaction.

Finally, $A = 4$ Σ -hypernuclei studies directly address the question of narrow Σ -widths. Because $\Sigma^- p \rightarrow \Lambda n$ conversion is weak in the singlet state, due to spin-isospin selectivity (Do 82a), narrow $\Sigma^- + {}^3\text{H}$ states may exist if one can substitute the Σ^- for a proton in ${}^4\text{He}$. (See section 2.4.2.)

2.2.2.5 The "H" Dibaryon. One of the most spectacular and intriguing predictions of the MIT bag model is the existence of a stable, flavor-singlet, strangeness -2 dibaryon or "H" particle (Ja 77). Such an "H" particle is expected to have the same quark composition as two Λ 's (uuddss). However, this six-quark object is predicted to have a binding energy such that its mass lies 80 MeV below the $\Lambda\Lambda$ threshold. It could decay only weakly to ΣN or ΛN . Searches for such a massive six-quark object are of paramount importance. If found, the existence of a "di-lambda" would constitute a triumph for bag models and would serve as a severe constraint upon all quark models of hadrons.

Searches for the di-lambda have been proposed using the two-step $K^-p \rightarrow K^+ \Xi^-$, $\Xi^- d \rightarrow \text{"H"} n$ reaction and the single-step $K^-^3\text{He} \rightarrow K^+ n \text{"H"}$ reaction. In either case the counting rate is expected to be only a few per day even for LAMPF II kaon fluxes. A definitive experiment will require LAMPF II unless the "H" is stable and the production cross section is larger than expected.

The concept of a stable "H" is very appealing, because the uuddss quark system could reach optimal symmetry without regard to Pauli effects. However, the observation of double Λ -hypernuclei in old emulsion experiments has been reported in the literature ($^{10}_{\Lambda\Lambda}\text{B}$, $^{11}_{\Lambda\Lambda}\text{Be}$ and $^6_{\Lambda\Lambda}\text{He}$). How could these hypernuclei exist without decaying quickly to more stable dibaryon states? Is the "H" only barely bound with respect to the $\Lambda\Lambda$ threshold, in contrast to bag model predictions? Such puzzling questions remain to be answered as we learn more about quark confinement and the structure of baryons. (The H dibaryon is also discussed in section 2.4.3.2.)

2.2.3 Λ -Hypernuclei

2.2.3.1 Present Status. Hypernuclei were discovered some 30 years ago in emulsion studies using stopped K^- beams. Since then, some 30 species have been identified, and knowledge of their ground-state binding energies has been used to estimate the Λ central potential well depth (some 30 MeV compared to the 50-60 MeV for the nucleon) (Po 78). The advent of magnetic spectrometer systems made possible investigation of hypernuclei via the strangeness-exchange (K^-, π^-) reaction. At a "magic" momentum of approximately 530 MeV/c, recoilless production of lambdas in nuclei can be achieved. Under such conditions, low-spin substitutional states, in which a neutron is replaced by a Λ in the same shell-model orbital, dominate the spectrum (Do 82a). Such experiments led the CERN group to the surprising conclusion that the mean-field spin-orbit potential for a Λ in the nuclear medium is much weaker than that for the nucleon. The Λ spin-orbit strength is less than 2 MeV compared to a value of some 10 MeV in the nucleon case (Ga 83). This difference was subsequently explained in the context of both meson-exchange and asymptotically-free quark models of the Λ -N force. However, these models predict very different sizes for the ΣN and ΞN spin-orbit forces, to which we shall return in section 2.2.4.5.

2.2.3.2 The (K^-, π^-) Reaction. The interpretation of the early (K^-, π^-) data was confirmed by the state-of-the-art experiments on ^{13}C performed at BNL (Do 82a,b). Spectra at nonzero angles showed excitation of states having different spin; data at 0° and 15° are shown in Fig. 2.-2. The relative intensities of the peaks change with angle, and small but significant energy shifts occur which are directly related to properties of the Λ -nucleus interaction. The 10, 16, and 25 MeV peaks at 0° are primarily due to $\Delta L = 0$ transitions which convert $p_N \rightarrow p_\Lambda$; at 15° the 10 MeV peak is $\Delta L = 2$ ($p_N \rightarrow p_\Lambda$),

Fig. 2.-2. Lambda-hypernuclear spectra from an AGS experiment (Do 82B).

whereas the 16 and 25 MeV peaks are combinations of $\Delta L = 1, 2$, and 3 ($p_N \rightarrow s_\Lambda$, $p_N \rightarrow p_\Lambda$, and $p_N \rightarrow d_\Lambda$). Deviations from a weak-coupling picture [coupling a Λ to the 0^+ ($T = 0$) ^{12}C core ground state plus the 2^+ ($T = 0$), 1^+ ($T = 0$), 1^+ ($T = 1$), and 2^+ ($T = 1$) excited states of ^{12}C] provide information about the strength of the quadrupole-quadrupole force as well as the size of the Λ spin-orbit splitting. However, high-resolution data on a variety of p-shell targets will be required before one can sort out the details of the spin-spin and spin-orbit parts of the ΛN force. The large deviation of the strengths of the dominant peaks, from predictions based on neutron-pickup strengths from $^{12}\text{C}(d,p)^{13}\text{C}$, indicates that hypernuclei tend to form states with a higher degree of spatial symmetry than is possible in normal nuclei. If one uses the states with [54] and [441] symmetry as a basis, the [54] symmetry possible in $^{13}_\Lambda\text{C}$ is forbidden in a system of 13 nucleons by the Pauli principle. This is analogous to the s-shell nature of $^5_\Lambda\text{He}$ contrasted with the need for one of the 5 nucleons in the ^5He system to be in a p-shell (or higher) state. Thus, the evidence for a dynamical selection rule emerges.

Full exploitation of the structure information available from the spectra of Λ -hypernuclei requires considerable improvement in energy resolution, which is possible only with more intense K^- beams. A preliminary design shows that one could obtain a beam flux 100 times greater than the 10^5 K^-/s available today with an energy resolution 10 times better than the 3 MeV presently attainable and with a $\pi:K$ ratio of 1:1 (Ph 83). The impact of these improvements can be judged from Fig. 2.-3 which shows pion-scattering spectra prior to and during the meson factory era. Thus, one can hope to achieve experimental conditions necessary to characterize the excited states of the accessible hypernuclei and

Fig. 2.-3. Comparison of inelastic pion-scattering spectra prior to and during the meson factory era.

properly determine those aspects of the ΛN force (for example, the spin-spin and tensor components) not yet explored.

We note that study of mirror hypernuclear systems will require a high-resolution π^0 spectrometer (Ph 83). The (K^-, π^0) reaction has not yet been exploited for this reason. This reaction will selectively populate the $T + 1/2$ hypernuclear states which are otherwise hidden under the $T - 1/2$ states in the (K^-, π^-) reaction.

2.2.3.3 High-Spin States. The (K^-, π^-) reaction primarily populates low-spin, substitutional states, unless one moves to larger angles where the production cross section falls off. A complementary production process, the (π^+, K^+) reaction, can excite the higher spin states at forward angles, because the reaction involves a larger momentum transfer (Ph 83). Although the cross section is smaller than the (K^-, π^-) reaction for the same momentum transfer, the π^+ beam flux can be 10 to 30 times larger than that of the K^- . The momentum transfer in the (π^+, K^+) reaction can be chosen to match the momentum structure of the form factor for high-spin, natural-parity "stretched" states. These are the analogs of the most interesting high-spin states seen in normal nuclei and, for similar reasons, should be some of the most important excitations explored in hypernuclear physics.

2.2.3.4 Hypernuclear γ Spectroscopy. The $(K, \pi \gamma)$ coincidence measurements provide superb energy resolution in the study of particle-stable hypernuclear states. The technique is being pushed to its present limit at BNL (Ma 84) in an effort to extract the size of the effective spin-spin interaction in the shell-model description of Λ -hypernuclear structure (Do 82b, Mi 84). Such coincidence measurements provided the excitation energies of the M1 spin-flip states in the $A = 4$ system. Spin and parity information for such particle-stable states, as well as unique values for the spin-spin, symmetric and antisymmetric spin-orbit, tensor, and three-body components of the Λ interaction, can be obtained from similar high-resolution measurements on heavier targets.

As an example of the utility of such measurements, consider one aspect of ${}^7_\Lambda\text{Li}$ compared to normal ${}^6\text{Li}$ (Fig. 2.-4) (Do 82b, Gi 82). Observation of the $5/2^+$ to $1/2^+$ transition has confirmed that the low-lying 3^+ (2.18 MeV) continuum state of the ${}^6\text{Li}$ system can be stabilized with the addition of a Λ . This ${}^6\text{Li}$ core state is difficult to model because there is no bound ${}^5\text{He}$ or ${}^5\text{Li}$ state from which it can be formed by the addition of a nucleon. It is not well represented as a hole in ${}^7\text{Li}$. Thus, isotope and isotone studies do not provide realistic tests for nuclear models of ${}^6\text{Li}$. Because the methods of treating continuum states differ from those for bound-states, one has been limited to comparison with ground state properties of ${}^6\text{Li}$. However, the addition of a Λ to form particle-stable states in ${}^7_\Lambda\text{Li}$ can now be used to test models of ${}^6\text{Li}$. With the insertion of the relatively weakly interacting Λ , the ${}^6\text{Li}$ core remains intact while the 3^+ and 0^+ continuum levels become particle stable.

Fig. 2.-4. Level diagrams of ${}^6\text{Li}$ and ${}^7_\Lambda\text{Li}$.

description of ${}^6\text{Li}$ can now be evaluated in terms of the dynamics of a system with several bound levels. Success in describing the spectrum of ${}^7\text{Li}$ depends crucially upon the correct modeling of ${}^6\text{Li}$. Gamma-rays other than the observed 2 MeV $5/2^+ \rightarrow 1/2^+$ transition are needed to take full advantage of the power of this technique.

2.2.3.5 Nonmesonic Weak Λ Decay. In the early studies of hypernuclei it was realized that Pauli blocking in a heavy system would inhibit the $\Lambda \rightarrow \pi N$ free-space decay mechanism. Therefore, the dominant decay mode of these hypernuclei should be the nonmesonic decay corresponding to the reaction $\Lambda N \rightarrow NN$. Non-leptonic weak vertices such as this are also studied in nuclear parity-violation experiments. For this reason, one hopes to investigate the structure of the four-Fermion, nonleptonic, strangeness-changing weak interaction through studies of the decay of Λ -hypernuclei with mass number greater than 5 (Mc 84). Indeed, such studies may be the only means of examining the weak $\Lambda N p$ coupling. Existing meson-exchange model calculations of the nonmesonic decay rate, Γ_{nm} , show a sensitivity to this factor.

Furthermore, there exist calculations of $\Gamma_{\text{nm}}/\Gamma_{\text{free}}$ in both a meson-exchange model and a quark model. The value of $\Gamma_{\text{nm}}/\Gamma_{\text{free}} \approx 1$ from the meson exchange models is three times larger than the $\Gamma_{\text{nm}}/\Gamma_{\text{free}} \approx 0.3$ prediction of the quark model (Ch 83). (A hybrid model calculation involving both pions and quarks obtains a value of 2-4.) The only measurement ($A = 16$) yielded a value of 3 ± 1 for 22 events. A new BNL measurement in ${}^{12}_{\Lambda}\text{C}$ will push the technological limits of the present machines. Systematics will be required, however, before definitive conclusions can be drawn about either the relevance of any of the proposed models or the properties of the four-Fermion weak amplitude.

2.2.4 Beyond the Λ

2.2.4.1 Present Status. Although only a few Σ -hypernuclei have been observed, Σ -hypernuclear studies lie at the forefront of this field of investigation. Forward production of Σ 's has been studied in p-shell targets from Li to C, using the (K^-, π^+) reactions at 720 MeV/c. [Note that the (K^-, π^+) reaction is specific for forming Σ^- -hypernuclear states from valence protons.] Surprisingly narrow Σ -states were observed (Po 82, Ga 83, Be 84). A large width due to strong $\Sigma \rightarrow \Lambda$ conversion had been anticipated. These presumably correspond to coherent substitutional transitions leading to 0^+ final states, which would dominate near the Σ "magic" momentum of approximately 300 MeV/c.

Interesting questions arise in the interpretation of these data. Why are some Σ -states relatively narrow? What are the single-particle properties (well depth, spin-orbit potential, etc.) of a Σ in a nucleus? Do Σ states exhibit good isospin? The data are yet too crude to permit definitive answers. Angular distributions for (K, π) reactions are needed, and an intense low-momentum K^- beam is a major component of any such experiment. Without major improvements in present beams one can only look longingly at the tantalizing hints which have come from the rudimentary, exploratory measurements.

2.2.4.2 Σ -Hypernuclei Narrow Widths. Why do some Σ -hypernuclei exhibit narrow states? Have we seen only those states which are narrow? Do narrow states occur only in light systems, or are they representative of a rich, but yet undiscovered, Σ -hypernuclei spectroscopy of heavier systems?

Several alternative explanations exist for why Σ -states should have widths substantially smaller than the semiclassical estimate of 20-30 MeV for nuclear-matter densities (Ga 83). Unstable bound states is one suggestion. Ordinary Σ resonance states in the continuum, which can occur in the absence of $\Sigma \rightarrow \Lambda$ conversion, are described by nonnormalized functions corresponding to poles in the fourth quadrant of the momentum plane. Including absorption leads to possible additional poles in the second quadrant, which correspond to normalizable "bound states in the continuum" eigenstates. For such states, the $\Sigma N \rightarrow \Lambda N$ transition width and the Σ escape width can partially cancel leading to Σ -states which are narrower than a conventional resonance picture would indicate. Because such unstable bound states are coupled only weakly to the usual continuum, they would not be seen in Σ -nucleus scattering but would show up strongly in (K^-, π^-) studies. The narrow states in such a picture would be limited to those whose orbits place them near the surface, i.e. narrow Σ -states are not universal in this picture.

A second explanation involves the spin-isospin selectivity of the $\Sigma \rightarrow \Lambda$ conversion process. The low-energy $\Sigma^- p \rightarrow \Lambda n$ reaction is dominated by the $T = 1/2$, 3S_1 ΣN channel. (The $T = 1/2$ is due to the Λn final state; the $S = 1$ may be the result of a one-pion-exchange mechanism.) Because the nuclear cores of most hypernuclei are not spin-isospin saturated, suppression of the $\Sigma N \rightarrow \Lambda N$ transition, leading to selectively narrow states, may be very A dependent. In such a picture, this suppression would be most pronounced in s -shell and p -shell hypernuclei (for heavier systems, a larger fraction of the core nucleons is in spin-isospin saturated subshells). Thus, the narrowest Σ -widths would occur in the lightest Σ -hypernuclei, particularly those having maximum spin.

Finally, one might have narrow Σ -states due to Pauli and binding effects, short range of the $\Sigma N \rightarrow \Lambda N$ transition interaction, or ΣN initial-state correlations. Each of these explanations suggests more universally narrow Σ -states. This quenching of the widths would be most interesting and imply that, indeed, a rich Σ -hypernuclear spectroscopy awaits our search.

2.2.4.3 Metastable Σ -Hypernuclei. The spin-isospin selectivity explanation of narrow Σ -widths leads one to ask whether one might not also observe metastable Σ -hypernuclei. Such bound systems, not the trivially stable $\Sigma^- nn$ and $\Sigma^- \Sigma^- nn$ $T_z = T$ states, might result from the $S = 1$ dominance of the $\Sigma N \rightarrow \Lambda N$ transition reaction (Do 82a). For example, if a Σ^- were inserted in ^4He in place of one proton such that the remaining Σp pair had spin zero, then suppression of the $\Sigma N \rightarrow \Lambda N$ conversion would be maximal. If the $\Sigma N \rightarrow \Lambda N$ transition potential for the 1S_0 were zero, then this coherent substitutional state would be strictly bound. Whether such metastable states are found to exist, the Σ -decay widths of such systems should be very narrow.

2.2.4.4 Isospin Mixing in Σ -Hypernuclei. The question of isospin mixing in Σ -hypernuclei is open (Ga 83). If the size of the Lane potential ($t_\Sigma \cdot T$) in the Σ -nucleus interaction is somewhat larger than the charge-dependent factors (such as $\Sigma^{\pm,0}$ mass differences and Σ Coulomb energies), then Σ -hypernuclei will be best described in an isospin basis as are normal nuclei. Present data do not allow one to make a definite conclusion. Thus, the question of the size of the symmetry potential in Σ -hypernuclei remains unanswered until more precise data become available. $SU(3)$ model estimates indicate that the Σ Lane potential

could be as large as half of the nucleon Lane potential, so that one might anticipate sizable isospin splittings in Σ -hypernuclei which have $N \neq Z$.

2.2.4.5 ΣN and ΞN Spin-Orbit Forces. The surprisingly small Λ spin-orbit potential has received much attention because a simple additive quark model, in which the baryons can exchange only light quarks, predicts the ΛN spin-orbit potential to be identically zero (Po 82). The one-boson-exchange potential model and the mean-field approach to baryon-baryon scattering have also been shown to yield small (nonzero) values for the ΛN spin-orbit interaction compared to that of the NN system.

Because the asymptotic QCD quark-quark interaction has a strong spin-spin component, and because the nucleon-nucleon spin-orbit interaction is of shorter range than the central and tensor components, it has been suggested that the investigation of hypernuclear spin-orbit potentials might show evidence for quark-gluon degrees of freedom in nuclei (Do 82b).

The additive quark model yields for the spin-orbit strengths (Po 82, Do 82b)

$$v^N : v^\Lambda : v^\Sigma : v^\Xi = 1 : 0 : 4/3 : -1/3 \quad .$$

In contrast, one of the one-boson-exchange models (which has been fit only to NN , ΛN , and ΣN data) predicts a ratio of strengths given by (Do 82b)

$$1 : 0.2 : 0.4 : 0.15 \quad .$$

The mean-field approach (which, in certain extreme limits dominated by ω exchange, can be made to yield results identical to the additive quark model) yields (Do 82b)

$$1 : 0.2 : 0.5 : 0.05 \quad .$$

LAMPF II will be able to provide vastly improved data from which the ΣN spin-orbit interaction can be extracted. It will also make feasible the first experiments on Ξ -hypernuclei and ΞN scattering. These will be crucial in distinguishing among the above models of spin-dependent forces in hyperon-nucleon systems.

2.2.4.6 Y^* Resonances. Propagation of the Δ resonance in the nuclear medium has been a subject of much interest (Al 79). The shift in the pole position as a function of target mass, the change in the width due to spreading

interactions, collision damping, and dispersion effects, and the coupling of the resonance to open channels such as $\Delta N \rightarrow NN$ have been keenly studied in order to understand the dynamics of this unstable state which dominates the πN spectrum. The Δ plays an important role in the observed large ratio between the $T = 0$ and $T = 1$ pion absorption cross section and in the systematic reduction of the Gamow-Teller and magnetic-transition strengths.

Investigation of the Δ resonance and how it propagates in the nuclear medium has been inhibited by the natural width (120 MeV) of the Δ . Its probability for decay within the target nucleus after being produced on a given nucleon is just as large as its probability to interact with a second nucleon. Because of this problem, the Λ^* (1520) is considered to be a better candidate with which to investigate resonance propagation in nuclear matter (Al 79). It is well isolated from the nearby resonances at 1405 MeV and 1670 MeV. The narrow width, 16 MeV, ensures that it will travel much farther before decay. However, one must avoid overwhelming the resonance with Fermi broadening, which will occur in K^- elastic scattering. Use of a (K^-, π^-) tag will avoid the problem.

If the $\Delta^* N$ interaction proves to be similar in nature to that of the Λn , then one would expect to see Δ^* -hypernuclei. Because the Δ^* is a $3/2^-$ state, the physics of such systems would be a first. We have no other way to study negative parity baryons in a nuclear medium. The narrow width of the Λ^* (1520) would ensure their observation.

Radiative decays of the γ^* resonances will address the issue of the quark structure of hyperons (Ph 83). There exist several models of the structure of the baryons: the simple three-quark model (q^3), the SU(6) basis model, the $q^3 + q^4 \bar{q}$ basis, etc. In the three-quark picture, the radiative widths for such transitions as Λ^* (1405) $\rightarrow \Sigma^0 (1193) \gamma$ and Λ^* (1520) $\rightarrow \Sigma^* (1385) \gamma$ are strictly forbidden. Quantitatively different predictions for the various models have been made for Λ^* (1520) $\rightarrow \Sigma \gamma$ and $\Lambda \gamma$, for example. Therefore, the low-lying excited states of the $S = -1$ and -2 baryons provide a useful laboratory in which to study the question of confinement.

2.2.4.7 $S = -2$ Hypernuclei. The (K^-, K^+) double-strangeness-exchange reaction opens the door to the study of doubly-strange ($S = -2$) hypernuclei (Ξ , $\Lambda\Lambda$, $\Lambda\Sigma$, $\Sigma\Sigma$). Investigation of the spectroscopy of $S = -2$ hypernuclei represents a logical next step in the evolution of hypernuclear physics (Do 82a). Narrow Ξ states are likely to exist, just as in the case of the Σ , although the widths will depend upon the $\Xi^- p \rightarrow \Lambda\Lambda$ conversion rate. ($\Xi N \rightarrow \Lambda\Lambda$ should be smaller than $\Sigma N \rightarrow \Lambda N$, making conversion less of a problem.) Our only hope of investigating the $\Lambda\Lambda$ interactions is through study of the spectroscopy of $\Lambda\Lambda$ -hypernuclei. Such information is necessary if we are to extend our knowledge of the SU(3) structure of the baryon-baryon forces. It would provide further tests of the one-boson-exchange potential model, which is fitted to NN, ΛN , and ΣN scattering data, but makes specific predictions about the $\Lambda\Lambda$ and ΞN interactions.

One-boson-exchange considerations indicate that none of the possible few-body systems are bound: $\Xi^- n$, $\Xi^- nn$, $\Lambda\Lambda n$ or $\Lambda\Lambda nn$. However, reactions on many p-shell targets should lead to $\Lambda\Lambda$ -hypernuclear bound states; an example would be the reaction $^{12}\text{C}(K^-, K^+)_{\Lambda\Lambda}^{12}\text{Be}$. Cross sections for production of specific discrete $\Lambda\Lambda$ -hypernuclear states are expected to be small. Observation of them will

provide a formidable challenge to experimentalists even at LAMPF II; at existing facilities such experiments are not feasible.

2.2.4.8 Charmed Hypernuclei. The production of a charmed state of a bound many-baryon system is now within the realm of possibility. Simple SU(4) extensions of the one-boson-exchange potential models (that describe NN and YN scattering) to include C^+N scattering ($C^+ \sim cud$) have been used in model calculations of few-body systems (Gi 83). They predict the existence of ${}^4_C\text{He}(C^+nnp)$ and ${}^5_C\text{Li}(C^+nnpp)$. These bound states are estimated to have binding energies similar to ${}^4_\Lambda\text{H}$. In addition, bound states have also been predicted for heavier systems, just as they exist for heavy Λ -hypernuclei. Because the charmed baryons decay rapidly, investigation of such C^+ -hypernuclei are our only means of investigating even crudely the C^+N interaction.

2.2.4.9 Hyponuclei. No baryons with $S = +1$ have been observed. Nevertheless, there exists the possibility that nuclear matter with positive strangeness can be produced (Go 82). This hyponuclear matter would be formed by introducing an $\bar{s}$ quark into a nucleus and making it "stick." Five-quark objects ($q^4\bar{q}$) containing an $\bar{s}$ have been postulated, but no such Z^* resonances have been observed in K^+N studies. However, five-quark configuration-mixing has been employed in descriptions of strange-baryon resonances, and there is no fundamental inhibition against $q^4\bar{q}$ objects in nuclear matter.

Both (π^+, Λ) and (K^+, p) reactions would deposit a stopped or slow K^+ in the nucleus. The question is open as to whether the K^+ can form bound or resonant states, i.e., a kaonic nucleus with $S = +1$. The nuclear medium could modify the K^+N interaction to such an extent that a net attractive K^+ -nucleus interaction might emerge. If such a state were to exist, it could decay only weakly because all strong decay channels would be closed.

2.3. MEDIUM- AND LARGE-LENGTH SCALE PHYSICS

2.3.1 Introduction

A language which has proved useful in describing phenomena at the medium-length scale (0.8-2.0 fm) is that of meson and baryon degrees of freedom. Impressive theories of the nucleon-nucleon interaction have been based on exchanges of π , 2π , η , ρ , ω mesons, and intermediate excitations of the Δ_{33} resonance (Ma 84, Vi 79). Current-generation meson factories have provided an abundant source of the least massive of the mesons, the π meson, and one has been able to examine in detail the nature of the propagation and interaction of the pion and Δ_{33} resonance within a collection of nucleons. Theoretical and experimental studies have made it possible to assess the role of pions and the Δ_{33} resonance in the equation-of-state of dense neutron-star matter, the spin-isospin collective response of nuclei, and the quenching of the Gamow-Teller resonance.

Although mesons and baryons may be the dominant degrees of freedom in the medium-energy/momentum transfer regime of nuclear reactions, certain experimental studies could provide key evidence about the role played by quarks and gluons in the structure of nuclei. Two such possibilities are examined in subsection 2.3.2. A neutrino-nucleus scattering experiment is sensitive to an enhanced nuclear pion field, proposed recently to explain nuclear quark distributions from ultra-high-energy lepton scattering experiments. Pion double-charge exchange and reactions producing higher mass mesons, such as η , η' , ρ , ω , and ϕ mesons, may provide extreme conditions where the usual (non-quark) models of nuclear reactions are inadequate.

The programs discussed in subsection 2.3.3 rely on high-intensity K^+ , antiproton, and higher-energy π^+ beams to extend the microscopic frontiers of nuclear structure and reactions. What can antiprotons, for example, tell us about the emerging new models of nucleon- and antinucleon-nucleus interactions based on the Dirac equation? Intense beams of K^+ 's, the "hadronic electron," will be sensitive to the distribution of neutrons in the nuclear interior. Pion scattering and charge exchange at higher resonances will produce views of nuclear structure which are complementary to medium-energy proton and electron scattering experiments.

2.3.2 Mesonic and Baryonic Structure of Nuclei

Studies of mesons and their interactions with hadrons that can be made at LAMPF II are essential to progress in unraveling the dynamics of nuclei. One must understand how mesons and baryons behave in nuclei before the more subtle effects of quarks and gluons can be isolated. Experiments in this regime will have applications in the quantitative description of medium modifications to weak and electromagnetic interactions in nuclei. For example, meson-exchange currents play a decisive role in the extraction of nuclear-structure information from data to be obtained at electron accelerators. The meson-exchange corrections have important implications for strong nuclear reactions; they are a source of medium-range 3-body interactions affecting knock-out reactions and optical-model behavior, as well as properties of the nuclear ground state. Meson theory underlies, at an ever-increasing depth of application, nuclear models common to studies at all medium-energy facilities.

2.3.2.1 Charged-Current Neutrino Reactions: Enhanced Pion Field in Nuclei?

The anomalous A dependence of deep-inelastic muon scattering (Au 82), now known as the European Muon Collaboration (EMC) effect (see section 2.4.2 for a detailed discussion) implies that nucleon-quark distributions are different when the nucleons are embedded in nuclei. The EMC effect was discovered in an experiment that probed nuclear structure at very short wavelength. However, critical tests of many of the ideas advanced to explain this effect rest on observations of coherence in quark wave functions which are present only for longer-wavelength probes.

A prime example is a test of the enhanced-pion-field model of the EMC effect (Er 83). Enhancements in the pion field of a heavy nucleus, long anticipated on theoretical grounds, provide a natural mechanism for the increased scattering cross section seen in the low-momentum transfer region in the EMC experiment. At the wavelength of 200-GeV muons there is of course no hint that the coherence implied by the enhanced-pion model is present; pions are merely the source of additional quark-antiquark pairs. This is where the copious quantities of ~ 1 -GeV neutrinos produced at LAMPF II will play a unique role.

Charged-current neutrino reactions offer the following advantages in probing the nuclear pion field:

- (1) Unlike the electromagnetic interaction, the weak interaction couples directly to the pion field via the pion pole term.
- (2) The wavelength probed by 1-GeV neutrino interactions is appropriate for "seeing" the pion as opposed to an incoherent $q\bar{q}$ pair at high energies.
- (3) The entire nuclear volume is probed, unlike recent hadron-scattering measurements of the pion field (Ca 84a).
- (4) Statistical precision, even at the level of the EMC experiment, will be attainable because the cross sections are large by neutrino physics standards.

An enhanced nuclear pion field will manifest itself as an anomalous A dependence in the charged-current production cross section in the momentum transfer range $\sim 2 \text{ fm}^{-1}$ for nuclear-excitation energies of $\sim 70 \text{ MeV}$. The energy resolution obtainable in several proposed detection schemes is more than sufficient for this application.

2.3.2.2 Meson-Baryon Interactions in Nuclei. Studies of meson propagation in the nucleus have been previously made in high-energy pion and electron scattering. As a result of the measurements, one has learned about the interaction of the ρ -meson (Vo 81) and other composite 3π and 5π systems with the nucleus (Be 71). Figure 2.-5 shows, for example, an extraction of the ρ -nucleon cross section obtained in a coherent-production experiment with π^+ mesons of momentum 2.9 GeV/c, compared to measurements from coherent photoproduction. Recent measurements of η production in nuclei have been made at LAMPF I. One of the interesting and poorly understood findings from the higher-energy measurements at the CERN PS is that the mean free path of multipion systems is longer than naively expected (Be 71). This has led to speculations about the length scales over which hadronic matter hadronizes.

Fig. 2.-5. The rho-meson-nucleon total cross section as deduced from π^+ meson-nucleus scattering and coherent photoproduction (Co 84).

With the higher intensities of pions and K^- available at LAMPF II, it will be feasible to extend the knowledge about the interaction of heavier mesons, multimeson and excited baryon systems with the nucleus. For example, in pion- or proton-induced reactions, one may produce heavy mesons and identify them by detecting their decay products or the recoiling nucleus (Co 84). Detection of two-body decay modes would provide a clean signal, and even relatively small branching ratios to charged particles could be utilized (e.g., the 1% branching ratio of $\omega \rightarrow \pi^+\pi^-$). Antiproton-nucleus collisions are another means for producing mesons in nuclei, and the peak flux ($10^9 \bar{p}/\text{sec}$ at 3 GeV/c) is perfectly matched to production of the charmonium family (Gi 83). Antiproton-nucleus collisions in the energy range 8-10 GeV give rise to pions in the range of 0-1.5 GeV and are well suited to study the hadronization time scales.

3.2.3 Double Charge Exchange at High Momentum Transfer

Pion-induced double charge exchange (DCX) is a unique reaction because two nucleons must be involved in order for two units of charge to be transferred to the nucleus. With the flux available at LAMPF I, these cross sections are measurable for a few states of the nucleus, most notably the double isobaric-analog state. The analyses of these data show that large contributions appear to arise from short-range phenomena (Gr 84, Le 84). Possible explanations include exchange of heavy (vector) mesons and nucleon excitations. Various contributions to the DCX process are shown in Fig. 2.-6.

Fig. 2.-6. Examples of meson and baryon contributions to pion double charge exchange: (a) Sequential scattering through N^* and Δ resonances (A and B) as mediated by meson exchange (C); (b) DCX from the pion cloud in the nucleus; and (c) a term closely related to (b).

A more economical description may emerge from the quark-bag picture in which the distinct heavy-meson and nucleonic excitations are replaced by six-quark configurations in a region of space where the nucleons overlap in the nucleus. One may see more clearly the possibilities of this picture in Fig. 2.-7 which shows a slice through the center of the nucleus under the assumption that nucleons have a radius of 0.7 fm. The proper radius is not well determined and is a matter of great interest. From this simple picture, it is clear that a substantial fraction of the nucleons ($\approx 30\%$) are overlapping at central nuclear densities, and the nucleus thus provides a laboratory for studying extended quark systems. Recent hybrid quark-hadron model estimates (Mi 84, Ki 84) suggest that even at LAMPF I energies a sizable part of pion DCX take place in the six quark region. The DCX reaction is thus seen as a nuclear process which may involve new elements of both the hadronic and the multiquark structure of nuclei.

By measuring the two outgoing protons one can take a "snapshot" of the nucleus and study the DCX process as a function of the distance between the two nucleons. The important quantity to measure is the cross section as a function of q , the momentum transferred between the nucleons (in Fig. 2.-6a,

Fig. 2.-7 Cross section of a ^{207}Pb nucleus. The blue circles represent neutrons and the red, protons. The nucleus and nucleon radii were assumed to be 7.1 and 0.7 fm, respectively. Nucleons are assumed to be randomly distributed.

$\vec{q} = \vec{p}^- + \vec{k} - \vec{p}$), and to study the dynamics for different quark configurations. By an appropriate choice of kinematics of the incident pion and the outgoing nucleons one has the freedom to select states A, B and C (Fig. 2.-6). The onset of deconfinement may be identifiable as $|\vec{q}|$ is raised.

2.3.3 Large-Scale Nuclear Phenomena

Characterizing the structure of the nuclear ground and excited states is one of the most challenging aspects of the field of nuclear physics. Interplay between experimental programs and the developments of nuclear many-body theory has always been a major source of intellectual vitality for the field. LAMPF II will be able to participate directly in this process by providing beams of particles sensitive to nuclear properties that have been difficult to measure at more traditional facilities. A necessary outgrowth of nuclear structure physics at LAMPF II will be a deeper understanding of all aspects of nuclear interactions and dynamics.

The variety of spin and isospin selectivity of p , $\bar{p}$, K^+ , e , ν and $\pi^\pm$ make it possible to probe the spin and isospin properties of the nucleus in new and complementary ways. For example, calculations suggest that antiprotons would be selective probes of isoscalar spin-flip excitations (Do 83), the least understood of the elementary modes of excitation. One expects that the decreasing de Broglie wavelength of the projectile, the weakening of the projectile-nucleon cross sections and the relatively minimal influence of Coulomb scattering for the higher energy range proposed here, will permit transparent connections between nuclear structure and the basic projectile-nucleon interaction.

3.3.1 Antiproton-Nucleus Scattering

The use of the antiproton as a nuclear probe opens a wide range of research spanning such diverse fields as reaction theory, relativistic scattering theories, the fundamentals of antinucleon-nucleon and nucleon-nucleon interactions, and nuclear structure. The limitations of past utilizations of $\bar{p}$'s can be traced directly to the conditions of low incident flux and poor beam purity. A revolution came with the realization that cooled antiproton beams could be created with superb beam purity and phase space. This led to the LEAR project at CERN. The first generation of approved LEAR experiments is under way, and there appear to be many more proposals than the facility can accommodate.

LAMPF II will be a copious source of antiprotons with clear advantages over the present LEAR facility in intensity and in energies above 1.2 GeV. The characteristics of the present $\bar{p}$ beams at the AGS and LEAR and of the proposed LAMPF II are given in Table 2.-I.

Experiments designed to explore issues in nuclear reactions and structure will probably concentrate in the 0.5 to 2 GeV region. Initially, these experiments would use either separated beams with fluxes typically in the 10^7 $\bar{p}$ /sec range or unseparated beams in the 10^9 $\bar{p}$ /sec range. Many elastic, inelastic, or charge-exchange experiments could use the latter since antinucleon identification in the final reaction state would identify the beam particle as an antiproton. Subsequent additions to the facility would certainly include a

TABLE 2.-I

Comparison of Facilities

	LEAR	AGS	LAMPF II
E_{max} (GeV)	1.2	~20	~30
$\bar{p}$ Flux (per sec)	10^6	10^7	10^9
Separated Flux	10^6	10^5	10^{9*}
Momentum Spread	0.05%	2%	2%

* Assumes time-separated beam.

time-separated beam for high-flux experiments ($10^9 \bar{p}/\text{sec}$) and cooled beams for precision low-energy measurements.

At the heart of nuclear physics is the issue of how to solve the many-body problem. Experience with nucleonic probes has led to successful theoretical descriptions of elastic scattering and reactions such as (p,p') , (p,n) , (d,p) , etc. The antiproton behaves in many ways like a strongly absorbed nucleon, but it is fundamentally different. Do our ideas of reaction theory carry over to the extreme conditions that can be created with antinucleons? Can we, by going to this more exotic particle, force what are normally small corrections under "normal conditions" to become dominant effects in different regimes?

The mechanism of antiproton-proton annihilation in free space or the nucleus is not at all understood and is of great importance. In free space, the average pion multiplicity is five. Does this change when the proton annihilates in the nuclear medium, where the possibility exists for the antiproton to annihilate on several nucleons at once? Antiprotons of energy greater than 1 GeV have mean free paths of several Fermis and penetrate into the high-density region of nuclei. One could, for example, look for evidence of multiple nucleon annihilation by changes in the mean pion multiplicity.

Elastic and inelastic scattering of $\bar{p}$ from nuclei has considerable interest as a means of studying the response of the nucleus to an external field, using the selectivity of nuclear excitations to study particular components of the elementary interaction, and determining whether or not the interaction is different in the nuclear medium than it is in free space. Determination of the optical potential for antinucleons, which describes the single-particle behavior of the antiproton in the nuclear medium, is especially interesting in the context of the relativistic Dirac theory of nucleon-nucleus interactions (Se 84), where large couplings to virtual NN pairs seem to be important. The spin structure of the interaction for inelastic scattering is particularly important since qualitative relations between different spin-isospin components are often the clearest predictions of various quark or meson-exchange models.

The zero- (or negative-) energy $\bar{N}$ -nucleus interaction has an interest all its own. Do strongly bound states exist in any sense? How do low-energy antinucleons propagate through matter? This interaction is revealed through the use of $\bar{p}$ atoms and/or $(\bar{p}, p)$ reactions leading directly to atomic states or the putative deeply-bound states. Investigations in these energy regimes deal with the negative-energy interaction of the $\bar{N}N$ system. LAMPF II will enable a basic test of these considerations by utilizing an observation recently made by Miller (Mi 84). He has emphasized that the large scalar and vector interactions of the emerging relativistic Dirac theory give rise to a strong binding of the nucleons occupying the Dirac negative-energy hole states in nuclei. This lowers by 0.6-1 GeV the threshold energy at which knock-out reactions remove nucleons from the occupied states of the negative-energy sea. Because of the large binding, the $\bar{N}$ would annihilate rather than leave the nucleus. The experimental signature of such processes (a large multiplicity of pions coming from the annihilation) should be observable in coincidence measurements and would lead to a critical test of the Dirac-based mean-field theories.

2.3.3.2 The K^+ -Nucleus Interaction as a Probe of Nuclear Wave Functions.

Because of low K^+ fluxes, experiments with these particles are difficult with current facilities. LAMPF II will provide as much as two orders of magnitude more K^+ to experimental areas and therefore make a variety of new experiments feasible.

Among particle probes, the K^+ holds a very special position. Because the antiquark in its constitution is strange and since there are no (valence) strange quarks in the nucleon, quark-quark annihilation is suppressed. Thus the system cannot couple strongly to a 3-quark object. The K^+ -nucleon system is a true 5-quark quantity and there are no presently known resonances below 1 GeV/c. For this reason the K^+ -nucleon interaction has two important properties: it has a slow energy dependence and it is very weak--the weakest of all known strong interactions.

The implications for K^+ -nuclear scattering are significant (Do 82b). Because of the weakness of this interaction, the distortion effects are small, so an approximate treatment of them is of greater utility than for pion or nucleon projectiles. The lack of a strong energy dependence is very important for theories treating the scattering and distortion effects, because the direct use of measured phase shifts is possible. In contrast, for pions near the (3,3) resonance, the binding of the nucleons introduces large corrections. Since the principal partial wave is s (not p wave as for the pions), nucleon-recoil effects are smaller. Additionally, the smaller amount of multiple scattering causes the finite range of the interaction to play a small role.

Another result of the weakness of the K^+ coupling is that it penetrates deeply into the nucleus. Thus, the whole nuclear volume is sampled, not just the surface, as is the case for resonant pions. This allows one to obtain a more complete picture of the nucleus. The comparison of K^+ scattering to electron scattering provides a means for studying the neutron distribution in nuclei.

Valence neutron distributions can also be studied with the charge-exchange reaction (K^+ , K^0). The distortion effects in this reaction are much smaller than the corresponding pionic case so that nuclear structure information can be extracted with greater confidence.

In summary, the K^+ -nucleus interaction can be constructed from the K^+ -nucleon interaction with much greater reliability than any other strongly interacting particle. All of this assumes, of course, that the relevant degrees of freedom of the nucleus are entirely nucleonic. We know that this picture is oversimplified and requires some modification: the structure of the K^+ and the target nucleons must enter at some level. The meson cloud surrounding the nucleon is one source of these corrections. At a deeper level, it becomes of great interest to understand how the dynamics of the scattering are influenced by the structure of the hadrons themselves. There are preliminary indications that such modifications are important (Si 84), and that it may be more economical to view K^+ scattering as probing the overall quark distribution.

2.3.3.3 Elastic and Inelastic Scattering of Pions. The conceptual design for the LAMPF II experimental areas permits beam lines that view the production target at 0° . The high intensity of pions produced by a 45-GeV beam along with the small pion source size obtained by observing the production target at 0° should make possible a much higher intensity facility with better resolution than is presently possible at meson factories. Typically, increases of 10 can be expected for π^+ fluxes and 100 for π^- fluxes. These increased fluxes, coupled with an expected factor of two improvement in energy resolution, present a unique opportunity for exploring nuclear substructure as well as extending conventional nuclear structure studies with pions.

An important property needed in a probe for studies of nuclei is differential sensitivity to neutron and proton components in the excitation of nuclear states. A good example is ^{14}C , shown in Fig. 2.-8. Knowledge of these separate components provides stringent tests of our microscopic description of nuclear structure. Also, this differential sensitivity allows identification of simple modes such as pure neutron or proton excitations, the so-called "stretched" states, or isovector collective states. It also makes feasible

Fig. 2.-8. Spectra from inelastic π^+ and π^- scattering from ^{14}C . Figure from Ref. (Ho 84).

searches for regions of isospin mixing in the excitation spectrum. The large asymmetries in the π^+ and π^- cross sections for excitation of the 2_2^+ state and three 4^- states of ^{14}C are promising sources of information for determination of the proton and neutron transition amplitudes.

The ratio of the π^+p/π^+n -interaction strength (Fig. 2.-9) is seen to be a sharply energy-dependent function. This is true both for the central component of the force that excites natural parity states, such as giant electric resonances, and the spin-orbit component that excites unnatural-parity states. From Fig. 2.-9 one can see that the cross-section ratio is larger than 1 at $T_\pi \lesssim 300$ MeV, drops well below 1 for $300 \lesssim T_\pi \lesssim 1000$ MeV, and rises to ≥ 1 at higher energies. Thus the important quantity to measure is the π^+/π^- cross-section ratio at the peak of the angular distributions. This ratio, as a function of pion energy, provides a detailed dissection of a nuclear state. When combined with the electron- and kaon-scattering data, a determination of the radial forms of the separate neutron and proton transition densities will be possible.

Fig. 2.-9. Ratio of π^+ -proton excitation to π^+ -neutron excitation at $q = 0$ (solid line) and $q = 2 \text{ fm}^{-1}$ (dashed line). Figure is from Ref. (Ca 82).

2.4 QUARK-GLUON PHYSICS

2.4.1 Introduction

In this section we examine four general areas of experiments in which the quark/gluon substructure of matter must certainly manifest itself. Each area holds the potential for experimental discoveries which could revolutionize the description of nuclear systems. A high-intensity, high-energy primary proton beam is crucial for these studies either for achieving the necessary intensity of secondary beams or for examining cross sections which are minute due to the momentum transfer necessary to observe constituent (quark/gluon) scattering.

Subsection 2.4.2 deals with the frontier problem of nuclear structure physics: What role do quarks and gluons play in nuclei? Starting with the evidence provided by the EMC effect, three areas of experimental investigation are proposed which will yield information complementary to existing experiments. They are studies of dilepton production in high-energy hadron-nucleus collisions, deep-inelastic scattering of 30 to 40 GeV muons, and investigations of the A-dependence of high- p_T hadron production.

Subsection 2.4.3 focuses on extensions of hadron spectroscopy which should guide the application of QCD in the nonperturbative region. Surely, one of the most exciting prospects in strong interaction physics is the possible existence of objects containing no valence quarks--gluonium. Discovery of well-defined resonances of the dibaryon system would have enormous impact in guiding bag-model approaches to hadronic structure. The possible existence of a stable "H" dibaryon (discussed previously in section 2.2.2.5) is but one example of the physics promise here. High-intensity kaon and pion beams in the 1 to 5 GeV range are the backbone of this program.

The high intensities of π and K beams at 30 to 40 GeV will make feasible for the first time tests of elegant relations among exclusive pseudoscalar meson-induced reactions yielded by perturbative QCD. Polarized-beam and polarized-target reactions will provide direct tests of constituent scattering theories, which predict very simple asymptotic values for spin observables. Even if the perturbative region for exclusive processes proves to be out of reach, studies of the reactions described in subsection 2.4.4 at very high momentum transfer will teach us much about how to proceed theoretically.

Finally, subsection 2.4.5 presents two approaches to creating exotic conditions in nuclei which are complementary to relativistic heavy-ion collisions. Quark deconfinement over extended nuclear volumes is the goal of both approaches. One method is to implant the energy of a 7- to 8-GeV antiproton or antideuteron inside a heavy nucleus. Computer simulations based on two very different models suggest that conditions approaching those necessary for deconfinement may indeed be within reach. The other venture proposes the use of the high-energy beams in conjunction with low-momentum transfer reactions to excite proposed collective color vibrations of the nucleus. Both ideas are speculative but contain the possible seeds of entirely new research directions.

2.4.2 Quark and Gluon Structure of Nuclei

2.4.2.1 Present Status: The EMC Effect. Deep-inelastic scattering of 100- to 300-GeV muons at CERN has produced convincing evidence that the quark distributions of nucleons in nuclei differ in a small but significant way from those of free-nucleons (Au 82). This is now referred to as the EMC (European Muon Collaboration) effect after the first group to observe it. The EMC effect is the first credible evidence that quark degrees of freedom must be taken into account in nuclear structure. Further experimental evidence that can be brought to bear on this matter is clearly a high priority for nuclear science.

The capabilities offered by LAMPF II have an enormous potential for the elucidation of nuclear quark distributions through the application of the electromagnetic, strong-, and weak-interaction probes. As in the EMC experiment, the studies proposed here rely on precision measurements of anomalous A dependences. For this reason, the high-intensity beams of LAMPF II are crucial. Since LAMPF II will span the region from traditional medium-energy to high-energy physics, quark phenomena in nuclei can be studied at a variety of distance scales. The ability to use probes of different wavelengths may prove pivotal in the understanding of the role of quarks in nuclei since some models invoked to explain the EMC effect in ultra high-energy experiments also have implications for nuclear structure on the scale of mesons and nucleons. Studies of models of the EMC effect with implications at longer wavelengths have been discussed in section 2.3.2.1.

2.4.2.2 The Drell-Yan Process. The annihilation of a quark and antiquark from a hadronic reaction yielding a dilepton pair, known as the Drell-Yan (DY) process (Dr 70), bears a simple space-time relation to deep-inelastic lepton scattering, as can be seen in Fig. 2.-10. It is complementary, since either the q or $\bar{q}$ can come from the beam hadron and a variety of beams will be available. Extensive studies of the DY process have been made with protons and pions at energies in excess of 100 GeV (Ke 82). It now appears that the range 20-50 GeV offers a number of opportunities for important new applications to the study of quark-structure functions of nuclei.

A most exciting application of the DY process to the study of quark structure of nuclei has been suggested recently by Bickerstaff, Birse and Miller (Bi 84). They show that the cross section for the DY process for protons on nuclei may be written as*

$$\frac{d^2\sigma}{dx_1 dx_2} = \frac{K4\pi\alpha^2}{9s} \frac{1}{x_1 x_2} \left[\sum_a e_a^2 v_a^P(x_1) \bar{q}_a^T(x_2) + \frac{1}{x_2} \bar{q}_a^P(x_1) F_2^T(x_2) \right] \quad (1)$$

*There are, of course, approximations involved in using equation 1. It is valid only for the $q\bar{q}$ annihilation process in the factorization approximation; analysis of experiments would use a more complete calculation.

Fig. 2.-10. Feynman diagrams for three processes with a single intermediate photon: (a) the Drell-Yan process, (b) e^+e^- annihilation into hadrons, and (c) deep-inelastic lepton scattering (Ke 82).

where x_1 and x_2 are, respectively, the momentum fractions of the quarks in the projectile and target. For small values of x_1 , the valence quark distribution in the proton, v_a^P , is small, resulting in dominance of the cross section by the term on the right which contains the electromagnetic F_2 structure function of the target. Hence, at small x_1 , the DY process should be very similar to deep-inelastic lepton scattering. In contrast, at large values of x_1 , $v_a^P(x_1)$ is much larger than $\bar{q}_a^P(x_1)$, the antiquark distribution of the projectile. This leads to a dominance of the left-hand term and a consequent sensitivity to $\bar{q}_a(x_2)$, the target sea-quark distribution, not found in the lepton-scattering experiments.

The complementary nature of the DY process at large x_1 becomes evident when one examines the predictions of several models proposed to explain the EMC effect. Three distinct models considered by Bickerstaff, Birse, and Miller are an enhanced nuclear pion field (Er 83), six-quark clusters (Ca 83), and rescaling of the nucleon size (Ja 83); a fourth model incorporates some of the physics of the enhanced pion field and six-quark cluster models. The predicted A dependence of the DY process is shown in Fig. 2.-11. For small x_1 , the DY A dependence closely resembles that of the F_2 function. At large x_1 , however, there are remarkable differences arising from the different assumptions made about the sharing of momentum between valence-quark and sea-quark distributions. Clearly, the DY process could provide essential clues about the origin of the EMC effect and the role quarks play in the structure of nuclei.

Fig. 2.-11. Calculations of the ratio of Drell-Yan cross sections of iron and deuterium at two values of x_1 (Bi 84).

The requirements necessary to exploit the complementary features of the DY process are met by a high-intensity proton beam with energy in the 30 to 50 GeV range. Counting rates in the range of 10 dilepton pairs per second with $x_1 x_2 > 0.15$ can be achieved with microamp-level beams on targets of 1 interaction length using a spectrometer similar to existing ones (Ke 82). The advantage of using energies lower than the previous CERN SPS and FNAL fixed-target experiments is clear from equation 1. For given values of x_1 and x_2 , the cross section varies as $1/E_{\text{beam}}$. The lower bound on the beam energy is given by the requirement that dilepton production be consistent with the DY process. Beams of 40 GeV seem safely within this bound and offer enough margin to use dilepton pairs with $m > 4 \text{ GeV}/c^2$, hence above the region of strong vector-meson resonances.

Other new opportunities for DY experiments at LAMPF II could be realized through the properties of $K^\pm$ beams. First, one could perform kaon-proton DY experiments of much higher quality than now exist in order to vastly improve the knowledge of the structure functions associated with these beams. Once they are known, kaons can be employed in elegant tests of nuclear structure functions. For example, $q\bar{q}$ annihilation by K^+ valence quarks can come only from the target-quark sea since the K^+ contains $\bar{s}u$; the $\bar{u}$ component of the K^- , in contrast, can annihilate with a valence u quark from the target. The use of an rf separator to achieve high-purity high-energy kaon beams appears to be feasible (see Chapter 5).

2.4.2.3 Deep Inelastic Muon Scattering. A recent detailed study at SLAC (Ar 84) of deep-inelastic electron scattering from a series of nuclei using beams of 8 to 24.5 GeV has confirmed the anomalous A dependence found in the EMC data for values of $x > 0.25$. However, the excess of scattering cross section for heavy nuclei in the region $x < 0.25$ observed in the EMC experiment is not found in the SLAC results. Speculation about the differences between the two experiments usually centers on the effect of nuclear shadowing, i.e., at sufficiently small Q^2 , the virtual photon exchange may be shielded from the nuclear interior. Although the SLAC data show no obvious Q^2 dependence, the small-x data were taken at much smaller values of Q^2 than the EMC data.

It is unlikely in the next few years to be able to extend the electron results to higher energies. Experiments planned at FNAL using ~700 GeV muons will provide exciting new vistas in deep-inelastic lepton scattering. Beams of muons up to 40 GeV, obtainable from π and K decays at LAMPF II, will fill an important gap in the study of the quark structure of nuclei. Muons have the additional advantage, with respect to electrons, that shadowing corrections are more easily handled. Intensities in the range of 10^8 to 10^9 /sec with good duty factors offer the possibilities of more exclusive experiments such as $(\mu, \mu'X)$ where X may be a single hadron or more complex final state. An example of such an experiment discussed recently (Bo 84, Gu 84) would be $(\mu, \mu'\pi)$, in which pion enhancement models of the EMC effect would be examined using the relation

$$\frac{d^2\sigma}{dx \cdot dQ^2} = G_{\pi/N}(x) \frac{d\sigma}{dQ^2} (\mu\pi \rightarrow \mu\pi) \quad , \quad (2)$$

where $G_{\pi/N}(x)$ is the probability of finding a pion with fraction x of the nucleon's momentum. The pionic-enhancement models (Er 83) predict $G_{\pi/N}$ to be larger in a complex nucleus due to many-body effects in the nuclear pion field (another way to look for such an enhancement using neutrino beams is discussed in section 2.3.2.1).

2.4.2.4 High- p_T Hadronic Reactions - Implications for the Quark and Gluon Structure of Nuclei. Purely hadronic processes with high-transverse-momentum outgoing particles have enormous potential for studying the modification of nucleon structure in nuclei. Here, in contrast to the electromagnetic experiments discussed previously, both quark and gluon degrees of freedom enter directly. Thus, in probing EMC-type anomalous A dependences, one has the possibility of new physics related to the fundamental field of QCD, as well as an additional complication to disentangle. This is in familiar analogy to lower-energy hadronic reactions, where ingenuity must be applied to disentangle nuclear-structure effects from reaction-mechanism effects. As in the low-energy regime, the additional observables available with polarized beams and targets (of light nuclei) should help enormously.

Although many issues connected with the penetration of nuclei by energetic hadrons are still under debate (Hw 84, Bu 84), it is clear for high- p_T particle production that a coherent wave of hadronic matter penetrates to the center of even the heaviest nuclei. The A dependence of the hadron-nucleus cross section is commonly written as $\sigma_{hA} \propto A^\alpha$. For small p_T processes, α is near 2/3, the

expected value characteristic of nuclear-surface dominance due to strong absorption. Beginning at values of $p_T \geq 2$ GeV/c, α climbs toward unity, indicating the absence of shadowing. Then, in apparent violation of common sense, α rises significantly above unity. The anomalous A dependence exemplified by the Cronin effect (Cr 75) preceded the EMC effect by some years and yet did not raise many eyebrows in the nuclear physics world. Recently, Baldin (Ba 84) has proposed a model for high p_T pion production from proton beams on nuclei that produces a ratio of "structure functions" which is very reminiscent of the EMC effect (Fig. 2.-12).

Although a number of details of Baldin's analysis need further investigation, the richness of possible hadronic reactions invites imaginative experimentation to test the details of this and other models. Concentration on the highest p_T processes possible appears to be the best possibility for achieving conditions where perturbative QCD may offer theoretical guidance. This fact alone argues for the highest possible primary beam intensities.

2.4.3 Hadron Spectroscopy

The original evidence for quarks came from the structure of hadron spectra. This venerable field, although often described as having had a "golden age" nearly 15 years ago, offers frontier physics opportunities for the study of confinement aspects of QCD. Two areas, the search for gluonium and for

Fig. 2.-12. Ratios of "structure" functions derived from high- p_T pion production (triangles) compared to ratios for deep-inelastic lepton scattering (circles) (Ba 84).

multiquark structures such as dibaryons and baryonium, are qualitatively different from the spectroscopy of the late 1960's.

The experimental program required to unambiguously identify these new hadronic states will necessarily benefit conventional areas of hadron spectroscopy. The wealth and detail of the data obtained will enable us to both deepen and extend our understanding of hadrons composed only of quarks and antiquarks.

The LAMPF II facility opens new vistas in spectroscopy in two major ways. First, by producing kaon beams of an intensity and quality comparable to pion beams during the "golden age," LAMPF II not only advances the traditional spectroscopy of mesons and baryons but also makes feasible a large class of experimental searches for structure in the strange or doubly-strange dibaryon systems (such as the "H" particle discussed in section 2.2.2.5 and again in 2.4.3.2). The second major qualitative advance of LAMPF II would be the production of (electromagnetic- or time-) separated $\bar{p}$ beams in the 1 to 5 GeV range with at least two orders of magnitude higher intensity than LEAR. The first round of baryonium physics at LEAR has not produced definitive evidence for resonances of bound states. It is unlikely that this oversubscribed and undersupported facility will deliver enough antimatter to exhaust the possible richness of this field in the likely event that structures manifest themselves only as small parts of inclusive cross sections.

2.4.3.1 Gluonium. The existence of gluonium--colorless, flavorless composites of two or more gluons--is one of the basic tenets of QCD. These proposed quarkless states of matter should display a rich spectrum of excited states of various spins and parities. Estimates of their mass spectra have been carried out in lattice gauge theories, bag models and constituent massive gluon models. Unfortunately, their predictive powers are minimal; none of them is the analog of the potential models which have been so successful in treating heavy quarkonia with correct detailed predictions of masses and couplings. Consequently, it falls to experiment to offer guidance in the quest for this new form of matter.

Exciting current candidates are three $J^{PC} = 2^{++}$ resonances in the 2 GeV/c² mass range which have been found in the Zweig-rule-violating process $\pi^- p \rightarrow \rho p n$. To confirm the $\rho\rho$ resonances as gluonium states it is necessary that these resonances have other decay modes from which we can verify their flavor-singlet property. LAMPF II, because of its high beam intensity and consequent high event rate, is ideal for performing analogous Zweig-violating experiments to both confirm the present results and to find new conditions.

Other glueball candidates arise in the 0^{-+} and 2^{++} states formed in the radiative decays of the J/ψ (3100). Recent experiments with the Mark III detector have produced an incredibly rich spectrum of 0^{-+} and 2^{++} states heretofore unknown (Me 84). It is now clear that these states cannot be treated singly, because they are strongly mixed with each J^{PC} sector. The careful elucidation of the interrelationship of these states is necessary to establish the gluonium nature of any of them. Again, the careful study of these states with the intense pion and kaon beams of LAMPF II should complement the radiative formation of these states and help in their identification as gluonium or as other interesting states, e.g., quark-antiquark-gluon composites, or multiquark states.

2.4.3.2 The Dibaryon and Baryonium Systems. Elucidating the nature of the nucleon-nucleon interaction, both in free and in bound systems, has been the central problem of nuclear physics for half a century. Great strides have been made in the long- and medium-range parts of the force. The boson exchange picture seems to provide an adequate description for both the NN and the $\bar{N}N$ systems through the G-parity transformation. For distances shorter than a Fermi, however, little progress has been made and most models resort to pure phenomenology.

The QCD/quark revolution in strong interaction physics of the last decade gives us an entirely new viewpoint on this problem. The more general problem is that of six quarks, or three quarks and three antiquarks, interacting via gluon exchange. We shall never have a complete understanding of the short-range NN ($\bar{N}N$) interaction without an overall comprehension of the two-baryon (baryon-antibaryon) system.

At present, the most reliable guides we have to the implications of QCD in multiquark systems [$N(q + \bar{q}) > 3$] are the bag models. The apparently necessary condition of quark confinement is put into these models in an ad hoc manner. The P-matrix analysis is then applied to relate states of the multiquark bag to physical states of mesons and baryons at large separation. Bag models generally predict a rich spectrum of six-quark structures in the NN system and $2q-2\bar{q}$ states in the $\bar{N}N$ system. In neither system have any of these predicted new states been convincingly identified. Intriguing structures suggestive of resonances have been seen in total cross sections for pure nucleon spin states (Au 78). Interpretation of these structures is still under intensive theoretical investigation.

LAMPF II will provide beams of kaons, pions and antiprotons, and tagged beams of antineutrons, hyperons (possibly polarized) and antihyperons which will make feasible a broad-range program of dibaryon and baryon-antibaryon spectroscopy. For the first time it will be possible to examine reactions with extremely small cross sections--reactions, which by virtue of large isospin or strangeness transfer, are more likely to lead to narrow states.

Non-Strange Dibaryons. The need for qualitatively different experimental capabilities in the search for multiquark phenomena is illustrated in a tantalizing spectrum from a recent Saclay experiment (Fig. 2.-13) (Ta 84). The possible dibaryon structure is the broad peak near 2.24 MeV; it is consistent with the structure observed in the pure spin cross section experiments. Obviously, complete experiments with polarized beams as well as decay measurements are necessary to determine whether or not this modest structure is the beginning of a new spectroscopy. Low-energy (1-5 GeV) polarized-proton beam experiments could become a major program at LAMPF II.

There are also a number of new possibilities created by LAMPF II secondary beams in the area of non-strange dibaryon resonances. The formation of an $I = 2$ dibaryon system via $\pi^+ + {}^2\text{H} \rightarrow \pi^- + X$ would have the advantage of an isospin-forbidden decay into two nucleons. Antiprotons or antineutrons can be used to achieve different kinematic conditions for the formation of dibaryons via reactions such as $\bar{N} + {}^4\text{He} \rightarrow N + X$. Noting that the signal-to-noise conditions may be typified by Fig. 2.-13, such experiments are not feasible without LAMPF II intensities.

Fig. 2.-13. Missing-mass spectrum for the $^3\text{He}(p,d)X$ reaction (Ta 84).

Strange Dibaryons. One of the most direct experimental confrontations with quark models involves the search for the $S = -2$ H particle (discussed previously in section 2.2.2.5). Using the MIT bag model, Jaffe (Ja 77) predicted the existence of a stable S-wave flavor-singlet dibaryon with $J = 0$ at 2.15 GeV. This completely symmetric system consists of two u quarks, two d quarks, and two s quarks. Since all of the quarks are in a relative s state and have their spins antiparallel, a significant binding energy is involved. This particle is analogous in the six-quark sector to the alpha particle in nuclear physics but with three different constituents instead of two.

A number of methods have been proposed to look for this particle, including the two-step process $K^- + p \rightarrow K^+ + \Xi^-$ followed by $\Xi^- + d \rightarrow H + n$ and the one-step process $K^- + ^3\text{He} \rightarrow K^+ + n + H$. One might also create the H via $\bar{p} + p \rightarrow \bar{H} + H$. Counting rates, even for the kaon and antiproton fluxes available at LAMPF II, are estimated to be only a few per day. Thus, a definitive experiment without a facility like LAMPF II seems to be out of the question unless the "H" is stable and the production cross section is much larger than expected.

Although the $S = -1$ two-baryon system is not expected to bind as tightly as the flavor singlet $S = -2$ state, there are bag-model $S = -1$ states which should have small coupling to Y^*N channels. Consequently, the ambiguity found in searches for the NN dibaryons due to coupled-channel effects will be much less important; in addition, the $S = -1$ dibaryons are expected to be quite narrow. At present, there is intriguing preliminary evidence for dibaryon structure in the $S = -1$ system seen both in the ${}^2\text{H}(\text{K}^-, \pi^-)X$ and ${}^2\text{H}(\pi^+, \text{K}^+)X$ (Pa 81) reactions (Fig. 2.-14). Clearly, there is urgent need for additional data of higher quality.

To search for baryons with one or more units of strangeness, experiments involving the scattering of Λ 's, Σ 's, and Ξ 's from nucleons are of critical importance. At Fermilab energies, relativistic time dilation allows beams of these particles to be made so that the scattering problem can be addressed more easily; however, the very high energies put one outside the range of resonant structures. The more difficult low-energy scattering experiments would use the intense meson and antiproton fluxes available at LAMPF II to great advantage. Tagged hyperon beams can be made using, for example, the $\text{K}^- + \text{p} \rightarrow \pi^- + \Sigma$ or the

Fig. 2.-14. Missing mass spectrum from the ${}^2\text{H}(\pi^+, \text{K}^+)X$ reaction at 1.4 GeV/c (Pa 81). The largest peak near 2.14 GeV is a threshold phenomenon. The peak just above it (near 2.17 GeV) is a possible dibaryon state.

$\bar{p} + p \rightarrow \bar{\Sigma} + \Sigma$ reactions. These types of reactions have been studied in a hydrogen bubble chamber to obtain the few hundred low-energy hyperon-nucleon events presently available. LAMPF II will allow this kind of experiment to be done using counters, giving a large increase in event rate and simplifying the data analysis. With the use of polarized hydrogen targets, beams of polarized hyperons can also be produced and scattered from polarized protons (Fig. 1, Ma 83). This opens up an exciting area of physics and gives us an entirely new and perhaps more selective view of the six-quark system.

Baryonium. After a year of searches for bound and low-lying resonance states of the $\bar{p}p$ system at LEAR, no clear evidence of narrow baryonium states exists. Measurements of the $\bar{p} + p$ total cross section have found no evidence of the so-called S meson (Cl 84), and discrete gamma transitions from $\bar{p}p$ interactions at rest have not been observed. Still, there are suggestions of discrete pion lines in the $\bar{p}p$ and $\bar{p}d$ systems, which demand further investigation. It is very likely that baryonium will remain a subject of active theoretical and experimental interest for many years.

LAMPF II will provide beams of antiprotons of up to three orders of magnitude higher intensity than LEAR. For many experimental investigations, the intensity increase will more than offset the disadvantages of using separated rather than cooled beams. LAMPF II, unlike the LEAR of present and near-term future, would be dedicated to this area of physics for six months rather than approximately one month per year.

As an example of an experiment which would be feasible with LAMPF II and only marginal with the present day LEAR, consider the very interesting suggestion of Moalem et al. (Mo 84). They propose direct formation of bound NN states via single-nucleon pickup in reactions such as $\bar{p} + {}^4\text{He} \rightarrow {}^3\text{H} + \text{baryonium}$. Typical predicted peak cross sections are 300 $\mu\text{b/sr}$. Assuming conditions typified by the SPES II spectrometer system (Ga 82), only ~ 3 counts/hour would be obtainable at LEAR versus 3000/hour at LAMPF II. In the likely event that these states would be observed as small bumps on a large annihilation background, proof of the existence of bound baryonium states via this reaction would be exceedingly difficult at LEAR. More complex experiments involving decays among baryonium states are clearly impossible without LAMPF II intensities.

Finally, non-zero strangeness and/or isospin may provide additional clues about the nature of baryonium structures in a manner analogous to the dibaryon system. Intense beams of antiprotons may be used as sources of either tagged antineutrons or antilambdas via the reactions

The $\bar{n} + p$ system is a better way to study the pure $I = 1$ baryonium system than with an antiproton beam on a deuterium target. Tagged $\bar{\Lambda}$'s will open entirely new dimensions to the search for baryonium states.

2.4.4 Hadronic Reactions at High Momentum Transfer

In this section, we discuss experiments at high momentum transfer and the possibility of confronting QCD predictions. Although the momentum transfers available at LAMPF II might not be high enough to reach the asymptotic region (where QCD perturbation theory applies), a systematic experimental and theoretical program is necessary for the study of the transition from confinement to asymptotia. The high intensity of LAMPF II beams will make possible many valuable and possibly crucial experiments. The greatest immediate impact of these discoveries is in nuclear physics, since the effective N-N interaction and nuclear structure involve the same distance scales as the confinement-asymptotic freedom transition.

2.4.4.1 High Momentum Transfer Spin Physics. Polarization physics at large momentum transfer will become one of the forefront areas of medium/high-energy physics. This is perhaps a surprising statement since at asymptotically large momentum transfer, where QCD perturbation theory applies, polarization phenomena should be very small. In particular, since quark helicity is conserved in massless QCD perturbation theory, one obtains very simple relationships for spin-correlation parameters with sufficiently high momentum transfer. In contrast to the simplicity of perturbative QCD, current experiments indicate that there is a rich set of spin phenomena which more than challenge current theory. One of the most striking signals of this new physics is the data on hyperon polarization in high-energy pp collisions. Since Λ polarizations are readily measured by the Λ -decay process, the polarization of the Λ in $p + p \rightarrow \Lambda + X$ has been measured over a wide range of energies. The large (15-20%) analyzing power at $p_T \approx 2-3.5$ GeV/c cannot be understood in simple QCD models, and has led to considerable theoretical interest.

Even more startling are the recent AGS results on the analyzing power in pp scattering, which is found to go from $\sim 16\%$ at $p_T = 1.7$ GeV/c to $\sim 51\%$ at 2.6 GeV/c (Ph 83). This dramatic increase cannot be explained with present theory. Even the 16% asymmetry at $p_T = 1.7$ GeV/c is not understood. Although it is conceivable that one can explain this phenomenon by including additional processes within the conventional QCD framework, it is likely that a new quark/gluon phenomenon has been uncovered. Similarly, the large p_T ZGS results on spin-spin correlation in 12 GeV proton-proton elastic scattering are still not understood after 6 years. There, it was found that the ratio of spin-parallel to spin antiparallel scattering increased sharply from about 1 to about 4 as one moved to $p_T = 1.9$ GeV/c. These two totally unexpected results remain a challenge to any theory which attempts to predict the short-range behavior of hadronic interactions.

The experimental results seem to imply that there are strong quark-gluonic processes not seen in the quark-quark phenomenological forces which have been derived from QCD perturbation theory and fits to hadron spectra. One of the most striking observations from quark models of baryons is that the spin-orbit forces obtained from fits to experiment are much smaller than predicted by one-gluon exchange. The precise form of the interaction is not known, but powerful non-perturbative gluonic forces, possibly related to the confining mechanism, seem to be coming into play.

The extension of spin-observable measurements to higher momentum transfer is of high priority. The cross sections fall rapidly with momentum transfer, so the experiments will require as much intensity as possible. The 45 GeV high-intensity LAMPF II accelerator would thus play an essential role in exploring this fundamental new physics.

2.4.4.2 Pseudo-Scalar Meson-Nucleon Reactions. When and if it applies, perturbative QCD brings a beautiful simplicity to the class of reactions involving pseudoscalar mesons and nucleons. Farrar (Fa 84) has shown that, in perturbation theory, 12 basic diagrams are sufficient to describe 66 different reactions. The result is that there are numerous relations among the various possible reactions. As an example one finds the following cross section relation:

$$1/\sqrt{3} (\pi^- + p \rightarrow \pi^+ + \Delta^-) = (\pi^- + p \rightarrow K^+ + Y^-) = (K^- + p \rightarrow \pi^+ + Y^-) .$$

Tests of such predictions are inconceivable without a high-intensity accelerator. Although it has been suggested (Is 84) that perturbative QCD may not be applicable to any exclusive reaction at accelerator (even SSC) energies, LAMPF II tests of these relations could be valuable in determining aspects of the symmetries of QCD that are maintained in the effective quark Lagrangian.

It has been pointed out by Isgur and Llewellyn-Smith (Is 84) that "processes such as meson-baryon scattering are sufficiently rich that they allow elegant tests of the applicability of perturbative QCD which are free from ad hoc assumptions about hadronic wave functions" (necessary in the description of form factors).

2.4.4.3 Parity Violation. Parity-violation experiments test the fundamental weak interaction of quarks and quark nuclear structure. At sufficiently high momentum transfer, gluonic corrections become manageable and reactions such as $\bar{p} + p \rightarrow \pi + X$ could be interpreted in terms of weak quark-quark scattering. Even with the improvements in polarized-beam quality and intensity at LAMPF II, it is likely that semi-exclusive reactions at large P_T will not be immediately practical. In spite of this, high-energy parity-violating total cross section experiments of the type performed at LAMPF I and the ZGS (Lo 80) will provide new ground for testing quark models of NN and N-nucleus interactions (Go 83). In time, more exclusive measurements may yield even more detailed tests of the short-range structure of hadronic matter.

2.4.5 Exotic Forms of Nuclear Matter

2.4.5.1 The Quark-Gluon Plasma via Antimatter-Matter Annihilation. One of the most promising areas in which to observe the effects of quarks in nuclei is to subject nuclear matter to conditions of extreme temperature or pressure. Many calculations have predicted that under such conditions normal nuclear matter will undergo a phase transition to a quark-gluon plasma in which individual nucleons lose their identity. Various signatures for this transition have been suggested, mostly involving the production of pairs of leptons or strange mesons with a different probability or a different spectral shape than

would be produced by normal nuclear matter under these conditions. If the transition from normal matter to quark-gluon matter is sharp, then the line representing the boundary of the phase transition (drawn in temperature-density space) is predicted by lattice-QCD calculations to have a zero-density intercept around 180 MeV. The zero-temperature intercept is much less certain with predictions ranging from ~ 3 times normal nuclear-matter density to 10-20 times normal density. Colliding heavy-ion beams have been suggested as a method for achieving the conditions necessary for the formation of the postulated plasma.

A complementary method of achieving extreme conditions is the annihilation of antiprotons or antideuterons in nuclei. The release of energy in the form of a number of pions rapidly heats a small volume of nuclear matter. Studies to date on the annihilation of $\bar{p}$'s in nuclei have been done only with low incident energies (Di 84).

As the momentum of the incident $\bar{p}$ is increased, two events conspire to make the energy deposition more effective. First, the total antinucleon-nucleon cross section decreases strongly so that the annihilation can take place further into the nucleus (we assume approximately 1 fm inside the nuclear surface for the statements that follow). Thus, those pions that scatter back on their first collision are still within the nucleus. The second effect is that the pions are focused forward, so no high-energy pions are traveling backwards in the laboratory frame, and, hence, the pions must traverse the entire nucleus in order to escape. The number of pions produced also increases. Finally, one must allow for the formation or hadronization of the pions, a period during which the highest energy primordial pions may travel a significant fraction of a nucleon diameter without interacting. All of these effects increase significantly the fraction of energy deposited. In intranuclear cascade (INC) calculations, about two-thirds of the available energy is deposited for the above conditions. This fraction is roughly constant for incident momenta ranging from 2-6 GeV/c.

As the energy increases and the fraction of energy left in the nucleus stays constant, the absolute amount of energy deposited increases and, with it, the temperature. Figure 2.-15 shows relativistic hydrodynamic calculations that display clearly the temperature evolution of various regions of the nucleus. In these calculations, the energy was deposited in a manner determined by the mean-free paths of pions, the momenta of which were obtained from a Monte Carlo calculation. The least certain ingredient is the equation of state that relates pressure to density and temperature.

By looking at the number of nucleons as a function of their kinetic energy in the INC calculations, "hot" components can be clearly discerned and its temperature obtained from the slope curve. Figure 2.-16 shows the development of this temperature as a function of time for various incident momenta. If the momenta become too large, then hadronization occurs outside the nucleus. Thus, we have not considered cases above 10 GeV/c. One way to overcome this problem is to use an antideuteron as the energy source. In this case the velocity of the center of mass is much less and the average pion energy is less for the same total energy available. A curve of this type is also shown in Fig. 2.-16.

As can be seen, the use of energetic $\bar{p}$'s or $\bar{d}$'s could yield temperatures high enough to form the hypothetical quark-gluon plasma. Even if the predictions of a phase transition were to be withdrawn, measurements of the

Fig. 2.-15. Relativistic hydrodynamics calculation of the annihilation of an antiproton in uranium. The figure (on opposite page) shows the time history of the annihilation in six frames at intervals of 1.15 fm/c. Temperature is indicated by color, with a scale given at the left of the figure. At time zero, an antiproton annihilates. After a hadronization delay, the pions produced in the annihilation interact with the uranium. Notice the white hot region which persists for 5 frames. In this region, the temperature may be hot enough for a long enough time to cause a transition to the quark gluon plasma phase.

6 GeV $\bar{p}$ annihilation on ^{238}U

8 pions Central two fm

T_{MeV}

200

0.0

3.4

150

1.1

4.6

100

2.3

5.7

50

Fig. 2.-16. Comparison of the temperatures obtained as a function of time for the two models considered. The "at rest" absorption occurs on the surface and the annihilation for the other cases takes place on the beam axis just inside the nuclear surface (0.7 fm for the hydrodynamic calculations and 1 fm for the INC). The numbers in parentheses give the actual energy deposited in the nucleus in GeV. The momentum quoted for the $\bar{d}$ is per antinucleon.

behavior of nuclear matter at moderate density and high temperature would still be of fundamental importance. The calculations indicate antimatter annihilation to be an excellent method to investigate the thermal degrees of freedom of the equation-of-state and thus the effective mass of the nucleon in regions of high temperature.

2.4.5.2 Color Excitations of Nuclei. Nuclei are made of quarks interacting through the color field. However, in the nuclear response to probes at low energy-momentum transfer, the color degree of freedom and its sources (quarks and gluons) are hidden. A highly speculative but extremely exciting task for the nuclear physicists of the 1990's and beyond will be to search for collective color modes of excitation of nuclei.

All currently-known nuclear excitations involve only the flavor degrees of freedom of quarks. The traditional description of such excitations involves composite hadronic degrees of freedom (N , Δ , π , p ...). The character of intrinsically-colored excitations of nuclei is profoundly influenced by the phenomenon of color confinement. Consider, for example, color generalizations of the giant-dipole resonance in nuclei, where the neutrons oscillate harmonically against the protons. Let, say, the red quarks oscillate in step against the blue quarks. Just as the conventional giant dipole excitation is isovector, so must the "color dipole" carry color. The operators describing such excitations are necessarily of complicated many-body character; a classical description is probably inescapable.

Two examples of color excitations have been discussed. One is a localized hot spot of quark-gluon plasma. This may be produced in ultrarelativistic heavy-ion collisions or, as we have seen, by shaking heavy nuclei with antimatter. A second color mode has been suggested recently (Ca 84a)--a localized ring of nucleons in a massive nucleus within which quarks move freely. Estimates of the excitation energy of such modes yield 50 to 100 MeV per nucleon--a few GeV in a medium-weight nucleus. The ring excitations may be so different from lower-energy states of nuclei that they have small widths (a few tens of MeV). The limited spatial extent of the ring excitations results in small angular momentum. Such excitations may be accessible in high-energy, low momentum-transfer hadronic reactions--such as (π^+, p) , (p, d) , (π^+, Λ) , wherein the outgoing particle is more massive than the projectile.

Two types of color excitation must be considered--those involving changes in the color configurations of quarks and the excitations of the gluon field. Excitations of the glue are likely to be just as important as quark excitations, since about half the energy-momentum content of free nucleons resides in the glue. But leptonic probes couple only to the quarks. Thus in the search for color excitations, leptonic, hadronic and heavy-ion probes will complement each other. Leptons will excite quark excitations via interactions of great simplicity; hadrons will excite gluonic excitations and reveal aspects of the quark excitations complementary to those seen by leptons; heavy ions will create large spots of high-density, high energy-density hadronic matter.

Studies of color excitations of nuclei could be a major focus of interest in the 1990's and beyond. If detected, they would shed light on the central problem of many-body hadron physics--quark confinement in the many-body environment. We can not at present foresee the details of the experimental program that will reveal the color excitations. It is, however, clear that a combined attack, with leptonic, hadronic and heavy-ion probes, will be most likely to lead to definitive results.

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CHAPTER 3.

ELECTROWEAK INTERACTIONS

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3. ELECTROWEAK INTERACTIONS

3.1 INTRODUCTION

3.1.1 The Standard Model and Its Family Problems

The ultimate goal of physics is a complete description of the fundamental particles and forces of nature. The recent unification of the electromagnetic and weak forces is the culmination of the efforts of many experimental and theoretical physicists during much of this century. It is a milestone comparable to the unification of the electric and magnetic forces by Maxwell a hundred years ago. This theoretical model of the electroweak interactions is consistent with all known physics. However, there remain many unanswered questions so that it is probable that this model is incomplete. This is analogous to Maxwell's electromagnetism being correct but incomplete in that it was unable to deal with radioactivity.

In this chapter we will show how a facility with the capabilities of LAMPF II is essential for continued progress in this area. The high fluxes of high-quality beams that LAMPF II will produce will allow precise tests of higher-order corrections to the predictions of existing models; such tests will demonstrate that we do have a theory of the electroweak force with great calculational and predictive power. Even more important, LAMPF II will be an important tool to investigate the possibility of new particles and interactions, especially those related to the problem of quark-lepton families, and it will enable us to probe new physics at mass scales far beyond those achieved in direct experiments, even with the next generation of high-energy accelerators. This combination of precision tests and searches for physics at very high mass scales forms the high precision frontier.

The theoretical framework for the following discussion is called the "standard model." It consists of the $SU(3)_C$ gauge field theory called quantum chromodynamics and the spontaneously broken $SU(2)_W \times U(1)_B$ gauge field theory known as the "electroweak model." Recent experiments at CERN have directly produced the W and Z particles, which are the massive intermediate vector bosons of the weak force. These new particles, together with the photon, are responsible for the electroweak force.

The fermions in the standard model come in repeating sets called families (or generations). Each family consists of a pair of quarks and a pair of leptons; at the present time three families have been observed. They are shown in Table I and their mass spectra are shown in Table II. For future reference we note that there are large differences between the masses of different families.

All of the mesons and baryons and their strong, weak, and electromagnetic interactions, and the nuclei and atoms formed from them, can be described, in principle, by combinations of these fermions and their couplings to the gauge bosons, but many physicists feel that a set of 58 fundamental elements is still too many. There are other problems as well. Why are there three "gauge groups," $SU(3)_C$, $SU(2)_W$, and $U(1)_B$, instead of one? Where does gravity fit in?

Are the quarks and leptons really separate, or is there a dynamics that relates them? Do they have a common substructure, that is, are they composite also? Why are there three families, or are there more?

TABLE I Families of the standard model.

$\begin{pmatrix} u_r, u_b, u_g \\ d_r^-, d_b^-, d_g^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L$	$(d_r, d_b, d_g)_R$	$(u_r, u_b, u_g)_R$	$(e^-)_R$
--	--	---------------------	---------------------	-----------

$u \equiv$ up quark, $d \equiv$ down quark, $e \equiv$ electron, $\nu_e =$ electron-type neutrino

$\begin{pmatrix} c_r, c_b, c_g \\ s_r^-, s_b^-, s_g^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L$	$(s_r, s_b, s_g)_R$	$(c_r, c_b, c_g)_R$	$(\mu^-)_R$
--	--	---------------------	---------------------	-------------

$c \equiv$ charm quark, $s \equiv$ strange quark, $\mu \equiv$ muon, $\nu_\mu \equiv$ muon-type neutrino

$\begin{pmatrix} t_r, t_b, t_g \\ b_r^-, b_b^-, b_g^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L$	$(b_r, b_b, b_g)_R$	$(t_r, t_b, t_g)_R$	$(\tau^-)_R$
--	--	---------------------	---------------------	--------------

$t \equiv$ top quark, $b \equiv$ bottom quark, $\tau \equiv$ tau, $\nu_\tau \equiv$ tau-type neutrino

$r \equiv$ red, $b \equiv$ blue, $g \equiv$ green: the three colors of QCD
 u, c, t have electric charge $+ 2/3 e$.
 d, s, b have electric charge $- 1/3 e$.

Theories that identify the known dynamical symmetry groups as subgroups of a larger, encompassing group are known as Grand Unified Theories (GUTs). Here both quarks and leptons appear in the same representation of this "unifying" group, and so one predicts dynamical relations between them. The most startling of these predictions is that a proton will decay into leptons and mesons. Major non-accelerator experiments are searching for evidence of proton decay.

Inclusion of gravity is a sort of "super" unification; it requires the introduction of "supersymmetry," a dynamical symmetry between bosons and fermions. Technical problems with this approach (it has too few degrees of freedom) led to serious consideration of composite quarks and leptons where the multiplicity of families reflects the existence of a substructure.

Here we believe a more conservative approach is in order. Following the lines that led to the standard model, we separate the questions of families from those of GUTs, etc., and simply ask if there are dynamical relations between the different families and what symmetry group describes them. A general theorem due to Goldstone offers two mutually-exclusive possibilities for the realization of a broken symmetry in field theory.* One is the existence of a massive vector boson; the other is the appearance of a massless scalar (spin 0) boson. Present experiments already provide constraints on these two possibilities, but there is insufficient information to choose between these alternatives or to suggest

other possibilities. In the absence of positive experimental results, theorists do not even have a guide as to how to assign family quantum numbers.

TABLE II Fermion masses.

<u>Family 1</u>	
$m_u \sim 5 \text{ MeV}/c^2$	$m_e = 511 \text{ keV}/c^2$
$m_d \sim 10 \text{ MeV}/c^2$	$m_{\nu_e} < .060 \text{ keV}/c^2$
<u>Family 2</u>	
$m_c \sim 1500 \text{ MeV}/c^2$	$m_\mu = 105.6 \text{ MeV}/c^2$
$m_s \sim 150 \text{ MeV}/c^2$	$m_{\nu_\mu} < 0.265 \text{ MeV}/c^2$
<u>Family 3</u>	
$m_t \sim 40 \text{ GeV}/c^2$	$m_\tau = 1782 \text{ MeV}/c^2$
$m_b \sim 5 \text{ GeV}/c^2$	$m_{\nu_\tau} < 164 \text{ MeV}/c^2$

3.1.2 Unlocking the Family Secrets

The challenge is to provide the capability for measurements of such precision that the keys to understanding the families of particles will be uncovered. Two classes of experiments that require a LAMPF II are family-nonconserving decays and neutrino mass measurements. These are highlighted here with more complete discussions contained in later sections.

Decays from one family to another, which depend on the presumed "family gauge bosons," have not been observed. Table III lists a few representative processes that could provide the needed information and the present experimental limits for them.

The observation of one of the decay modes listed in Table III would constitute direct evidence for a family-nonconserving interaction. It would also provide a measurement of the mass scale associated with family symmetry breaking. The weak interactions are rare processes with slow rates due to the large ($\sim 100 \text{ GeV}$) dynamical scale set by the W and Z boson masses. We can interpret the absence of observed family changing decays as being due to a similar, but larger, dynamical scale associated with the breaking of "family symmetry," that is, the mass, M_F , of the "family gauge vector boson" is large. The stringent limit on $K_L \rightarrow \mu e$ shown in Table III allows us to put a lower bound on M_F .

*We do not consider here explicit symmetry breaking or discrete symmetries. Neither provides well defined accessible experimental tests.

TABLE III Experimental limits on some family nonconserving processes. The branching ratio is the fraction of the total number of decays occurring to a specific final state.

Process	Branching ratio limit	Ref
$\mu^+ \rightarrow e^+ \gamma$	$< 1.7 \times 10^{-10}$	Ki 82
$\mu^+ \rightarrow e^+ \gamma \gamma$	$< 3.8 \times 10^{-10}$	Gr 84a
$\mu^+ \rightarrow e^+ e^+ e^-$	$< 2.4 \times 10^{-12}$	Ei 84
$\pi^0 \rightarrow \mu^+ e^- \text{ or } \mu^- e^+$	$< 7 \times 10^{-8}$	Bry 82
$K^+ \rightarrow \pi^+ \mu^\pm e^\pm$	$< 7 \times 10^{-9}$	Di 76
$K^+ \rightarrow \pi^+ \mu^+ e^-$	$< 5 \times 10^{-9}$	Di 76
$K_L \rightarrow \mu^+ e^- \text{ or } \mu^- e^+$	$< 2 \times 10^{-9}$	Cl 71

If we assume that the strange and (anti)down quarks in the kaon annihilate to form a virtual massive family boson that emits the final observed muon and electron, the rate (Γ) for this decay process will be given (roughly) by

$$\Gamma \sim \frac{g^4}{M_F^4} M_K^5. \quad (1)$$

The fourth power of M_F appears just as M_W and M_Z do in ordinary weak processes (in G_F^2 , the Fermi constant squared). Our GUT prejudice appears in the assumption that the "family coupling constant" is similar to that for the weak (g) and electromagnetic interactions.

The branching ratio limit in Table III can be used in Eq. (1) to obtain a lower bound on M_F :*

$$M_F \gtrsim 10^5 \text{ GeV}. \quad (2)$$

The scale in Eq. (2) is not directly accessible by high energy accelerators in the foreseeable future. The Superconducting Super Collider (SSC), which is presently being considered for construction next decade, will reach 10^4 GeV and is estimated to cost several billion dollars. We cannot expect to be able to do something yet ten times larger for a long time.

Our experimental knowledge of the scale in Eq. (2), however, does not arise from such a brute force approach but rather from a precise measurement at a level of a few parts in 10^9 . Since Γ varies as M_F^{-4} , it is necessary to improve the present branching ratio limit by four orders of magnitude to search for a value of M_F in the 10^3 TeV (10^6 GeV) range. This will be quite feasible at LAMPF II where kaon fluxes on the order of 10^8 /sec will be produced. (Raw fluxes would be larger but some flux will be traded for higher quality beams to reduce backgrounds and minimize systematic errors.) A typical solid angle times efficiency factor for an inflight decay experiment is on the order of 10%.

*Certain chirality properties of the family interaction could require that two of the five powers of M_K in Eq. (1) be replaced by the muon mass, m_μ . However, after taking a fourth-root to obtain Eq. (2), there is little difference.

Thus, 10^7 K's per second could be examined for the decay mode of interest. A branching ratio larger than 10^{-12} could be found in one day of running. A full year of beam time would be sensitive down to the 10^{-14} level. Of course, we do not know if a positive signal will be found at even this tiny level. Nonetheless, the need for such an observation to elucidate family dynamics drives us to make the attempt.

We must also investigate the possibility that family symmetry is governed by a massless scalar boson, which has been called the familon. As is generally true for such scalars, the strength of its coupling falls inversely with the mass scale (F) at which the family symmetry is spontaneously broken. Cosmological arguments suggest a lower bound on F very near to its upper bound from laboratory experiments:

$$10^9 \text{ GeV} < F < 10^{12} \text{ GeV}. \quad (3)$$

The familon should appear in the two-body decays such as $\mu \rightarrow e + f$ or $s \rightarrow d + f$. The latter appears in $K^+ \rightarrow \pi^+ + f$, where the familon is unseen and the π^+ appears at a momentum of 227 MeV/c. This is a special case of the class of experiments called $K \rightarrow \pi + \text{nothing}$. Improved searches for the familon will require very precise and high resolution measurements of meson-decay spectra.

The neutrino sector offers another arena to search for evidence of nonconservation of family quantum numbers. Neutrino radiative decay has been sought in ν_μ beams with the result of a bound on the lifetime of $\tau_\nu > 10^5 \text{ sec} (m_\nu/\text{MeV})$ for decay into a lighter neutrino plus a photon. If the muon family were not distinct from any other, the naive expectation for τ would be $10^3 \text{ sec} (\text{MeV}/m_\nu)^5$. Better evidence may be found in neutrino scattering experiments. The decay of positive muons produces muon-family anti-neutrinos and electron-family neutrinos as long as family quantum numbers are preserved. If these neutrinos are subsequently scattered on nuclear targets, they should produce only positive muons and electrons, respectively. A LAMPF experiment has confirmed that no negative muons or positrons are produced at a level of about 5%. Since positive pions decay overwhelmingly ($10^4:1$) into positive muons and muon-family neutrinos, subsequent scattering of these neutrinos should produce only negative muons; this, too, has been accurately confirmed.

From the above, it is clear that at present no evidence exists in any (electric charge) neutral process mediated by the Z^0 boson for a nonconservation of a family quantum number, i.e., neutral family-changing interactions. Is it possible that these quantum numbers are exactly preserved? Are the family superselection rules as inviolate as those for electric charge and angular momentum? The answer to this is unequivocally NO! We know beyond a shadow of a doubt that family must be a broken symmetry. To see this, one needs only to examine the mass spectrum of the fermions in Table II.

All of the fermion masses violate the electroweak symmetry. However, the pattern of mass splittings within each family and between families all show that family symmetry is also broken. More importantly, because we do not know the mass scale or understand the pattern of the family symmetry breaking, we also cannot determine the mass scale for the electroweak breaking in the fermion sector. If the interactions responsible for family symmetry breaking subtract from the electroweak masses, the scale of the latter could be consistent with the W and Z masses.

Other evidence for family symmetry breaking can be found in weak decays mediated by the charged bosons W^+ and W^- . We know from π and K decays that both the strange quark s and the down quark d decay into the up quark but with significantly different strengths; the down quark has the greater strength. Cabibbo described this by saying (in modern terms) that in the charged weak current, the up quark couples to the "rotated" combination $d' \equiv (d \cos \theta_c + s \sin \theta_c)$ of down and strange quarks. Today we express this fact by saying that the eigenstates of the weak interaction (d', s') are rotated combinations of the eigenstates (d, s) of the quark mass matrix:

$$\begin{pmatrix} d' \\ s' \end{pmatrix} = \begin{pmatrix} \cos \theta_c & \sin \theta_c \\ -\sin \theta_c & \cos \theta_c \end{pmatrix} \begin{pmatrix} d \\ s \end{pmatrix}$$

where the Cabibbo angle, θ_c , is experimentally found to be $\sin \theta_c = 0.23 \pm 0.01$. However, the difference between the mass and weak-interaction eigenstates shows that the full Hamiltonian (of the world) cannot be diagonalized in a manner that defines conserved family quantum numbers.

The generalization of this formalism to the three-family case produces an important relation between family nonconservation and CP nonconservation (time reversal noninvariance). The generalized form, due to Kobayashi and Maskawa (KM), is a (3×3) unitary matrix, which mixes the d , s , and b quarks and includes a phase δ , which allows CP-nonconserving effects to occur. If family symmetry were exact, the angles in the KM matrix would vanish; then the phase δ would be absorbed in the b -quark field and could not contribute to the observed CP nonconservation.

Returning to the discussion of fermion masses, the neutrino masses are especially interesting, as they appear to be very small. The question of neutrino mass is complicated by their lack of electric charge and a peculiarity of the CPT theorem. The latter guarantees the existence of a spin-down (negative helicity) anti-particle state as a partner for every spin-up (positive helicity) particle state but it does not require a spin-down particle state. (We choose the spin quantization axis parallel to the momentum.) When the particles are massless, we can speak of positive and negative helicity or, equivalently, chirality, which just refers to the handedness (as in the right hand rule for the Poynting vector) of the combination of spin rotation and linear momentum. However, helicity is not a Lorentz invariant quantity for a massive state, as the particle may be brought to rest and then accelerated in the opposite direction thus changing the sign of its helicity. The chirality, on the other hand, is a Lorentz invariant. Thus, if a particle has mass it requires that either an additional spin down particle state exists (Dirac case) or that particle and antiparticle are indistinguishable (Majorana case).

The latter is impossible when there is a conserved charge, such as electric charge, which distinguishes the particle and antiparticle. Neutrinos, however, carry only weak charge, which is not conserved due to the spontaneous breaking of electroweak symmetry. Indistinguishability of the neutrino and its antineutrino means that lepton number is not conserved; the mass term reflecting this indistinguishability is termed a Majorana mass. Because of the possibility of this Majorana mass in addition to the standard Dirac mass, the neutrino mass matrix is more complicated than that for the other fermions.

Thus the question of neutrino mass is very important and different types of searches are being undertaken. One class of experiment looks for massive neutrinos in decays such as $\pi \rightarrow \mu \nu$ and $K \rightarrow \mu \nu$; the effect appears as a small spike in the lepton energy distribution above the background of multibody decays. Another possibility is to detect a distortion of the π^+ spectrum in $K^+ \rightarrow \pi^+ \nu_i \bar{\nu}_1$ caused by the closing of phase space due to the ν_1 mass. This method is sensitive to neutrino masses in the range 50 to 150 MeV/c², independent of the overall rate normalization. (De 84) So far, the only positive report of a neutrino mass comes from the tritium decay experiment of Lubimov et al. (Lu 83); they report $m(\nu_e) \approx 20\text{--}45 \text{ eV}/c^2$.

Another consequence of massive neutrinos is the possibility of neutrino oscillations. If the neutrino types have mass and mix in the same general manner as quarks, beams of neutrinos will oscillate from one type into another as one moves away from the neutrino source, just as occurs in the $K^0 - \bar{K}^0$ system. The probability of finding the neutrino state $|b\rangle$ orthogonal to $|a\rangle$ at a distance x is:

$$P_{a \rightarrow b} = \sin^2(2\theta) \sin^2\left(\frac{\pi x}{L}\right) \quad (4)$$

for relativistic neutrinos with mixing angle θ , where the neutrino oscillation length, L , is given by

$$L \approx 2.5m \left(\frac{E}{\text{MeV}}\right) \left(\frac{\Delta m^2}{\text{eV}^2}\right)^{-1} \quad (5)$$

and $\Delta m^2 = m_a^2 - m_b^2$.

Although the general case of three neutrino types is more involved, Eqs. (4) and (5) already contain the generally important implications for experiments that attempt to discover neutrino mass effects in this way. From Eq. (4), it is clear that the largest signal is obtained at $x = L/2$. Even there, the signal is small if θ is small. Thus, experiments require many neutrino scattering events and as small a background as possible for maximum sensitivity to small values of θ . However, the flux of neutrinos always falls off as x increases for any fixed size detector due to the declining solid angle subtended. Thus, for a given Δm^2 , the smallest value of E (and so, L) is desirable. Unfortunately, neutrino cross sections fall rapidly with decreasing energy. For any observation process with an energy threshold, the optimum energy will be not far above that threshold. For example, to detect muon neutrinos, they must have sufficient energy to produce muons. Finally, as for small θ , small Δm^2 reduces the signal strength at any given distance, x . This again puts a premium on maximizing the number of detectable events using the high flux from LAMPF II.

3.1.3 Precision Tests of the Standard Model

The questions about family have been considered in the context of possible extensions to the minimal standard model. Even if we were completely convinced of the validity of the standard model it would be necessary to subject it to the most detailed tests possible. In fact it is widely believed that the standard model is inadequate and that extensions will be required. In this situation, a new generation of precision experiments at LAMPF II should give us clues as to the direction for extensions.

As mentioned earlier, the standard model consists of the $SU(3)_C$ gauge field theory called quantum chromodynamics and the spontaneously broken $SU(2)_W \times U(1)_B$ gauge field theory known as the "electroweak model". These theories define spin 1 (vector) gauge fields, which describe the transmission of the strong and weak forces, just as the photon transmits the electromagnetic force. The eight massless "gluons" fill an adjoint representation of the $SU(3)_C$ group, and the three massive intermediate vector bosons of the weak force, the W^+ , W^- , and Z , fill the adjoint representation of the $SU(2)_W$ group.

The Z and the photon are actually orthogonal mixtures of a component from that adjoint and the single vector boson associated with the $U(1)_B$ gauge group. This mixture is characterized by the weak angle, θ_W . The relation between the mass of the W and Z and θ_W is $M_W^2 = M_Z^2 \cos^2 \theta_W$. In addition to these vector bosons, the standard model includes several fermions and one scalar field called the Higgs boson. This last is an as yet unobserved remnant of the spontaneous symmetry breaking, which is the technical means used to generate masses for the W and Z bosons.

Each of the fermion families shown in Table I has the following form: it includes one left-chiral (i.e., negative helicity, for a massless or highly relativistic particle) neutral lepton, a neutrino, in an $SU(2)_W$ -doublet with a left-chiral projection of a charged lepton. The right-chiral projection of this charged lepton appears as a separate $SU(2)_W$ -singlet. All of these leptons are $SU(3)_C$ -singlets. There are also charge $+2/3$ generically "up," and $-1/3$ generically "down," quarks. These appear in $SU(3)_C$ -triplets with their right-chiral projections in $SU(2)_W$ -singlets as for the charged lepton. The left-chiral projections of these up and down quarks form an $SU(2)_W$ -doublet.

This gives a complete accounting of the known "elementary" fermions and their electroweak interactions. Detailed tests of the predictions and parameters of the standard model are now essential to demonstrate its correctness. Three such tests are highlighted here. Perhaps the most outstanding precision test is $\nu_\mu e$ scattering. It would allow a clean, high precision measurement of θ_W , which would test the electroweak theory at the one-loop level; this is a stringent test, analogous to the Lamb shift or $g-2$ for quantum electrodynamics.

The problem of CP nonconservation is one of the great mysteries of our generation. It has been seen only in the K system and the experimental evidence does not dictate a unique theoretical interpretation. One of the strengths of the standard model is a natural parameterization of CP nonconservation. It is very important, therefore, to make measurements with adequate precision to see if this represents the proper explanation.

In principle, CP nonconservation can be studied in other, heavier quark systems, namely $D^0 (=c\bar{u}) - \bar{D}^0$, $B^0 (=b\bar{d}) - \bar{B}^0$, and $T^0 (=t\bar{u}) - \bar{T}^0$, but there are two obstacles to this. One problem is that the CP-nonconserving effects arise only in communicating channels between particle and antiparticle, which are a smaller fraction here of the total decay amplitude than for kaons. The K^0 and $\bar{K}^0$ both largely decay to two and three pions, so that unitarity promotes $K^0 - \bar{K}^0$ mixing. Such modes are Cabibbo (or KM) suppressed for the heavier mesons, and the dominant channels do not communicate to this order in the weak interaction; e.g., $D^0 = c\bar{u} \rightarrow s\bar{u}$, while $\bar{D}^0 = \bar{c}u \rightarrow \bar{s}u$ and the s and $\bar{s}$ do not mix under the strong interaction. The other problem is a small production rate, $< 10^9/\text{yr}$ vs.

the 10^9 K's/sec possible at LAMPF II. Thus, one must search for a smaller fractional effect in a smaller available data set. Clearly, the strange quark offers the better window on family symmetry and CP nonconservation.

The process $K^+ \rightarrow \pi^+ + \text{nothing}$ else seen was introduced as a way to search for the familon. It is also an important test of the standard model via the rate for the allowed mode $K^+ \rightarrow \pi^+ + \nu_i + \bar{\nu}_i$ for all light neutrino types, i. This process occurs through a one-loop quantum field theoretic correction to the standard electroweak theory. It is interesting in itself for two reasons: 1) it depends on quark mixing differently from CP nonconservation and so allows independent study of the quark mixing matrix elements; 2) it depends on the number of light neutrino types, N_ν . Because the latter number is expected to be determined in Z^0 decay studies, an uncertainty in the hadronic overlap matrix element can be determined. Present estimates place the branching ratio in the range $N_\nu \times (10^{-9} - 10^{-10})$. If this experiment yielded a discrepancy with the N_ν value determined from Z^0 decay it would be evidence for new physics or that at least one of the neutrinos has a mass greater than about 200 MeV/c².

3.1.4 Summary

Here we have addressed the unresolved problem of family symmetry and how a high-intensity medium energy machine, LAMPF II, can effectively attack that problem. The search for answers to the family problem and the process of testing the detailed predictions of the standard model may lead to surprises. In addition, one of the important functions of the LAMPF II accelerator will be to make explicit searches for new phenomena. The fact that LAMPF II will be a facility with a broad range of beams will be an asset in responding to new developments that we cannot foresee. However, we can already identify some experiments that are strong candidates for revealing new physics. Those that have already been mentioned above include the decay $K \rightarrow \pi + \text{nothing}$ and CP-nonconservation experiments. Also, new light particles may appear in a variety of rare processes and evidence for new heavy particles may be revealed in decays such as the family nonconserving decays.

These experiments will require a variety of high quality beams to be available at LAMPF II, just as muon beams at LAMPF have enabled the world-leading searches for family nonconservation in the decays $\mu \rightarrow e\gamma$ and $\mu \rightarrow e\gamma\gamma$. Of particular importance are the high flux and high quality stopping kaon beams and the high flux and well defined multi-GeV kaon beams. The extra flux at LAMPF II also provides unique opportunities to obtain desirable experimental conditions for other meson decay studies. For example, a momentum-analyzed K^+ beam could be used to produce, by charge exchange scattering, a narrow momentum bite neutral kaon beam with a flux in the latter comparable to present day unanalyzed beams. Similarly, a narrow-band neutrino beam at LAMPF II could equal or exceed the flux in present wide-band beams.

For neutrino oscillation experiments, the high flux of neutrinos at modest energies allows maximization of the statistical power of experiments while minimizing backgrounds. The LAMPF II complex can provide high flux, low-energy ν_μ beams with exceptionally little ν_e contamination for experiments searching for $\nu_\mu \rightarrow \nu_e$, clean beams at the intermediate energies optimum for studying ν_μ disappearance, and unparalleled intensities of ν_e and $\bar{\nu}_e$ at energies suitable for studying ν_μ appearance.

In summary, the family symmetry problem is as fundamental and important a problem as Grand Unification, and it may well be a completely independent one. The known bound of 100 TeV on the scale of family dynamics is an order of magnitude beyond the direct reach of even foreseeable accelerators, such as the Superconducting Super Collider. These dynamics, however, may be accessible in studies of rare kaon decays, meson and hyperon decays, neutrino oscillations, and careful studies of CP nonconservation. Thus the high-precision frontier is complementary to the high-energy frontier. To undertake these experiments at the required levels of precision requires intense fluxes of particles from the second or later families. Only particles from the second family can be produced at the required intensities. The LAMPF II high intensity, medium-energy accelerator is a highly cost-effective means to meet these experimental needs.

In the remainder of this chapter specific experiments will be discussed. These are organized by whether they provide keys to the family problem, precision tests of the standard model, or whether they search for new physics. Of course the experiments given are only a sample. It must be emphasized that LAMPF II is a facility that will provide a broad range of probes that will allow us to respond to developments in physics that we cannot anticipate at this time.

3.2 KEYS TO THE FAMILY PROBLEM

As emphasized in the introduction, there already exists significant evidence for family symmetry breaking in the quark mass spectrum, and in transitions between quarks mediated by the charged weak bosons W^+ and W^- . Keys to the family problem must now be sought in the properties and interactions of leptons. The two principal areas of investigation are the neutrino mass matrix and the breakdown of lepton-family conservation in charged lepton processes. Either of these areas may provide the necessary insight to solve the family problem.

3.2.1 Searches for Massive Neutrinos

The study of the neutrino mass matrix is of great importance to physics beyond the standard model. Neutrino masses and mixings offer a unique window on sectors of the fundamental theory that are not directly accessible through charged particles. High-intensity neutrino beams of modest energies offer experimental advantages for studying the expected ranges of masses and mixing parameters.

It has been known for a long time that the charged weak current mixes quarks of different generations: the strange quark decays into the up quark with an amplitude proportional to $\sin\theta_c$, where θ_c is the famous Cabibbo angle. On general grounds such as quark-lepton symmetry and grand unified theories that mix quarks and leptons in the same multiplets, we expect that a similar situation will hold for leptons: the charged weak current will mix the generations because the weak flavor eigenstates are not mass eigenstates. This means that at least one neutrino must have a nonzero mass and that flavor eigenstates, for example ν_f , must be linear superpositions of mass eigenstates ν_1, ν_2, ν_3 :

$$\nu_f = \sum_{i=1}^3 U_{fi} \nu_i .$$

However, when we come to consider the neutrino mixing matrix and mass eigenvalues in detail, we find that we have no definitive theoretical guidance. Numerous general arguments, as well as those within specific models, suggest that the lepton mixing matrix should be similar, or even identical to, the quark mixing matrix. However, these arguments are not compelling. Likewise there are attractive arguments suggesting that neutrino masses should be small ($\sim 10^{-6} - 1$ eV), but there are no firm predictions about them. Thus we must make a search over the entire range of masses and mixing parameters.

The present information on ν -mass measurements has been summarized in Table II. Prospects for improvements at LAMPF II are discussed following a review of neutrino oscillations, which are a less direct way of searching for neutrino mass, but one that also gives information about mixing parameters.

3.2.1.1 Neutrino Oscillations. Suppose that ν_e and ν_μ are linear superpositions of two mass eigenstates ν_1 and ν_2 ,

$$\nu_e = \nu_1 \cos\theta + \nu_2 \sin\theta \quad \nu_\mu = \nu_1 \sin\theta - \nu_2 \cos\theta .$$

As the muon neutrino propagates in time, the phase relationship between ν_1 and

ν_2 will change unless ν_1 and ν_2 are perfectly degenerate in mass. This means that an original ν_μ state will develop a new admixture corresponding to the flavor ν_e and will suffer a diminution in the component of the initial flavor. The mere observation of either the new flavor ν_e ("appearance" experiment) or the diminution of the original flavor ν_μ ("disappearance" experiment) would be sufficient to establish the qualitative result that ν_1 and ν_2 must have different masses and that at least one mass must be different from zero.

The probability for the appearance oscillation $\nu_\mu \rightarrow \nu_e$ is given by

$$P(\nu_\mu \rightarrow \nu_e) = \sin^2 2\theta \cdot \sin^2\left(\frac{\pi x}{L}\right), \quad L = 2.5 \frac{E_\nu}{\Delta m^2}$$

where $\Delta m^2 = m_1^2 - m_2^2$ measured in $(\text{eV})^2$, θ is the mixing angle given above, x is the flight path measured in meters, and E_ν is the neutrino energy in MeV. The probability for the disappearance oscillation $\nu_\mu \rightarrow \nu_\mu$ is

$$P(\nu_\mu \rightarrow \nu_\mu) = 1 - P(\nu_\mu \rightarrow \nu_e) .$$

In principle, we should consider three flavors and three mass eigenstates; but the corresponding analysis is much more complicated than in the two-state case, and it has become customary to analyze data in terms of the simpler case.

From the expression for $P(\nu_\mu \rightarrow \nu_e)$, we see that the figure of merit determining the sensitivity of an experiment to Δm^2 is the ratio $E_\nu(\text{MeV})/x(\text{meters})$ or equivalently $E_\nu(\text{GeV})/L(\text{km})$. Typical figures of merit for various sources of neutrino are

Quantity	Source of ν			
	Solar	Cosmic Rays	Reactors	Accelerators
E_ν/L	$10^{-10} - 10^{-11}$	10^{-4}	10^{-1}	$1 - 100$
$\Delta m^2(\text{eV})^2$	$10^{-11} - 10^{-12}$	10^{-4}	10^{-2}	$10^{-1} - 10^3$

In the limit of very large Δm^2 , the appropriate oscillating factor $\sin^2(\pi x/L)$ can be replaced by its average value of $1/2$, and experiments, especially appearance ones, become very sensitive to small values of $\sin^2 2\theta$. Limits of order 10^{-2} for $\sin^2 2\theta$ have been achieved in some experiments, and a reduction to 10^{-3} is expected in others. The current status of limits on Δm^2 and $\sin^2 2\theta$ from a variety of experiments of both the appearance and the disappearance types is shown in Fig. 3.-1. Proposed limits for a facility at LAMPF II could be more than an order of magnitude more sensitive.

While the emphasis in future experiments is likely to be on smaller mixing angles and smaller mass differences, we should mention the claim by a group working at Le Bugey reactor (Ko 84) to have seen an "indication of oscillations" with parameters $\sin^2 2\theta \approx 0.25$ and $\Delta m^2 = 0.2(\text{eV})^2$ in ν_e disappearance. This is in contradiction to the published limits from the Gösgen reactor (Vu 82). It is by no means clear at this time that the hint at Le Bugey is correct, but if it should be, the energy range and intensity available at LAMPF II should make it an ideal place to explore and exploit this finding.

Fig. 3.-1 Present experimental limits on neutrino mixing parameters.

Appearance experiments gain from high intensity and relatively low energy. Disappearance experiments are more difficult to perform because of systematic uncertainties in the knowledge of the beam characteristics. The main problem with these experiments is that the dominant source of variation of interaction rate with distance from the source comes from geometrical effects that are poorly known. The need for intense beams to yield many interactions is at variance with the loss of intensity that is implied by detailed knowledge of the production of the neutrino beam. The availability of intense proton beams for ν production will enable a more satisfactory compromise to be reached than is available at present.

There are a number of advantages that LAMPF II would have in conducting the next generation of neutrino-oscillation experiments. With such a low duty factor (3×10^{-6} s/pulse), these experiments will have very low cosmic-ray backgrounds. The extremely high flux allows a much increased counting rate, improved signal/noise, and the potential for direct measurement and removal of many neutrino-induced backgrounds (via use of the narrow-band beam). These advantages should provide greatly improved limits on neutrino oscillations.

3.2.1.2 Direct Neutrino Mass Measurements. Direct measurements of neutrino masses are complementary to neutrino-oscillation experiments. A sensitive class of experiments of this kind involves measuring the lepton momentum spectrum in two-body leptonic decays of pseudoscalar mesons (Sh 80,

Shr 81). The presence of massive ν states in addition to the dominant one would show up as additional peaks in the charged lepton momentum spectrum, with relative intensity proportional to $|U_{fi}|^2$. If the neutrinos of the dominant decay were massive the momentum of the charged lepton would be shifted down by $m_\nu^2/2M$.

In the decays $K_{e\nu}$ and $K_{\mu\nu}$ the relatively large mass of the kaon allows a wide range of possible ν masses to be explored. A recent KEK experiment (Ya 82, Ha 82) has obtained the limit on $|U_{fi}|^2$ versus neutrino mass shown in Fig. 3.-2. Increased numbers of events per mass bin would improve the limits on the mixing amplitudes. This may be possible at current machines with more sophisticated event analysis and larger acceptance spectrometers. It would be much easier at LAMPF II even using present-day techniques. There is also interest in improving the mass limit of the muon neutrino extracted from the $\pi \rightarrow \mu \nu_\mu$ reaction, sensitive only to m_{ν_μ} (Deu 80).

Fig. 3.-2. (a) Upper bounds of $|U_{\mu i}|^2$ versus neutrino mass. (b) Upper bounds of $|U_{e i}|^2$ versus neutrino mass made in (Ya 82), but originally from TRIUMF and CERN.

It was noticed recently (Deu 82) that muon capture and, in particular, a precision measurement of the triton recoil energy in the reaction $\mu^- + {}^3\text{He} \rightarrow {}^3\text{H} + \nu_\mu$ yield useful limits on eventual neutrino mixing in m_ν ranges that are not accessible to two-body meson decays. A precision measurement of the triton energies in the reaction $\mu^- + {}^6\text{Li} \rightarrow {}^3\text{H} + {}^3\text{H} + \nu_\mu$, near their endpoint, provides a sensitivity linear in m_{ν_μ} (Wi 71). A precision of 40 keV on the sum of the two triton energies would provide a precision (or limit) of 100 keV on m_{ν_μ} (to be compared to the actual limit of 265 keV!). Unfortunately, the corresponding region of the Dalitz plot is very scarcely populated ($\sim 10^{-10}$ /stop) and so requires very high stop rates. Moreover, to avoid a prohibitive energy loss of the tritons, the target thickness should be limited to ~ 10 mg/cm², requiring a stop density of 10^8 to 10^9 /g/s. An experiment such as this is only practical with the greatly enhanced low energy μ flux available at LAMPF II.

3.2.2 Family Nonconserving Decays

In the standard model of the electroweak interactions with a single Higgs doublet, family is conserved in the couplings of the Z^0 and of the Higgs meson. Still, (lepton-) family conservation would not be expected to hold, unless the neutrinos are massless (or degenerate). However, in view of the stringent experimental limits on the masses of ν_e , ν_μ , and ν_τ , the rates of family nonconserving K decays would be negligible* because of $(m_\nu/m_W)^2$ suppression. As a consequence, searches for family nonconserving processes probe the existence of new interactions and/or new leptons beyond those of the three-generation standard model. The possibilities include the existence of horizontal gauge interactions mediated by new flavor-changing neutral gauge bosons, existence of flavor-changing neutral Higgs bosons, family nonconservation mediated by leptoquarks, existence of further generations of leptons, and further electroweak interactions beyond the standard model.

3.2.2.1 Family Nonconserving Kaon Decays. The branching ratios of $K_L \rightarrow \mu^\pm e^\mp$ and $K^{(\pm)} \rightarrow \pi^{(\pm)} \mu^\pm e^\mp$ mediated by a horizontal gauge boson Y^0 associated with an Abelian horizontal gauge symmetry are given by (He 81)

$$B(K_L \rightarrow \mu e) \approx (9.4 \times 10^{-1}) \left[\frac{1}{m_Y(\text{TeV})} \right]^4 \tilde{g}^4 \times (\text{Re}\beta_A)^2 (|\sigma_V|^2 + |\sigma_A|^2) \quad (6)$$

$$B[K^{(\pm)} \rightarrow \pi^{(\pm)} \mu^\pm e^\mp] \approx (9.6 \times 10^{-3}) \left[\frac{1}{m_Y(\text{TeV})} \right]^4 \tilde{g}^4 \times (\text{Re}\beta_V)^2 (|\sigma_V|^2 + |\sigma_A|^2) \quad (7)$$

where $\tilde{g}$ is the horizontal gauge coupling constant and m_Y is the mass of Y^0 ; $\tilde{g}_{A,V}^\beta$ and $\tilde{g}_{A,V}^\sigma$ are the effective couplings of Y^0 (the subscripts denote the type of coupling) to sd and μe , respectively. $\beta_{A,V}$ and $\sigma_{A,V}$ include the various mixing angles in the quark and leptonic sectors. Note that $K_L \rightarrow \mu e$ ($K^\pm \rightarrow \pi^\pm \mu e$) is sensitive to an axial vector (vector) quark current. The experimental limits yield lower bounds on m_Y of

$$m_Y \geq 147 [|\text{Re}\beta_A|^{1/2} (|\sigma_V|^2 + |\sigma_A|^2)^{1/4} \tilde{g}] \text{ TeV} \quad (\text{from } K_L \rightarrow \mu e), \text{ and} \quad (8)$$

$$m_Y \geq 34 [|\text{Re}\beta_V|^{1/2} (|\sigma_V|^2 + |\sigma_A|^2)^{1/4} \tilde{g}] \text{ TeV} \quad [\text{from } K^{(\pm)} \rightarrow \pi^{(\pm)} \mu e]. \quad (9)$$

*See, for example, the review in the Panel P-2 report, (Al 79).

The experimental value of the K_L - K_S mass difference (Δm_K) implies a bound on m_Y that, in general, is much more stringent than Eqs. (8) and (9) (Her 79, Ka 80). Nevertheless, branching ratios for $K_L \rightarrow \mu e$ of the order of 10^{-11} , and perhaps even larger, are possible, provided that the mixing angles in the quark sector are much smaller than are those in the leptonic sector (Her 79, Ka 80, He 81). For equal couplings $B(K(\pm) \rightarrow \pi(\pm)\mu e)$ would be smaller by a factor of 100 [cf. Eqs. (6) and (7)]. In models with non-Abelian horizontal gauge symmetries, the branching ratios for $K_L \rightarrow \mu^\pm e^\mp$ and $K^\pm \rightarrow \pi^\pm \mu^\pm e^\mp$ could be as large as the existing experimental limits, as mechanisms are possible (Ma 79, Ca 80, Le 60) that lead to cancellations among the contributions of the various horizontal gauge bosons in Δm_K [but not in $B(K_L \rightarrow \mu e)$, $B(K^\pm \rightarrow \pi^\pm \mu^\pm e^\mp)$].* It should be noted, however, that the existence of such a mechanism requires that some special conditions be satisfied in the model (He 81, Sh 81a, Sh 81b).

A search for family nonconservation via horizontal-gauge bosons in Σ^+ decay may have advantages over such a search in kaon decay because the decay does not require a coupling to a pseudoscalar current, which would be sensitive to helicity suppression (Ka 80). However, the limit on the possible branching ratio for $\Sigma^+ \rightarrow p \mu e$ decay for such a process, based on Δm_K and the current $K_L \rightarrow \mu e$ branching ratio, is only $BR(\Sigma^+ \rightarrow p \mu e) \lesssim 10^{-6}$. Such a large branching ratio is allowed by the poor knowledge of $BR(K_S \rightarrow \mu e) \lesssim 10^{-3}$ and requires an accidental arrangement of the coupling constants. The clean and intense beams obtainable at LAMPF II would allow a great improvement of this limit and a restriction of the possible form of the coupling.

In the standard model with more than one Higgs doublet, lepton-family nonconserving transitions may occur via Higgs-lepton couplings (Bj 77). Such transitions may also be mediated by Higgs bosons that are associated with models based on gauge groups beyond $SU(2)_L \times U(1)$. In general Δm_K imposes a stringent limit on the couplings and masses of neutral strangeness-changing Higgs bosons (Her 79). Still, family nonconserving Higgs exchange may give rise to $K_L \rightarrow \mu e$ and $K^\pm \rightarrow \pi^\pm \mu e$ rates comparable to the present experimental limits, provided that the ratio of the effective Higgs-quark and Higgs-lepton couplings (which include the mixing angles) is small (He 81, Her 79, Mc 81, Sha 82).

Family conservation could also be violated through the exchange of leptoquarks (bosons causing quark-lepton transitions). Leptoquarks appear in theories that unify the strong and the flavor interactions. In a subclass of models due to Pati and Salam based on $\otimes [SU(2n)]^4$ ($n > 3$), they can be sufficiently light to cause observable effects in K decays (Pat 82). As leptoquark exchange does not generate a nonleptonic interaction in lowest order, leptoquark masses are not severely constrained by Δm_K . Consequently, the rates for $K_L \rightarrow \mu e$ and $K^\pm \rightarrow \pi^\pm \mu e$ are allowed to be comparable to the present experimental limits (Sha 81). Calculation of the unification scales to two loops (Go 81) in these models leads to the conclusion that for $n = 4, 5$, family nonconserving K decays may have measurable rates, while for $n = 4$, and perhaps also for $n = 5$, the proton lifetime (τ_p) and the neutron-antineutron oscillation time ($\tau_{n\bar{n}}$) are both much larger than can be measured in the foreseeable future (Go 81). The reason for this difference is that family nonconserving K decays are first-order processes in the leptoquark-induced hypoweak interaction whereas

*The rates for $K^\pm \rightarrow \pi^\pm \mu^\mp e^\pm$ are expected to be suppressed in such models. See R. N. Cahn and H. Harari, (Ma 79, Ca 80, Le 60).

the other two processes are third-order hypoweak. [The $n = 6$ case is already in conflict with the limits for τ_p and $B(K_L \rightarrow \mu e)$, while for $n = 3$, all of the processes are unobservably small.]

Relatively light leptoquarks also appear in extended technicolor models* and are expected to give rise to $K_L \rightarrow \mu e$ with a branching ratio near the present experimental limit, and to $K^\pm \rightarrow \pi^\pm \mu e$ with a branching ratio of the order of $10^{-2} B(K_L \rightarrow \mu e)$. We note that in models involving heavy neutrinos, the rates for $K_L \rightarrow \mu e$ and $K^\pm \rightarrow \pi^\pm \mu e$ could be large even in the absence of a direct family nonconserving interaction. An example is the extended electroweak model (Ri 81) based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ or possibly the standard model with additional lepton families.

3.2.2.2 Other Decays and Processes. The current experimental limits for rare neutrinoless muon decays are presented in Table III. The Crystal Box experiment currently underway at LAMPF hopes to push the limits for $\mu^+ \rightarrow e^+ \gamma$ and $\mu^+ \rightarrow e^+ \gamma \gamma$ to the level of 10^{-11} . A future LAMPF $\mu^+ \rightarrow e^+ \gamma$ experiment will try to reach a level below 10^{-13} . With the more intense surface muon beams potentially available at LAMPF II ($>10^8/s$) one might hope eventually for a limit of 10^{-15} for the $\mu^+ \rightarrow e^+ \gamma$ decay process (Co 83).

Closely related to these rare decays is the $\mu^- Z \rightarrow e^- Z$ conversion experiment (Br 84). The current limit for this reaction is $<2 \times 10^{-11}$. To significantly improve upon this, one requires very low beam-associated backgrounds coupled with a high μ^- flux ($\approx 10^7$ stops/s) at low (≈ 60 -MeV/c) momentum. This could be achieved using the short pulse mode of LAMPF II whereby one could obtain pulses with a width of 1.1 μs and an average stopping rate in excess of $10^7/s$ in a 250-mg/cm² target.

The conversion of muonium into antimuonium also tests family number conservation. Because the initial and final states are degenerate, the conversion amplitude grows linearly in time, yielding an expected branching ratio of about 10^{-5} for Fermi coupling. The experiment must be conducted so that interactions with magnetic fields or from collisions with other atoms do not lift this degeneracy. The production of muonium in vacuum is a crucial ingredient in this experiment. The use of the 1.1- μs pulse mode of LAMPF II (and possibly of a very short pulse beam structure--a few ns--if available) would be a major advantage for muonium-antimuonium searches as the timing information would provide a useful way to identify and track the muonium state.

To summarize, searches for family nonconserving decays probe physics beyond the three-generation minimal standard model of the electroweak interaction. Present theoretical schemes suggest many possible sources of family nonconservation. In different models various family nonconserving decays are the most sensitive probes of such interactions. The branching ratio of a family nonconserving decay could be near the existing experimental limit. Therefore, searches for family nonconserving decays with increased sensitivity are of great importance and may have an enormous impact on our understanding of the fundamental interactions.

*Phenomenological implications of technicolor schemes and some associated problems of these theories are discussed in (El 81, Di 81).

3.3 PRECISION TESTS OF THE STANDARD MODEL

LAMPF II will provide beams of neutrinos, muons, kaons, and hyperons with unprecedented intensities and beam quality. Previous discoveries made with these probes have played an essential role in the development of the standard model. For example, the discovery of the two-component neutrino and two varieties of neutrino in the 1950's helped lay the foundations of modern gauge theories, and the discovery of neutral currents in the 1970's established the physical relevance of these theories. Deep-inelastic scattering of neutrinos proved to be a valuable probe of hadronic structure and played an important role in the development of the quark-parton model. The observation of 2π and 3π decay modes of the charged kaons led to the discovery of parity nonconservation in the weak interactions. The breakdown of CP symmetry was found in the 2π decays of the K_L . More recently, the suppression of decays such as $K_L \rightarrow \mu^+\mu^-$ led to the prediction of the charmed quark (Gl 70, We 72). [Subsequently, the order of magnitude of the charmed quark mass was successfully predicted (Va 73, Ga 74, Ga 76) from the observed K_L - K_S mass difference.]

The discussion of precision tests of the standard model has been divided into sections corresponding to pure-leptonic and semileptonic weak processes. Pure-leptonic weak interactions are shown to determine many of the basic parameters of the model, while the understanding of semileptonic weak interactions must include the corrections required in the presence of strong interactions. Tests of the CVC hypothesis and tests of QED are also presented. High-precision tests of the standard model will either confirm our ability accurately to calculate higher-order corrections to lowest-order processes, or they will reveal deviations that open a window onto new physics.

3.3.1 Leptonic Weak Currents

We begin with a discussion of neutrino-electron scattering, a fundamental interaction between two point-like particles, which is expected to proceed through both charged and neutral currents. An accurate study of the differential cross section will enable us to probe the basic properties of the weak interactions in the low-energy regime. These theoretically clean reactions offer an unparalleled opportunity for stringent tests of the radiative corrections to the standard electroweak theory.

The basic reactions that can be studied at LAMPF II include $\nu_\mu e \rightarrow \nu_\mu e$, $\bar{\nu}_\mu e \rightarrow \bar{\nu}_\mu e$, and $\nu_e e \rightarrow \nu_e e$. The first two reactions only proceed through neutral currents (Z^0 exchange), whereas the third reaction can proceed through charged currents as well as through neutral currents. Detection of coherent interference between the charged- and neutral-current diagrams will demonstrate the identity of scattered and incident neutrinos, and a measurement of both the magnitude and sign of the interference will provide a test of the standard model as well as a possible indication for more than one Z^0 gauge boson.

The differential cross section for ν -e scattering has the general form

$$\frac{d\sigma}{dy}(\nu-e) = \frac{G^2 M_e E}{2\pi} [A + 2B(1-y) + C(1-y)^2], \quad y = E_e/E_\nu, \quad (10)$$

where y is the ratio of the scattered electron kinetic energy to the incident neutrino energy. For antineutrino-electron scattering, the differential cross section has the same form except that the roles of A and C are interchanged.

The coefficient A arises from the (V - A) part of the effective Hamiltonian, and since charged-current reactions are believed to be pure (V - A), it is the term that will show the interference between charged and neutral currents in ν_e -e scattering. B arises solely from the nonstandard interactions S, P, and T, and so the discovery of a nonzero value for this coefficient would demonstrate the presence of such interactions; should it also turn out that B is positive, then the neutrino must be a Dirac particle. C is the (V + A) term, and to the extent that charged currents appear not to involve (V + A), it is a pure neutral-current term; as such it can be used to test μ -e universality in ν_e and ν_μ neutral-current interactions.

In the standard model, the coefficient B vanishes for all processes and the coefficients A and C can be expressed in terms of the mixing angle θ_W and the ratio $\rho = (M_W^2/M_Z^2) \cos^2 \theta_W$, which is equal to 1 when the Higgs particles in the theory form weak isodoublets. The relevant expressions are

	A	C
$\nu_\mu e$	$\rho^2(2\sin^2\theta_W - 1)^2$	$\rho^2(2\sin^2\theta_W)^2$
$\bar{\nu}_\mu e$	$\rho^2(2\sin^2\theta_W)^2$	$\rho^2(2\sin^2\theta_W - 1)^2$
$\nu_e e$	$[2 + \rho(2\sin^2\theta_W - 1)]^2$	$\rho^2(2\sin^2\theta_W)^2$
$\bar{\nu}_e e$	$\rho^2(2\sin^2\theta_W)^2$	$[2 + \rho(2\sin^2\theta_W - 1)]^2$

An accurate measurement of the ratio of A to C will yield an accurate value for the electroweak mixing angle $\sin^2\theta_W$, whereas absolute values of these coefficients will yield an accurate value of the parameter ρ , which measures the relative strengths of charged- and neutral-current interactions. Comparison of the values of $\sin^2\theta_W$ and ρ as determined from relatively low-energy leptonic reactions with those obtained from reactions involving hadrons, especially the $\bar{p}p$ and e^+e^- experiments on W and Z, will provide an important test of the standard model and an even more important test of the radiative corrections to it. It is the successful measurement of radiative corrections that will establish that we do indeed have a renormalizable theory of electroweak interactions.

The masses of W and Z are related to the mixing angle by

$$M_W = K/\sin\theta_W ; M_Z = 2K/\sin 2\theta_W$$

$$K = \left[\frac{\pi\alpha}{\sqrt{2}G_\mu} \right]^{1/2} \frac{1}{[1 - \Delta r]^{1/2}} = \frac{37.2810 \pm 0.0003}{[1 - \Delta r]^{1/2}} \text{ GeV} , \quad (11)$$

where α is the fine-structure constant, G_μ the muon decay constant, and Δr is the entire radiative correction to order α in the standard model ($\rho = 1$). Taking $\alpha^{-1} = 137.035963 (\pm 15)$ from the Josephson effect, $G_\mu = (1.16634 \pm 0.00002) \times 10^{-5} \text{ GeV}^{-2}$ from the muon lifetime, $\sin^2\theta_W = 0.217$ (the 1981 central value), a top-quark mass of 36 GeV, and a Higgs mass equal to that of the Z boson, the theoretical value of the radiative correction is approximately 7%.

$$\Delta r = 0.0696 \pm 0.002 , \quad (12)$$

and it has the effect of adding approximately 3 GeV to the predicted masses of

the W and Z bosons. Present results from the UA1 and UA2 collaborations are compared with theory ($\sin^2\theta_W = 0.217 \pm 0.014$) below.

	UA1	UA2	Theory
M_W (GeV)	$80.9 \pm 1.5 \pm 2.4$	$83.1 \pm 1.9 \pm 1.3$	$83.0 \begin{smallmatrix} + 2.9 \\ - 2.7 \end{smallmatrix}$ [80.1 if $\Delta r = 0$]
M_Z (GeV)	93.9 ± 2.9	$92.7 \pm 1.0 \pm 1.4$	$93.8 \begin{smallmatrix} + 2.4 \\ - 2.2 \end{smallmatrix}$ [90.5 if $\Delta r = 0$]

The most accurate determination of $\sin^2\theta_W$ to date,

$$\sin^2\theta_W = 0.224 \pm 0.007 \pm 0.006, \quad (13)$$

and the demonstration that ρ is either equal to 1 or extremely close to it, come from deep inelastic neutrino-nucleon neutral-current interactions. The value of $\sin^2\theta_W$ includes a theoretical as well as an experimental error because one must make assumptions about quark and antiquark distribution functions inside the nucleon. Ideally, we would like to determine $\sin^2\theta_W$ in a manner that is free of hadronic complications, and the clean process of neutrino-electron scattering, especially with muon neutrinos, provides us with such an opportunity. Thus we can look forward to the day when (1) M_W and M_Z are both determined to within ± 100 MeV; (2) ρ is shown to be precisely equal to 1; and (3) $\sin^2\theta_W$ is measured to an accuracy of 1% in $\nu_\mu e$ scattering. We will then be able to use the formulae for M_W and M_Z above to determine the value of Δr and compare it with theoretical calculations. Any discrepancy between theory and experiment will be a clear signal for new physics.

Experiments with $\sim 10^4$ events are needed to achieve the requisite 1% accuracy in the measurements of ν -e scattering (Ba 83). This requires an improvement in statistics of roughly two orders of magnitude as compared with present experiments, but it is a goal well worth achieving because ν -e scattering forms an important complement to e^+e^- experiments in the elucidation of the structure of neutral currents. As emphasized by White, (Wh 83) normalization is an important practical question in these experiments.

For the same 10^4 number of events, it will be more advantageous to do this experimentation at low energies. At present the signal/background is significantly better at the Brookhaven AGS than at CERN (Wh 83). The advantages of low energies include (1) a good e/γ separation and (2) the ability to remove background from coherent π^0 production. In addition, the possibility exists at low energy to measure angular distributions that yield independent information on the couplings. The actual rate for obtaining events is ~ 60 νe or ~ 30 $\bar{\nu} e$ per 10^{19} protons (about one day at LAMPF II).

3.3.2 Semileptonic Weak Currents and CVC

The standard model describes the electromagnetic and weak interactions of both quarks and leptons. In this section, processes are considered that involve weak interactions between quarks and leptons. The description of these

interactions is complicated by the exchange of gluons between the quarks. Precise experiments will test the ability of the theory to predict the corrections due to the presence of the strong interactions.

3.3.2.1 Neutrino-nucleus Scattering. Neutrino-nucleus scattering will be an effective probe of the structure of the semileptonic weak currents. The importance of the nucleus arises from the fact that because one can study transitions between states of definite spin, isospin, and parity, one has a filter for specific parts of the interaction. For example, as $q^2 \rightarrow 0$, transitions from the $J = T = 0^+$ ground state of C^{12} to the $J = T = 1^+$ excited state at 15.1 MeV can be engendered only by the "Gamow-Teller" isovector part of the hadronic neutral current. In the standard model this is, of course, the axial-vector interaction, and one can learn about the corresponding form factor from measurements of the appropriate cross section; in nonstandard models, this transition could also be induced by an isovector-tensor interaction, and so one can also set limits on the presence of such interactions. Because of the conserved-vector-current (CVC) hypothesis, one expects certain form factors extracted from neutrino reactions to be the same as those observed in electromagnetic reactions such as electron scattering. Thus one will be able to test the CVC hypothesis as a function of q^2 over a wide range of parameters. Other tests of the CVC hypothesis that will become possible at LAMPF II are kaon beta decay and muon capture.

First neutrino-nucleus scattering will be discussed. The processes under consideration are shown in Fig. 3.-3. The starting point in our analysis will be the effective Lagrangian of the standard model,

$$\begin{aligned}
 L_{\text{eff}} = & \frac{iG}{\sqrt{2}} \{ [\bar{e}\gamma_\lambda(1 + \gamma_5)\nu_e + (e \leftrightarrow \mu)] \tilde{J}_\lambda^{(+)} + [\bar{\nu}_e\gamma_\lambda(1 + \gamma_5)e + (e \leftrightarrow \mu)] \tilde{J}_\lambda^{(-)} \} \\
 & + \frac{iG}{\sqrt{2}} [\bar{\nu}_e\gamma_\lambda(1 + \gamma_5)\nu_e + (e \leftrightarrow \mu)] \tilde{J}_\lambda^{(0)} \\
 & - \frac{iG}{\sqrt{2}} [\bar{e}\gamma_\lambda(1 + \gamma_5)e - 4\sin^2\theta_w \bar{e}\gamma_\lambda e + (e \leftrightarrow \mu)] \tilde{J}_\lambda^{(0)} .
 \end{aligned} \tag{14}$$

The first two terms in the Lagrangian density govern charge-changing, semileptonic interactions. The third term describes elastic and inelastic neutrino scattering. The last term describes the neutral-current weak interactions of the charged leptons. (We assume $G^\pm = G\cos\theta_c \equiv G$.) The hadronic

Fig. 3.-3 Semileptonic electroweak processes with nuclei.

currents entering into these expressions have the following structure in the standard model.

$$\tilde{J}_\mu = J_\mu + J_{\mu 5} \quad V - A \quad (15)$$

$$\tilde{J}_\mu^{(\pm)} = \tilde{J}_\mu^{V1} \pm i\tilde{J}_\mu^{V2} \quad \text{Isovectors} \quad (16)$$

$$J_\mu^Y = J_\mu^5 + J_\mu^{V3} \quad \text{e.m. current} \quad (17)$$

$$J_\mu^{(\pm)} = J_\mu^{V1} \pm iJ_\mu^{V2} \quad \text{CVC} \quad (18)$$

$$\tilde{J}_\mu^{(0)} = \tilde{J}_\mu^{V3} - 2\sin^2\theta_W J_\mu^Y \quad \text{Standard model.} \quad (19)$$

Equation (15) says that the current has a vector axial-vector structure. Equation (16) states that the charge-raising and charge-lowering currents are isovector operators. The third relation gives the isotopic spin structure of the electromagnetic current, and the fourth is the statement of CVC. The final relation is that of the standard model. It states that the weak neutral current is a linear combination of the appropriate isospin partner of the charge-changing currents and of the electromagnetic current itself. See (Wa 75) for details.

One can argue that CVC provided the original insight into the unification of the electroweak interactions, for it states that appropriate parts of the weak and electromagnetic interactions are governed by the same currents. One way to gain some insight into the implications of the standard model, as we shall discuss them, is to think of it as an extension of CVC. Thus the weak neutral current in the last of the above equations is simply related to all of the others. We emphasize that all of these relations are between current densities $J(x)$ and hold locally at each point x in the standard model.

We now discuss some tests of the standard model within this framework. The basic idea is to use the nucleus as a filter to select different pieces of the currents.

Consider first any nuclear transition from an isotopic spin $T = 0$ to a $T = 1$ state. In this case only the isovector parts of the currents contribute and we have the effective relations

$$\tilde{J}_\mu^{(0)} = J_\mu^{V3}(1 - 2\sin^2\theta_W) + J_\mu^{V5}; \quad J_\mu^Y = J_\mu^{V3}. \quad (20)$$

The following relations between electron and neutrino cross sections can then be immediately obtained in the extreme relativistic limit (ERL):

$$\frac{d\sigma_{\nu_\ell \nu_\ell^-} - d\sigma_{\bar{\nu}_\ell \bar{\nu}_\ell^-}}{[d\sigma_{\nu_\ell \ell^-} - d\sigma_{\bar{\nu}_\ell \ell^+}]^{\text{ERL}}} = \frac{1}{2} (1 - 2\sin^2\theta_W) \quad (21)$$

$$\frac{1}{2}[d\sigma_{\nu_\ell \ell^-} + d\sigma_{\bar{\nu}_\ell \ell^+}]^{\text{ERL}} - [d\sigma_{\nu_\ell \nu_\ell^-} + d\sigma_{\bar{\nu}_\ell \bar{\nu}_\ell^-}] = \frac{G^2 q^4}{4\pi^2 \alpha^2} (\sin^2\theta_W)^2 d\sigma_{ee}^{\text{ERL}}. \quad (22)$$

Consider $T = 0 \rightarrow T = 0$ transitions in the nuclear target. In this case, only the isoscalar parts of the currents can contribute and we have the effective relation

$$\mathcal{J}_\mu^{(0)} = -2\sin^2\theta_W J_\mu^\gamma. \quad (23)$$

The relation between cross sections is then

$$\frac{d\sigma_{\nu_\ell \bar{\nu}_\ell}}{d\Omega} = \sin^4\theta_W \frac{G^2 q^4}{2\pi^2 \alpha^2} d\sigma_{ee'}^{\text{ERL}}. \quad (24)$$

Now these are truly remarkable results. They are relations between cross sections that hold at all q^2 . They are independent of nuclear structure. Thus they must hold from wavelengths where only the gross features of the nucleus are observed, to wavelengths where the nuclear structure can be described in terms of point nucleons, down to wavelengths where the baryon and meson structure of the nucleus becomes important, and on to wavelengths where the structure of the system is to be described in terms of quarks and gluons. These relations provide a true test of the unification of the electroweak interaction. For example, Eq. (24) states that no matter how complicated each process is in detail, we should be able to lay the electron-scattering cross section (appropriately scaled) on top of the neutrino cross section and those two curves should be identical.

Some application of these results are given in (Wa 82). We note that the most direct way to test Eq. (24) would appear to be through elastic neutrino scattering from a $T = 0$ nucleus where the recoiling nucleus is observed in the laboratory (difficult, but possible at LAMPF II).

We emphasize that there are only a minimal number of assumptions in this analysis. The essential ingredients are the V-A structure of the currents, Lorentz invariance, the isospin structure of the currents, and isotopic spin invariance in the nuclear system. From the viewpoint of the quark model, another important assumption is that the nucleus is made up of u and d quarks only. At LAMPF II energies this is probably a reasonable assumption, but it is worth noting that if we include c and s quarks, then there is a term of the following form in the weak neutral current in the standard model,

$$\delta \mathcal{J}_\mu^{(0)} \equiv \mathcal{J}_\mu^s = \frac{i}{2} [\bar{c}\gamma_\mu(1 + \gamma_5)c - \bar{s}\gamma_\mu(1 + \gamma_5)s]. \quad (25)$$

This is an isoscalar contribution and, for example, gives rise to an isoscalar axial-vector term. The virtual presence of ($\bar{s}s$) pairs in the nucleus will thus modify our previous expressions to some extent. Estimates of the size of these effects suggest they are likely to be smaller than or comparable to the breakdown of isotopic spin invariance in the nuclear system.

3.3.2.2 Kaon Beta Decay. An independent way to test the CVC hypothesis involves the decays (see Ba 77)

$$(\bar{K})^0 \rightarrow K^\mp e^\pm (\bar{\nu}_e). \quad (26)$$

These are strangeness conserving transitions between states belonging to the same isospin multiplet. Unlike $\pi \rightarrow \pi^0 e \nu$, however, the decay of the ordinary down quark occurs in the presence of a strange (rather than a non-strange) quark.

If the CVC prediction also holds for this case, one expects

$$B(K_L \rightarrow K^+ e^- \bar{\nu}_e) = B(K_L \rightarrow K^- e^+ \nu_e) \approx 3 \times 10^{-9} \quad (27)$$

as the K_L is an almost equal mixture of K^0 and $\bar{K}^0$. Because both decays can be measured in the same beam and apparatus, many systematic uncertainties cancel in a measurement of the difference in the above branching ratios. The difference is proportional to $\text{Re } \epsilon$ (see Sec. 3.4.2). We note that a measurement of the charge asymmetry in $K_L \rightarrow K^\pm e^\mp \nu$ would provide information on $\text{Re } \epsilon$ that is independent of the information obtained from the lepton charge asymmetry in normal $K_{\ell 3}$ decay (K1 76). The latter depends on possible $\Delta S = -\Delta Q$ amplitudes (known from direct measurement (Be 69) to be less than 10% of the $\Delta S = \Delta Q$ amplitudes) whereas the former involves possible $\Delta S = 2\Delta Q$ contributions. Note, however, that according to present theories neither a $\Delta S = -\Delta Q$ nor a $\Delta S = 2\Delta Q$ contribution is expected at an appreciable level.

3.3.2.3 Kinematically Complete ν -Nucleon Scattering. It is an important feature of intermediate-energy neutrino experiments that they yield information that is lost at very high energies. As an example, consider a complete kinematic experiment on the nucleon carried out at all final lepton velocities β and scattering angles θ . From Lorentz invariance and the Dirac equation the single-nucleon matrix element of the weak current must have the following structure:

$$\begin{aligned} \langle p^- | J_\mu^{(\pm)} | p \rangle \propto & \bar{u}(p^-) [F_1 \gamma_\mu + F_2 \sigma_{\mu\nu} q_\nu - i F_S q_\mu \tau_3 \\ & + F_A \gamma_5 \gamma_\mu - i F_P \gamma_5 q_\mu + F_T \gamma_5 \sigma_{\mu\nu} q_\nu \tau_3] \tau_\pm u(p) . \end{aligned} \quad (28)$$

The form factors are all functions of q^2 . The target response is completely characterized by the following five quadratic combinations of form factors (Wa 82) with m (m_ℓ) the nucleon (lepton) mass

$$\begin{aligned} \tilde{F}_1 &= (1 + \frac{q^2}{4m^2}) F_A^2 + \frac{q^2}{4m^2} (F_1 + 2mF_2)^2 , \\ \tilde{F}_2 &= F_1^2 + F_A^2 + \frac{q^2}{4m^2} [(2mF_2)^2 + (2mF_T)^2] , \\ \tilde{F}_3 &= F_A(F_1 + 2mF_2) , \\ \tilde{F}_4 &= (F_A - \frac{q^2}{2m} F_P)^2 + q^2 F_S^2 (1 + \frac{q^2}{4m^2}) , \\ \tilde{F}_5 &= m [F_S(F_1 - \frac{q^2}{2m} F_2) + F_T(F_A - \frac{q^2}{2m} F_P)] . \end{aligned} \quad (29)$$

In the ERL ($m_\ell^2 = 0$), a measurement of the charge-changing neutrino cross sections allows a determination of the following quantities:

$$\left(\frac{d\sigma}{d\Omega} \right)_{\nu_\ell \ell^- \rightarrow \bar{\nu}_\ell \ell^+} \rightarrow \tilde{F}_1, \tilde{F}_2, \tilde{F}_3 . \quad (30)$$

The combinations $\tilde{F}_4$ and $\tilde{F}_5$ do not contribute to the cross section in the ERL.

A complete kinematic experiment at intermediate energy ($m_\ell^2 \neq 0$) as well as a measurement of the polarization, $P(\hat{n})$, of the final charged lepton in the reaction $(\bar{\nu}_\ell, \ell^+)$ (see Fig. 3.-4) allows a separation of all five of the structure functions as follows.

$$\begin{aligned} \left(\frac{d\sigma}{d\Omega}\right) \nu_\ell \ell^- &\rightarrow \left(2 - \frac{m_\ell^2}{q^2}\right) \tilde{F}_1 + \frac{m_\ell^2}{q^2} \tilde{F}_4 \\ \bar{\nu}_\ell \ell^+ &\rightarrow \tilde{F}_2, \tilde{F}_3, \tilde{F}_5 \\ P(\hat{n}) \bar{\nu}_\ell \ell^+ &\rightarrow \tilde{F}_1, \tilde{F}_4. \end{aligned} \quad (31)$$

The quantities F_S and F_T are second-class currents. They are absent in lowest order of the standard model. The quantity $\tilde{F}_5$ is evidently a pure second-class current detector, since it is proportional to F_S and F_T . Let us assume that second-class currents are absent. The term $\tilde{F}_4$ measures the matrix element of the divergence of the axial-vector current, a quantity of prime importance in the theory of strong interactions. The relative contribution of this quantity in the cross sections is characterized by the following parameter.

$$r \equiv \frac{1}{F_A^2} \frac{m_\ell^2}{q^2} \left(F_A - \frac{q^2}{2m} F_P\right)^2 \doteq \frac{m_\ell^2}{q^2} \left(\frac{m_\pi^2}{q^2 + m_\pi^2}\right)^2. \quad (32)$$

Here the Goldberger-Treiman relation and pion-pole dominance of the matrix element of the divergence of the axial-vector current have been invoked to obtain the final form of this expression. This quantity is sketched in Fig. 3.-5 starting at the threshold for μ production.

The assumptions invoked above form the cornerstone of the partially conserved axial-vector-current theory (PCAC). This theory provides one of the major insights into our understanding of the strong interactions. The characteristic q^2 dependence predicted above has never been measured. Note in Fig. 3.-5 how quickly r falls to zero with q^2 ; in fact, PCAC is designed to make this quantity die away with q^2 as quickly as possible. It can only be

Fig. 3.-4. A complete kinematic experiment. $\vec{\beta}$ is the velocity of the final, massive charged lepton and $\hat{n}$ is the direction of polarization.

Fig. 3.-5 The quantity r of Eq. (32).

measured precisely in an intermediate-energy experiment. The quantity $\tilde{F}_4$ can also be studied in the μ -capture reaction.

3.3.4 Tests of Quantum Electrodynamics

The measurement of the anomalous g -value for the muon and electron currently provides one of our most sensitive tests of quantum electrodynamics (QED). The current experimental value for a_μ [$a_\mu = (g_\mu - 2)/2$] is determined (Ba 79) to a precision of 8×10^{-9} , due almost entirely to the errors in counting statistics. The theoretical value a_μ^{th} includes both hadronic and weak contributions, in addition to those from QED; that is, $a_\mu^{\text{th}} = a_\mu^{\text{QED}} + a_\mu^{\text{had}} + a_\mu^{\text{weak}}$. The hadronic contribution is of order 67×10^{-9} and the weak contribution of order 2×10^{-9} . The largest uncertainty in a_μ^{th} comes from the hadronic vacuum polarization terms in a_μ^{had} . Recent high-precision measurements at CERN for the reaction in $e^+e^- \rightarrow \pi^+\pi^- \mu$ indicate that these hadronic terms could be measured with sufficient precision to make a new and better measurement of a_μ very attractive. One would require a repeat of the CERN $e^+e^- \rightarrow \pi^+\pi^-$ measurements at the higher energies available from Fermilab as well as improved calculations of the QED contributions to a_μ . The large proton current available from LAMPF II (about 100 times that used in the CERN storage-ring measurement of a_μ to 7 ppm) coupled with an improved apparatus would allow a measurement of a_μ^{weak} for the first time with a precision of 25%.

The studies of muonium at LAMPF by the Yale group have led to increasingly precise measurements of the muonium hyperfine splitting $\Delta\nu$ and of the ratio of muon to proton magnetic moments. The present uncertainties in these measurements are 0.036 ppm for $\Delta\nu$ and 0.37 ppm for the moment ratio (Ma 82). One may use these measurements, together with other precision measurements of the fine structure constant α , to test theoretical bound-state QED. The agreement is excellent at the level of 0.3 ppm, the primary limitations being uncertainties in some uncalculated QED terms from the theory and uncertainties from the magnetic-moment ratio.

The current technique employed by the Yale group involves the application of a microwave field to drive the transitions between the muonium hyperfine levels. The resonance condition is determined by measuring the decay positron asymmetry, which monitors the substate population. Further improvements in these measurements would require line-narrowing techniques such as could be obtained from the short pulse mode of LAMPF II. For example, one may accept a

1.1- μ s-wide pulse of muons in the target and then wait several (about three) muon lifetimes before counting the decay positrons. Because only long-lived muons are counted, the intrinsic line width is narrower. The details of these measurements are discussed in (Eg 82). He estimates that with intense pulse beams the uncertainties in Δv and the magnetic-moment ratio could be reduced to 5 ppb and 0.06 ppm, respectively.

Laser spectroscopy of muonic or pionic atoms is a field that is really just beginning. The appeal of lasers for studies of energy levels and lifetimes in atomic systems lies in the extraordinary precision achievable. This appeal intensifies as advances in laser technology continue to produce greater varieties of high-powered, tunable lasers. To use these techniques requires pulsed muon beams to match the laser operating conditions.

3.4 WINDOWS FOR NEW PHYSICS

An essential part of the justification for a new facility is its potential for new discoveries. In the area of the electroweak interactions, the potential of LAMPF II is great. In this respect it is complementary to the highest energy accelerators. LAMPF II has the capability to probe for new physics at mass scales even above that of the next generation of high energy accelerators through precision experiments that sense the presence of new mediating particles. Because CP nonconservation is not understood, new experiments hold an excellent potential for new insights. Several reactions and decays have been identified that may lead to surprises. Of course, almost by definition, the real surprises will come in areas not suspected at all, so the discussion here can only be suggestive.

3.4.1 Suppressed Decay Modes

Studies of first forbidden weak processes are motivated by the desire to test higher-order calculations, to constrain the parameters of the standard electroweak model, and to probe for the existence of new interactions. Because of their reduced natural rates these decays provide excellent windows for new physics. Ultimately, a systematic study of many decay modes will be necessary because usually several of these aspects are involved simultaneously in any single process. The examples described here illustrate the importance of LAMPF II for detailed studies of decays with branching ratios of order 10^{-9} or less due to suppressions such as the Glashow-Iliopoulos-Maiani (GIM) mechanism (Gl 70).

3.4.1.1 $K^+ \rightarrow \pi^+ + X$. The decay discussed in this section is one for which the observables measured in the final state do not fully identify the decay mode. These decays are of the form $K \rightarrow \pi + X$ where $K = K^+, K_L$, or K_S , and X consists of unobserved neutrals. Below, we shall phrase our discussion solely in terms of the only case studied experimentally ($K^+ \rightarrow \pi^+ + X$).

In the standard electroweak model the unobserved neutrals are the various $\nu\bar{\nu}$ pairs that can be emitted. More generally, the unobserved neutrals may involve new neutral particles: for example, the axion, the photinos or Goldstinos of supersymmetric models, or the familon. The present experimental limit is (As 81) $B(K^+ \rightarrow \pi^+ + X) < 1.4 \times 10^{-7}$ (90% C.L.).

Fig. 3.-6. Electroweak contributions to $K^+ \rightarrow \pi^+ \Sigma \nu_i \bar{\nu}_i$: (a) induced Z^0 contribution, (b) box diagram.

$K^+ \rightarrow \pi^+ \nu \bar{\nu}$ proceeds via the second-order weak diagrams shown in Fig. 3.-6. The amplitude is approximately proportional to

$$\sum_i U_{id}^* U_{is} m_{q_i}^2$$

(the effect of the GIM mechanism), where the sum is over all charge 2/3 quarks. U is the quark-mixing matrix. Ellis and Hagelin (El 83) derive the following bounds on the branching ratio of $K^+ \rightarrow \pi^+ \sum_{\ell} \nu_{\ell} \bar{\nu}_{\ell}$ in the standard model with three generations

$$4.2 \times 10^{-10} < B(K^+ \rightarrow \pi^+ \sum_{\ell=e,\mu,\tau} \nu_{\ell} \bar{\nu}_{\ell}) < 3.3 \times 10^{-9} \quad (33)$$

The upper bound follows from the constraint on the mixing angles obtained from the second-order weak part of $B(K_L \rightarrow \mu^+ \mu^-)_{\text{disp}}$ described below, and is relatively independent of the effects of heavier quarks. The lower bound in Eq. (33) is obtained from a lower bound on the top-quark couplings $U_{ts}^* U_{td}$, derived from Δm_K , and depends on the amount of the top-quark contribution to Δm_K needed to account for the observed Δm_K . The latter depends in turn on the value of the $K^0 - \bar{K}^0$ mixing matrix element that determines the charmed-quark contribution to Δm_K [the lower bound in Eq. (33) assumes the vacuum-insertion result for this matrix element] (El 83).

It follows from Eq. (33) that a branching ratio significantly above 3.3×10^{-9} would represent evidence for new physics. A branching ratio significantly below 4.2×10^{-10} would be evidence for the existence of additional generations and/or new interactions, provided that the $K^0 - \bar{K}^0$ matrix element is not much larger than the value obtained by vacuum insertion.

A possible source of a large branching ratio would be, for example, the presence of the decay $K^+ \rightarrow \pi^+ \tilde{\gamma} \tilde{\gamma}$ ($\tilde{\gamma}$ = photino), for which Shrock and Suzuki estimate (Su 82, Shr 82) that $B(K^+ \rightarrow \pi^+ \tilde{\gamma} \tilde{\gamma}) < 10^{-7}$. Shrock has emphasized that it may be possible to identify the photino by its decay $\tilde{\gamma} \rightarrow \gamma G$ (G = Goldstino). $\tilde{\gamma} \rightarrow \gamma G$ is the main decay mode of the photino (provided that the mass of the Goldstino is smaller than the mass of the photino), with a lifetime of (Fa 81)

$$\tau_{\tilde{\gamma}} = 1.7 \text{ s} \left(\frac{\Lambda_{ss}}{100 \text{ GeV}} \right)^4 \left(\frac{1 \text{ MeV}}{m_{\tilde{\gamma}}} \right)^5,$$

where Λ_{ss} is the scale of supersymmetry breaking, expected to be greater than or of order 10^2 GeV .

The existence of a light pseudoscalar particle known as the axion (a) follows from the assumption of a Peccei-Quinn (PQ) type global chiral $U(1)$ symmetry in the total Lagrangian. This PQ symmetry is introduced as a possible explanation of the apparent absence (or smallness) of CP nonconservation in the effective QCD Lagrangian.* An anomaly due to instanton effects is absorbed into the PQ-chiral current producing a small mass for the axion.

The present experimental limit (As 81)

*For a review and references see, for example (Pe 81). Supersymmetric theories often naturally include PQ symmetries (Ni 82).

$$B(K^+ \rightarrow \pi^+ a) \equiv \frac{\Gamma(K^+ \rightarrow \pi^+ a)}{\Gamma(K^+ \rightarrow \text{all})} < 3.8 \times 10^{-8} \quad (34)$$

is close to the conservative lower bound of 2×10^{-8} found in (Go 78) and starts to be uncomfortable for the so-called "standard" axion, which is based on a model involving two Higgs doublets. However, values of $B(K^+ \rightarrow \pi^+ a)$ even smaller than 2×10^{-8} are not completely ruled out (Fr 81). Furthermore, the axion may be associated with more complicated Higgs structures, increasing the possible range for mass and couplings of the axion. In special cases, the coupling may be so weak that the axion becomes "invisible" (Pe 81). Under these circumstances, broad searches ruling out ever larger areas of the mass and coupling constant parameter space seem well justified.

A truly massless spinless boson called the familon (f) has also been proposed (Wi 82). These particles accompany the spontaneous breakdown of possible global family symmetries, and their experimental manifestations are in some respects similar to those of axions. The $K^+ \rightarrow \pi^+ f$ branching ratio may be near the experimental limit in Eq. (34).

The tantalizing possibilities described here make it of great interest to improve the limit on the branching ratio for $K \rightarrow \pi + X$.

3.4.1.2 $K_L \rightarrow \ell^+ \ell^-$. The decay $K_L \rightarrow \mu\mu$ occurs with a branching ratio (Par 82)

$$B(K_L \rightarrow \mu^+ \mu^-)_{\text{exp}} = (9.1 \pm 1.9) \times 10^{-9}.$$

The lowest order contributions in the standard model are shown in Fig. 3.-7. The absorptive part of the $K_L \rightarrow \mu\mu$ amplitude, which is due to the 2γ -exchange contribution [Fig. 3.-7(a)] can be calculated reliably, yielding*

$$\Gamma(K_L \rightarrow \mu^+ \mu^-)_{\text{abs}} = (1.2 \times 10^{-5}) \Gamma(K_L \rightarrow \gamma\gamma) = (5.9 \pm 0.6) \times 10^{-9} \Gamma(K_L \rightarrow \text{all}).$$

* (Qu 68, Se 69, Ma 70). See also the review by H. Stern and M. K. Gaillard, Ann. Phys. 76, 580 (1973) and references quoted therein.

Fig. 3.-7. Electroweak contributions to $K_L \rightarrow \mu^+ \mu^-$: (a) 2γ -exchange contribution, (b) induced Z^0 contribution, (c) box diagram.

Subtracting the absorptive part from the experimental value, one finds that the dispersive part contributes to the $K_L \rightarrow \mu^+\mu^-$ rate by the amount

$$\Gamma(K_L \rightarrow \mu^+\mu^-)_{\text{disp}} = (3.2 \pm 2.0) \times 10^{-9} \Gamma(K_L \rightarrow \text{all}) . \quad (35)$$

A careful dynamical analysis of the dispersive matrix element for $K_L \rightarrow \mu\mu$ could provide a useful clue as to the existence of new interactions. One contribution comes from the dispersive part of the 2γ -exchange diagram. Further contributions come from the second-order weak diagrams [Figs. 3.-7(b) and 3.-7(c)]. Only high-precision measurements can place limits on the size of these contributions or the possibility of new interactions.

An interesting observable in $K_L \rightarrow \mu^+\mu^-$ is the longitudinal polarization of the muons. Polarizations of the order of one cannot be excluded; a polarization larger than about 10^{-3} would suggest the presence of a nonelectroweak contribution due to either Higgs mesons or leptoquarks (He 83).

The electroweak contribution to $K_L \rightarrow e^+e^-$ is the same as for $K_L \rightarrow \mu^+\mu^-$, except for $m_\mu \rightarrow m_e$. As a consequence, the electroweak (ew) contribution to $\Gamma(K_L \rightarrow e^+e^-)$ is suppressed by a factor of about 10^3 relative to $\Gamma(K_L \rightarrow \mu^+\mu^-)_{\text{ew}}$ making $K_L \rightarrow e^+e^-$ a sensitive probe of possible nonelectroweak quark-electron couplings. Examples are the couplings arising from neutral flavor-changing Higgs exchange or leptoquark exchange (the contribution from neutral flavor-changing gauge boson exchange is either proportional to the lepton mass or vanishes, depending on the type of coupling). The $K_L \rightarrow e^+e^-$ branching ratio due to either of these mechanisms could be as large (Sha 82) as the present experimental upper limit (Bo 67) $B(K_L \rightarrow e^+e^-) < 2 \times 10^{-9}$. An improvement of this limit therefore would be very valuable. We note that the information one would obtain from such an experiment is different from that which would result from a measurement of the muon polarization in $K_L \rightarrow \mu\mu$, in part because the nonelectroweak quark-electron and quark-muon couplings may be different.

Like kaon decays, hyperon decays that involve such currents should exist in second-order in the electroweak interaction and be calculable in any renormalizable gauge theory. A limit of $[\Gamma(\Sigma^+ \rightarrow p e^+ e^-)]/[\Gamma(\Sigma^+ \rightarrow \text{all})] < 7 \times 10^{-6}$ for such a process is given by Ang using a bubble chamber with a stopped K^+ beam (An 69).

3.4.1.3 $\Delta S = 2$ Hyperon Decays. The increase of cleanly tagged hyperons anticipated at LAMPF II can be used to improve limits on a multitude of suppressed or forbidden decays. No decay that violates the $|\Delta S| \leq 1$ rule in first-order electroweak interaction is known to exist. Searches for such processes consequently are of fundamental importance. For example, the $\Delta S = 2$ transition $\Xi^0 \rightarrow p \pi^-$ was sought by Geweniger et al. (Ge 75) using the CERN neutral-hyperon beam. They set the limit

$$\frac{\Gamma(\Xi^0 \rightarrow p \pi^-)}{\Gamma(\Xi^0 \rightarrow \Lambda \pi^0)} < 3.6 \times 10^{-5} .$$

Similarly, Bourquin et al. in the CERN charged-hyperon beam set the limit (Bou 79)

$$\frac{\Gamma(\Omega^- \rightarrow \Lambda \pi^-)}{\Gamma(\Omega^- \rightarrow \text{all})} < 1.3 \times 10^{-3}$$

for Ω^- decay. These limits and others could be improved at LAMPF II.

3.4.2 CP Nonconservation

3.4.2.1 CP Nonconservation in the Standard Model. The CP-nonconserving phase in the standard model, (Ko 73) leads to a nonvanishing imaginary part of the $K^0 - \bar{K}^0$ amplitude and, in contrast to the superweak model, to a CP-nonconserving component in the $K^0 \rightarrow 2\pi$ amplitude, which is in general not negligible. The suppression of the CP-nonconserving weak interaction relative to the CP-conserving ones is due to a product of quark mixing angles times the CP-nonconserving phase. An understanding of CP-nonconservation experiments requires some familiarity with the neutral kaon system and particular terminology and nomenclature. A brief review is given below. For more details see (K1 76, Wo 69).

The eigenstates of strangeness that are produced in strong interactions are $K^0 (S = +1)$ and $\bar{K}^0 (S = -1)$ where $|K^0\rangle = \text{CP} |\bar{K}^0\rangle$. The weak interactions do not conserve strangeness so the K^0 and $\bar{K}^0$ can mix by second-order weak transitions. The time evolution of the neutral K system is characterized by time-dependent amplitudes α and β for K^0 and $\bar{K}^0$, respectively. These satisfy a Schrödinger equation

$$-\frac{d}{dt} \begin{pmatrix} \alpha \\ \beta \end{pmatrix} = (iM + \frac{1}{2} \Gamma) \begin{pmatrix} \alpha \\ \beta \end{pmatrix}, \quad (36)$$

where M and Γ are each Hermitian matrices called the mass matrix and the decay matrix, respectively, and t is measured in the kaon rest system. The eigenvalues of Eq. (36) are $\gamma_S = iM_S + (\Gamma_S)/2$ and $\gamma_L = iM_L + (\Gamma_L)/2$ where M_i and Γ_i are the mass and the decay width of the K_i .

For CP nonconservation the crucial parameters are ϵ , which provides a measure of the breakdown of CP invariance in the mass matrix, and ϵ' , which is a measure of the CP noninvariance in the $\Delta S = 1$ decay channel. To a good approximation these parameters can be written as

$$\epsilon = \frac{e^{2i(M_L - M_S)/\Gamma_S} \frac{\text{Im} M_{12}}{\text{Re} M_{12}}}{\sqrt{2}}; \quad \epsilon' = \frac{1}{\sqrt{2}} \frac{\text{Im} A_2}{A_0} e^{i(\delta_2 - \delta_0)}$$

where M_{12} is the off-diagonal element of the mass matrix and A_0 and A_2 are the amplitudes for the decay of a K^0 into two pions with isospin 0 and 2, respectively. The quantities δ_0 and δ_2 are the corresponding s-wave pion-pion scattering phase shifts. Studies of these phase shifts constrain the phase of ϵ' to be $45^\circ \pm 10^\circ$ (Cr 81). The phase of ϵ is determined by the measured $K_L - K_S$ mass difference, Δm , and the lifetimes of K_L and K_S to be

$$\phi_n = \arctan(2\Delta m/\Gamma_S) = 43.7 \pm 0.2^\circ. \quad (37)$$

Under the assumption of CPT invariance, the two eigenvectors of Eq. (36) can be written in terms of ϵ alone

$$\begin{aligned} |K_S\rangle &= \frac{1}{\sqrt{2}} (1 + |\epsilon|^2)^{-1/2} [(1 + \epsilon)|K^0\rangle + (1 - \epsilon)|\bar{K}^0\rangle] \\ |K_L\rangle &= \frac{1}{\sqrt{2}} (1 + |\epsilon|^2)^{-1/2} [(1 + \epsilon)|K^0\rangle - (1 - \epsilon)|\bar{K}^0\rangle]. \end{aligned}$$

The three processes in which CP nonconservation has been observed are $K_L^0 \rightarrow \pi^+\pi^-$, $K_L^0 \rightarrow \pi^0\pi^0$, and the difference in the rates for $K_L^0 \rightarrow \pi^+\ell^-\nu$ and $K_L^0 \rightarrow \pi^-\ell^+\nu$, where ℓ is either μ or e . For recent reviews of the theory, see (Wo 80, Wo 81, He 82a). These processes are described by the following parameters

$$|\eta_{+-}|e^{i\phi_{+-}} = \frac{\text{AMP}(K_L \rightarrow \pi^+\pi^-)}{\text{AMP}(K_S \rightarrow \pi^+\pi^-)} = \epsilon + \epsilon'$$

$$|\eta_{00}|e^{i\phi_{00}} = \frac{\text{AMP}(K_L \rightarrow \pi^0\pi^0)}{\text{AMP}(K_S \rightarrow \pi^0\pi^0)} = \epsilon - 2\epsilon',$$

and

$$\delta_\ell = \frac{\Gamma(K_L \rightarrow \pi^-\ell^+\nu_\ell) - \Gamma(K_L \rightarrow \pi^+\ell^-\nu_\ell)}{\Gamma(K_L \rightarrow \pi^-\ell^+\nu_\ell) + \Gamma(K_L \rightarrow \pi^+\ell^-\nu_\ell)} = 2\text{Re}(\epsilon).$$

Experimentally one finds*

$$\eta_{+-} = [(2.27 \pm 0.02) \times 10^{-3}] \exp[i(44.7 \pm 1.2)^\circ],$$

$$|\eta_{00}|/|\eta_{+-}| = 1.02 \pm 0.04, \quad \phi_{00} - \phi_{+-} = (10 \pm 6)^\circ,$$

$$\delta_e = (3.33 \pm 0.14) \times 10^{-3}, \text{ and } \delta_\mu = (3.19 \pm 0.24) \times 10^{-3}.$$

More recent experiments on $K_L \rightarrow \pi^0\pi^0$ yield

$$(\epsilon')/(\epsilon) = (-4.6 \pm 5.3 \pm 2.4) \times 10^{-3} \quad (\text{Wl 84})$$

$$= (1.7 \pm 8.2) \times 10^{-3} \quad (\text{Sc 84})$$

These imply (Kl 76, Wo 69) $|\epsilon| = |2.28 \pm 0.05| \times 10^{-3}$, $\text{Re}\epsilon = (1.64 \pm 0.06) \times 10^{-3}$ (assuming $\Delta S = \Delta Q$), and $-0.0163 \lesssim \epsilon'/\epsilon \lesssim 0.0078$ (90% C.L.).

For future reference, we also quote the present experimental limits on CP nonconservation in $K_S \rightarrow 3\pi$ decay and the limit on T-nonconserving transverse polarization of the muon in $K \rightarrow \pi\mu\nu$, P_n^μ .

$$|\eta_{+-0}| \equiv \frac{\text{AMP}(K_S^0 \rightarrow \pi^+\pi^-\pi^0)_{\text{CP nonconserving}}}{\text{AMP}(K_L^0 \rightarrow \pi^+\pi^-\pi^0)} = 0.21 \pm 0.24 \quad (\text{Me 72}), \quad (38)$$

$$|\eta_{000}|^2 \equiv \left[\frac{\text{AMP}(K_S^0 \rightarrow \pi^0\pi^0\pi^0)_{\text{CP nonconserving}}}{\text{AMP}(K_L^0 \rightarrow \pi^0\pi^0\pi^0)} \right]^2 < 1.2 \quad (\text{Ba 73}),$$

$$P_n^\mu(K_L^0 \rightarrow \pi^-\mu^+\nu_\mu) = 0.0017 \pm 0.0056 \quad (\text{Mo 80}), \text{ and} \quad (39)$$

$$P_n^\mu(K^+ \rightarrow \pi^0\mu^+\nu_\mu) = -0.0030 \pm 0.0047 \quad (\text{Cam 81}). \quad (40)$$

In the standard model, both ϵ and ϵ' can be expressed in terms of $\sin \delta$, where δ is the one CP-nonconserving phase in the Kobayashi-Maskawa mixing matrix, together with absolute values of various K-M matrix elements and dynamical factors. With reasonable assumptions about these factors and using a top quark mass, $M_t \approx 45$ GeV, The model predicts (He 84)

*These numbers come from combining the results of many experiments. The data are most accessible in (Par 82).

$$\frac{\epsilon'}{\epsilon} > 3 \times 10^{-2}.$$

The experimental values given above are barely consistent with this bound raising questions about the validity of the standard model description of CP nonconservation. Even more serious a problem for the standard model is the absolute value of ϵ itself. It can be expressed as a product of K-M matrix elements U_{kl} , the dynamical factor B associated with the "box diagram," and a function of the top quark mass:

$$|\epsilon| = \left(\frac{B}{0.34}\right) |U_{cb} U_{ub}| \sin \phi F(M_t)$$

$$F(M_t) = -0.7 + 0.4 \ln \left(\frac{M_t^2}{M_c^2}\right) + 0.6 \left(\frac{M_t^2}{M_c^2}\right) |U_{cb}|^2 \left(1 - \left|\frac{U_{ub}}{U_{us} U_{cb}}\right| \cos \phi\right)$$

where ϕ is the CP-nonconserving phase introduced by Gronau (Gr 84b). Recent measurements indicate the U_{cb} and U_{ub} are both small: $|U_{cb}| \approx 0.05$; $|U_{ub}|/|U_{cb}| \lesssim 0.14$. Should $|U_{ub}|$ fall much below this limit, then the standard model may not be able to account for the magnitude of ϵ when B is in the generally accepted range of $1/3 \leq B \leq 1$. Obviously much more accurate measurements will be needed to resolve this issue. Further predictions in the standard model include $|\eta_{+-0} - \epsilon| = 2/3 |\eta_{+-} - \eta_{00}|$ and $D_n = (10^{-29} - 10^{-32})$ e-cm for the neutron electric dipole moment (Li 80).

3.4.2.2 Other Models of CP Nonconservation. It is still possible that the phenomenon of CP nonconservation can be completely explained by the standard model. However, the recent experimental results described above have made this possibility much less likely. The superweak model is perhaps the simplest, though not the most economical, explanation of the observed CP nonconservation. In this model, a new interaction is postulated that violates CP invariance and has a $\Delta S = 2$ component; the strength of the interaction is characterized by a coupling constant of the order of $10^{-9} G_F$ (Wo 64). An interaction with such properties can account for the magnitude of the observed CP nonconservation as it leads to

$$|\epsilon| \approx \left| \frac{\text{Im} M_{12}}{\Delta m} \right| \approx 10^{-3}.$$

Note that Δm is of the order of a second-order weak amplitude. The immediate predictions are

$$\epsilon = \frac{-1}{\sqrt{2}} \frac{\text{Im} M_{12}}{\Delta m} e^{i\phi_n},$$

where ϕ_n is given in Eq. (37), and $|\epsilon'/\epsilon| \approx 10^{-6} - 10^{-7}$. Thus, for all practical purposes, $\eta_{00} = \eta_{+-}$. Also, $\eta_{+-0} = \eta_{000} = \epsilon$.

CP-nonconserving effects in $\Delta S = 0$ and $\Delta S = 1$ processes other than those in K^0 decays are negligible, of the order of 10^{-9} in amplitude. In particular, the electric dipole moment of the neutron is estimated to be (Wo 74) $D_n \ll (10^{-29} - 10^{-30})$ e-cm, far below the present experimental limit (Al 81)

$$|D_n|_{\text{exp}} \leq 6 \times 10^{-25} \text{ e-cm}. \quad (41)$$

The predictions of the superweak theory are consistent with all available information on CP nonconservation. The most important task for present and future experiments in the field of CP nonconservation is to search for deviations from the superweak theory. Finding such a deviation would signal the existence of a different source of CP nonconservation. Of course, the superweak interaction could still be present simultaneously.

In the standard model with only one Higgs doublet, the interactions in the Higgs sector conserve CP. However, if the Higgs sector is extended to include more doublets, the couplings will be CP nonconserving, in general (Le 74, Si 76, We 76, Br 80). In the presence of only two doublets the models contain flavor-changing neutral-current interactions. To suppress the latter, the Higgs bosons must acquire very large masses and consequently the resulting CP nonconservation is superweak (Le 74, Si 76). An exception is the case when the Higgs contributions to the imaginary part of the $K^0 \rightarrow \bar{K}^0$ amplitude are suppressed due to some mechanism.* This requires however that some special conditions be met in the model (Sha 82).

With three or more doublets the flavor-changing neutral-current interactions can be eliminated, and the Higgs bosons could therefore be light (of the order of 5 to 10 GeV). The exchange of a Higgs boson generates an effective interaction with strength of order $G(M'M)/m_H^2$, where M and M' are the masses of the fermions to which the Higgs boson, of mass m_H , couples. For appropriate masses, this strength is $\approx 10^{-5}$. CP-nonconserving effects in the light Higgs model have been estimated in (An 78), with the results $|\epsilon'/\epsilon| = 1/50$ and $D_n \approx -2.8 \times 10^{-25}$ ecm. A more recent estimate (Kh 81) gives $D_n \approx 7 \times 10^{-25}$ ecm. In this model, $|\eta_{+-0}| \approx 10^{-3}$ and $|\eta_{+-0} - \eta_{+-}| \approx |\eta_{+-}|$, presumably, although no estimate is available to date. The prediction for $|\epsilon'/\epsilon|$ is clearly in disagreement with the most recent measurements.

CP- (and T-) nonconserving effects are expected also in semileptonic processes. In particular, a muon polarization, P_n^μ , normal to the decay plane in $K_L^0 \rightarrow \pi^+ \mu^+ \nu_\mu$ and $K^+ \rightarrow \pi^0 \mu^+ \nu_\mu$ is expected, with a size of (Zh 80) $P_n^\mu \approx 5 \times 10^{-3}$. This value is near the present experimental limits given in Eqs. (39) and (40). Note that $P_n^\mu = 0$ for an arbitrary combination of vector and axial-vector couplings.

Recent calculations in the model of (Li 80) find ϵ'/ϵ (Sa 81, Des 81, Do 81)** and D_n (Des 82) larger than the experimental upper limits for these quantities [Eqs. (38) and (41)]. This requires an increase of the relevant Higgs-boson masses by roughly a factor of 3 relative to the values needed to account for the experimental value of ϵ . If these estimates persist, the observed CP nonconservation must originate from a different mechanism. Even then, CP nonconservation in the Higgs sector could be present and lead to a value for P_n^μ of $P_n^\mu < 5 \times 10^{-4}$, or, if the $\Delta S = 1$ semileptonic interaction involves a different Higgs boson, to a value as large as the present experimental limit.

* (La 79). In this model $|\epsilon'/\epsilon| \approx 10^{-3}$ to 10^{-4} and $D_n \approx 10^{-26}$ ecm.

** See however (Ch 82a).

A model of CP nonconservation based on $SU(2)_L \times SU(2)_R \times U(1)$ has been proposed by Mohapatra and Pati (Mo 75). This model attributes the observed CP nonconservation to the existence of CP-nonconserving phases between the left-handed and right-handed couplings.

A four-quark version of this model predicts

$$\eta_{+-} = \eta_{00}, \quad \eta_{+-0} = \eta_{000}, \quad (42)$$

as in the superweak model.

With the third-quark family included, (He 82b, Bra 82, Ch 82b) Eqs. (42) hold only under special circumstances, even if the Kobayashi-Maskawa-type CP nonconservations are neglected. In general, $|\varepsilon'/\varepsilon|$ could be near the present experimental limit (He 84). The neutron electric dipole moment is of the order of 10^{-26} e - cm or less if mixing between the left-handed and right-handed bosons is neglected. If this mixing is considered, D_n could be as large as the experimental upper limit (Be 81, He 84).

In summary, three measurable quantities relevant to facilities with high-intensity kaon beams have emerged as particularly important in searching for deviations from the superweak theory.

- (1) The ratio $|\varepsilon'/\varepsilon|$ in $K^0 \rightarrow 2\pi$ decays.
- (2) η_{+-0} , η_{000} and also their deviation from η_{+-} .
- (3) The transverse polarization of the muon in $K_{\mu 3}$ decays.

At LAMPF II, these measurements can be performed in high intensity kaon beams with small phase space and low contamination. This will result in measurements with higher statistical precision, smaller backgrounds, and reduced systematic errors compared with what is now possible.

There are several additional effects yet to be studied. These include $[\Gamma(K^+ \rightarrow 3\pi) - \Gamma(K^- \rightarrow 3\pi)]$ (Sa 81, Des 81, Do 81), an asymmetry in the Dalitz plot for $K^0 \rightarrow \pi^+\pi^-\pi^0$ (Cab 81), $K_L - K_S$ interference effects in rare decay modes, (Wi 81) and CP nonconservation in rare decay modes (such as $K_L^0 \rightarrow \pi^0 e^+ e^-$) (Li 82).

In all nonsuperweak models considered, $|\varepsilon'/\varepsilon|$ could be near the present experimental upper limit. A large value for $|\eta_{+-0} - \eta_{+-}|$ (of the order of η_{+-}) would indicate the presence of a source of CP nonconservation different from the Kobayashi-Maskawa mechanism. $P_n^\mu \neq 0$ would signal the presence of CP nonconservation in the Higgs sector. Therefore, if a particular type of nonsuperweak CP nonconservation is found, the search for other possible types would have to continue, as several possible sources of CP nonconservation could be present simultaneously.

In addition to the mechanisms of CP nonconservation discussed above, there may be further possible sources of CP nonconservation. An example is CP nonconservation mediated by horizontal gauge bosons (introduced as attempts to understand the replication of fermion generations, and the large number of undetermined parameters in the standard model). This CP nonconservation is generally superweak, with $|\varepsilon'/\varepsilon| \approx 10^{-7}$, unless the contribution of the horizontal bosons to the $K^0 - \bar{K}^0$ amplitude is severely suppressed. (Sh 81a)

3.4.3 Right-Handed Currents

The advent of "left-right" symmetric models (Be 77) has renewed the interest in the search for deviations from pure V-A theory (St 80). At present there is a discrepancy between experiment and the predictions of the Cabibbo model in $\Lambda \rightarrow p e \bar{\nu}_e$ decays which cannot be explained by SU(3) symmetry breaking, but can be accounted for if right-handed currents exist (Ok 83). To incorporate the latter, the $SU(2)_L \times U(1)$ electroweak gauge group has to be enlarged to (at least) $SU(2)_L \times SU(2)_R \times U(1)$.

One test of right-handed currents is provided by precision measurements of neutrino helicity. Since the first measurement of the electron-neutrino helicity, (Go 58) it has been a challenge to perform a similar experiment for ν_μ . After other attempts (See Deu 81), a muon capture experiment succeeded in achieving the precision level of $\sim 10\%$.

It was a measurement (Ro 81) of $\langle k_{\text{recoil}} \cdot J_{12B} \rangle$ in the reaction $^{12}\text{C}(J=0) + \mu^- \rightarrow ^{12}\text{B}(J=1) + \nu_\mu$ similar to the approach of (Go 58), using the ingenious implantation technique. The contributions of the induced terms (absent in beta decay) are small and were measured in a separate experiment. They bring the helicity value from -0.96 ± 0.10 to -1.06 ± 0.11 .*

In addition to possible improvements of the experiment quoted above, (Ro 81) one may hope to obtain a precision of $\sim 2\%$ from ongoing experiments that measure the circular polarization of μ -mesic x rays.** Both use, however, other weak-interaction processes (for example, beta decay) with asymmetries measured to only about this precision.

This drawback may be overcome by a measurement of the neutron polarization from the $(\mu^- + p)_{F=0} \rightarrow n + \nu_\mu$ reaction (Deu 81) (but the calibration of the analyzer may be difficult to achieve on the per cent level) or the measurement of the muon polarization from $\pi \rightarrow \mu \nu_\mu$ at rest (surface muons) [P_μ] using an electromagnetic analyzer (Mott scattering) after spin rotation (Deu 81).

Already impressive results have been obtained from a muon decay electron-asymmetry measurement (Ca 84) with a result for the product $P_\mu \xi \delta / \rho > 0.9959$ (90% CL). This implies $M(W_R) > 380$ GeV. Our actual knowledge of muon decay parameters is compatible with $\xi - 1 \approx 10^{-2}$ (assuming $M_e = 0$ and $E_e/E_{e \text{ max}} \equiv X \rightarrow 1$). The necessary positive definite value*** of the decay spectrum, $dN \approx \rho - \xi \delta \cos \theta > 0$, implies

$$\xi < (\rho/\delta) = (0.752 \pm 0.003)/(0.755 \pm 0.009) = 0.996 \pm 0.013 < 1.009 \quad .$$

The situation may change, of course, when new precision measurements of ρ and δ become available.

A further test of the presence of right-handed currents in the $\Delta S = 1$ semileptonic weak interactions could be obtained by a precise measurement of the muon longitudinal polarization, $P_{K\mu}$, in $K^\pm \rightarrow \mu^\pm \nu$ decays, which is predicted to

* L. Ph. Roesch, et al., submitted to Am. J. Phys.

** L. Simmons, private communication.

*** F. Scheck, private communication and (Sc 78).

be $P_{K\mu} = \pm 1$ in the standard model. The present experimental value is (Ha 84) $(P_{K\mu})_{\text{exp}} = -0.97 \pm 0.047$. With the strength of right-handed couplings deduced from Λ -decay data, $P_{K\mu}$ is predicted to be (Ok 83) $P_{K\mu} = -0.85 \pm 0.06$ (including radiative corrections, which are expected to be small, and neglecting the neutrino mass). Whether or not the experimental situation in Λ decays persists, a measurement of $P_{K\mu}$ is a valuable probe of possible deviations from the V-A theory.

3.4.4 Nonleptonic Weak Interactions

An important component to the understanding of weak interactions between quarks will come from the systematic study of hyperon decays. One of the outstanding mysteries has been the $\Delta I = 1/2$ rule in nonleptonic hyperon and kaon decays. With the standard model there is the promise of an understanding of this rule. Specifically, the effective weak-interaction Hamiltonian for quarks is not that given by the GWS model alone. Finite, calculable (in principle) effects are induced by gluon exchanges. These come in two forms: initial- and final-state, etc., interactions and also "penguin" effects, such as shown in Fig. 3.-8. The latter have no physical effects when the strong interactions are absent, as they can be absorbed into a weak interaction wave-function renormalization of the quarks. Because of the vectorial coupling of the gluon, these penguin graphs induce effective nonleptonic right-handed charged-current interactions as well as effective strangeness-changing neutral-current interactions.

These calculations (Fi 78) are promising as they indicate enhancement of $\Delta I = 1/2$ processes. Unfortunately, these "short distance" calculations do not show sufficient relative (to $\Delta I = 3/2$) enhancement to explain the data. This alone may not be too unsatisfactory as there are additional "long distance" (current algebra) pieces that have not been included. However, there are reasons to question even the magnitude of the short distance enhancements (Go 83). Explicit one-loop calculations show that there may be significant omissions in the more popular "leading logarithm approximation" (or Operator Product Expansion) calculations (Ga 74, Al 75).

Nonetheless, this is an active area of theoretical effort (Des 80). If more experimental information on decay channels could be obtained, many of the ambiguities might be resolved phenomenologically (final-state polarizations

Fig. 3.-8. Penguin diagram contribution to strangeness-changing processes.

would be very important). Because they involve the dynamical interplay of the strong and weak interactions beyond the level of tree graphs, studies in this area offer deep tests of the full quantum field theory of the standard model.

Non-leptonic decays such as $\Lambda \rightarrow p\pi^0$, $\Sigma^- \rightarrow n\pi^-$, and $\Xi^- \rightarrow \Lambda\pi^-$ are self-analyzing. The asymmetry parameter α measures the angular distribution of the decay baryon in the hyperon rest frame: $I = 1 + \alpha \bar{P}_y \cdot \hat{q}$ where $\bar{P}_y$ is the hyperon polarization. This asymmetry parameter has been measured for many hadronic decay channels. However, the Ω^- is difficult to produce so that α is poorly known for its various decay modes. For example, Hemingway (He 78) measured $\alpha = -0.2 \pm .4$ for $\Omega^- \rightarrow \Lambda K^-$ decay. The result is in agreement with the calculation of Finjord, (Fi 78) who expects the decay to be nearly parity-conserving. Statistics are weak, however, so much could be learned by increased Ω^- production at LAMPF II.

Bourquin measured (Bou 79) $R = [\Gamma(\Omega^- \rightarrow \Xi^0 \pi^-)] / [\Gamma(\Omega^- \rightarrow \Xi^- \pi^0)] = 2.94 \pm 0.35$ for the decuplet-octet transition. Because the Ω^- contains three strange quarks, the $\Delta I = 3/2$ amplitude is expected to be 20% of the $\Delta I = 1/2$ amplitude, (Ov 76) giving $R = 3 \pm 0.3$ in agreement with experiment, but a more precise test is needed.

The radiative decays, such as $\Sigma^+ \rightarrow p\gamma$, $\Xi^0 \rightarrow \Lambda\gamma$, $\Xi^- \rightarrow \Sigma^-\gamma$, and $\Omega^- \rightarrow \Xi^-\gamma$ are also sensitive to the contribution from the penguin diagram. These decays fall into the general class of decay

$$B_1(P_1) \rightarrow B_2(P_2) + \gamma(k), \quad k = P_1 - P_2,$$

which are described by the transition matrix

$$M = \epsilon_\nu(k) \bar{U}(P_2) [a + b\gamma_5] \sigma_{\mu\nu} k_\mu U(P_1).$$

ϵ_ν is the photon polarization, and a and b are the parity-conserving and parity-nonconserving amplitudes, respectively. The amplitudes can be estimated by QCD calculations. They are related to the decay rate Γ and asymmetry parameter α by the equations

$$\Gamma = \frac{G^2}{\pi} \frac{(M_1^2 - M_2^2)^3}{2M_1} (|a|^2 + |b|^2), \quad \alpha = \frac{2\text{Re}(a^*b)}{|a|^2 + |b|^2}.$$

In the exact SU(3) limit assuming CP invariance, U-spin symmetry, and a current-current form of the weak interaction involving only left-handed currents, α is predicted to be zero (Va 76). Strong-interaction effects can be included via current algebra calculations and yield predictions typically about $\alpha \approx -0.25$ (Ho 71, Ra 71).

Other calculations to determine strong interaction effects employing dispersion relations to evaluate contributions of πN intermediate states have results ranging from -1.0 (Ki 73) to $+0.8$ (Fa 71) depending on input parameters. Theoretical reasons have been presented which rule against such large values of α from strong interactions and most predictions favor smaller values (Fa 71, Re 74, Sc 73).

However, two papers have been independently published that predict a large asymmetry, $\alpha \approx -1$, as a result of a right-handed current between the charmed and strange quarks (Fa 71, Ah 75). Such currents have been suggested to explain the nonleptonic $\Delta I = 1/2$ rule and the large ratio of nonleptonic to semileptonic weak decays (W1 75, DeR 75).

The only measurements of α are two bubble chamber experiments (Ge 69, Ma 80) yielding $\alpha = -0.72 \pm 0.29$ based on 107 events. The value of α has taken on new significance as a result of the above ideas. The statistical quality and systematic uncertainties involved in the present measured value are quite poor. Because α is a fundamental and important parameter in the theory on the radiative decay of the Σ^+ hyperon, a new measurement with good statistics and small systematic errors is needed.

Within the framework of QCD the three types of quark diagrams shown in Fig. 3.-9 are believed to contribute to the direct radiative decay. The four-quark diagram [Fig. 3.-9(b)], not allowed for Ξ^- or Ω^- decays, is allowed for Ξ^0 decay. A comparison of Ξ^0 decay with Ξ^- and Ω^- decay then allows the four-quark diagram to be isolated. This type of analysis has been proposed by a BNL experiment (Po 81). Their data will yield a few-hundred radiative decay events. Even with this addition, it is evident from Table IV that LAMPF II can be used to make a substantial impact on the understanding of these decays.

TABLE IV Radiative hyperon decays.

Reaction	Events	Eff. Denom.	Branching Ratio ^a	α
$\Sigma^+ \rightarrow p\gamma$	107		$(1.20 \pm 0.13) \times 10^{-3}$	-0.72 ± 0.29
$\Xi^- \rightarrow \Sigma^-\gamma$	0	2000	$<1.2 \times 10^{-3}$	
$\Xi^0 \rightarrow \Lambda\gamma$	1	200	$(5 \pm 5) \times 10^{-3}$	
$\Xi^0 \rightarrow \Sigma^0\gamma$	0	60	$<6.5 \times 10^{-2}$	
$\Omega^- \rightarrow \Xi^-\gamma$	0	2400	$<3.1 \times 10^{-3}$	

^aUpper limits are at 90% confidence level as listed in (Par 82).

Fig. 3.-9. Contributions to radiative hyperon decays.

3.5 EXPERIMENTAL CONSIDERATIONS

From the discussion above, we see that LAMPF II permits studies addressing a large number of important physics issues testing the validity of the standard model. The experiments to be done fall into two general classes: precision studies of known processes and sensitive searches for new, rare phenomena. The common denominator of these measurements is the need for beams of unprecedented intensity and purity. Fulfilling this need is the *raison d'être* of LAMPF II.

The raw flux of secondary particles, such as pions, kaons, muons and neutrinos, scales roughly as the product of the flux of protons on target and their energy (i.e., the beam power). The beam power at LAMPF II is more than 50 times that of any machine that has ever operated in this energy region. However, the secondary particle flux delivered to an experiment also depends on the properties of the production target and the secondary beam line. Improvements in existing beam line designs can be expected, as is discussed in Section 5, so that an even larger gain in secondary particle fluxes will be available at LAMPF II.

Experiments using neutrino beams will benefit from the full increase in flux. This will translate into improved counting statistics for cross section measurements, and to improved sensitivity to small neutrino mass differences and mixing angles for neutrino oscillation experiments. However, the full power of these improvements will be realized only if backgrounds and systematic uncertainties (such as beam flux normalization) can also be reduced. The greatly increased neutrino flux at LAMPF II will permit much more detailed, precise measurements of backgrounds and systematic uncertainties thus allowing greatly improved precision in the measurements. This is discussed in more detail in Sec. 5.3.1.

Using present-day technology most experiments studying secondary particle decay will be unable to utilize the full available particle fluxes. This is because of excessive rates in particle detectors and the need to suppress background processes. A more precisely defined incident beam would improve the event definition and so help reject backgrounds. As a result, some of the available flux will be traded for improvements in beam purity and emittance. Which beam parameter needs to be improved depends upon which experiment is to be done.

Some specific examples can be found in the experiments described in Physics with LAMPF II. For the $K^+ \rightarrow \pi \nu \bar{\nu}$ experiment, a stopping K^+ beam is used: beam purity and momentum bite are the crucial beam parameters. (Lo 83) For the measurement of the CP-nonconserving parameter η_{00} , the ability to produce a high flux K_L^0 beam at LAMPF II with a small beam spot and low neutron contamination allows the use of a detector with particularly large acceptance and small systematic uncertainties. (Cr 83) For the measurement of K^0 beta decay (Hi 83), it is most important to have a large kaon flux in a small beam spot.

The best secondary beam energy for these measurements cannot be precisely specified (except for measurements of decays at rest). Extremely high energies are not appropriate as particle identification becomes poor. At very low energies, the detection apparatus needs to be quite large and problems induced by multiple Coulomb scattering become severe. The optimum energy appears to be

from several GeV to about 10 GeV. LAMPF II is ideally suited to producing kaon beams in this region.

The high intensity of LAMPF II will provide the flexibility to tailor the secondary beam parameters to the needs of an individual experiment while maintaining particle fluxes that are unattainable today. This will result in significantly improved measurements that will test the standard model with unparalleled sensitivity. Improvements in detector technology will permit even more sensitive experiments capable of running with beams of higher intensities.

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CHAPTER 4.

DESIGN OF THE ACCELERATORS

CHAPTER 4. DESIGN OF THE ACCELERATORS

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4.1 INTRODUCTION AND SUMMARY

4.1.1 Summary of Performance

LAMPF II consists of a 6-GeV, 60-Hz, 170- μ A rapid-cycling booster and a 45-GeV, 3-Hz, 34- μ A main ring. The main-ring magnets have a flattop to provide 50% duty factor slow-extracted beam. The booster provides only fast-extracted beam. Four booster pulses of twenty are used to fill the main ring. The booster will be 100% utilized, because the remaining 16 pulses are delivered to the neutrino and pulsed-muon areas. This plan minimizes costs by eliminating the need for stretcher or accumulator stages. Future addition of a 45-GeV stretcher would allow an increase of a factor of 2 in the main-ring current and a corresponding improvement in the duty factor to 100%.

The figure of merit for any kaon factory proposal is events per year. In Tables 4.1.1-I and 4.1.1-II we have compared several existing and proposed facilities capable of producing kaons, antiprotons, and neutrinos in the energy range of LAMPF II. The operating schedule is an essential ingredient of this comparison, since only LAMPF II produces simultaneous slow- and fast-extracted beam. The other facilities will typically operate for half of the year with fast extraction and for the remainder of the year with slow extraction, or will divide the available current between slow- and fast-extracted users.

To construct Table 4.1.1-I we have assumed that the secondary beam lines for each accelerator are identical. Production cross sections were obtained from radial scaling fits to the existing data as discussed in Chap. 5. For neutrinos, we have assumed ideal horns for each energy and identical unbiased detectors. The neutrino flux calculations are also discussed in Chap. 5.

An examination of Table 4.1.1-I shows that LAMPF II is a superb source of high-energy kaons and antiprotons, while satisfactorily meeting the needs of the neutrino and low-energy kaon users. Most important, the high-energy main ring provides more opportunities for high-energy polarized proton, kaon, or antiproton experiments than any of the other facilities. These experiments will be particularly important for elucidating the role of quarks and gluons in nuclei as is discussed in Chap. 2.

TABLE 4.1.1-I

NORMALIZED FLUX FROM EXISTING AND PROPOSED "KAON FACTORIES"

Accelerator	K^-/Yr (1 GeV/c)	K^-/Yr (7 GeV/c)	$\bar{p}/\text{Yr}$ (7 GeV/c)	$\nu + e$ Scatters/Yr
AGS (1984) ^a	1	1	1	1
SIN II ^a	96	46	35	37
TRIUMF II ^a	139	161	176	68
LAMPF II ^b	104	193	302	37

^a Assumes 20 weeks/yr - fast-extracted protons
20 weeks/yr - slow-extracted protons

^b Assumes 40 weeks/yr - fast extraction and slow extraction simultaneously

TABLE 4.1.1-II

PARAMETERS FOR EXISTING AND PROPOSED "KAON FACTORIES"

Accelerator	Proton Energy (GeV)	Protons/ Pulse ($\times 10^{13}$)	Beam Rate (Hz)	Beam Current (μ A)	Beam Power (kW)	Duty Factor
CERN PS						
Fast Extraction ^a	24	2.0	0.67	2.2	53	4×10^{-6}
Slow Extraction ^a	24	2.0	0.38	1.2	29	50%
Brookhaven AGS (1984)						
Fast Extraction ^b	27.5	1.5	0.67	1.6	44	4×10^{-6}
Slow Extraction ^b	27.5	1.2	0.38	0.74	20	50%
SIN II						
Fast Extraction ^b	20	1.0	50	80	1600	3×10^{-5}
Slow Extraction ^b	20	1.0	50	80	1600	100%
TRIUMF II						
Fast Extraction ^b	30	2.0	30	100	3000	5×10^{-5}
Slow Extraction ^b	30	2.0	30	100	3000	100%
LAMPF II						
Fast Extraction ^c	6	1.75	48 ^d	136 ^d	816 ^d	5×10^{-5}
Slow Extraction ^c	45	7.0	3	34	1530	50%

^aNo longer available for experiments.

^bFast and slow extraction not available simultaneously.

^c6 GeV and 45 GeV available simultaneously.

^dAvailable for use at 6 GeV during simultaneous operation.

The accelerator designs are simple and conventional. The booster is injected by the 800-MeV H^- beam of LAMPF. Because of its high energy, high peak current, and high beam availability, the LAMPF linac is the best injector for a high-intensity proton synchrotron presently available anywhere in the world. The large apertures of the booster and main ring, which are much larger than the nominal beam size presented in this report, give us the flexibility to solve space charge and transverse stability problems without large beam losses. Several innovations are included in the accelerator design. We are developing ferrite-tuned cavities utilizing perpendicular bias to eliminate the high power losses observed in the tuners of all recent proton synchrotrons. Perpendicular bias is common in the microwave industry, but seems to have been overlooked by the accelerator community. A unique dual-frequency resonant magnet power supply substantially reduces the peak rf power requirement of the booster, thus minimizing cost. The booster is designed with a transition energy above its extraction energy, and the main ring has a transition energy below its injection energy. This design reduces the beam instabilities associated with operation near transition, and eliminates entirely the need for complicated transition-crossing hardware. The long straight sections of the main ring facilitate efficient slow

extraction without expensive large-aperture requirements in the dipoles. Both machines are protected from activation by collimator systems. Full-aperture, fast-abort systems are provided in both rings to dump the beam in the case of faults in the rf or magnet systems.

Los Alamos has the necessary infrastructure to design, build, and successfully operate a high-intensity accelerator facility. The LAMPF staff has unmatched experience in dealing with high-intensity beams in accelerators and experimental areas. The existence of the infrastructure, the staff, and the LAMPF linac argues that Los Alamos is the best place to build a facility of the class of LAMPF II.

4.1.2 Techniques Used to Reduce Beam Losses at High Intensity

High initial intensity is of little use if beam losses cannot be made small. It is our intention to keep the beam losses to less than 0.1% of the beam intensity to avoid activation of the accelerators. All losses larger than 0.1% will be localized in collimators, beam dumps, or septa specially designed to handle the power levels and activation. At LAMPF reducing activation decreased the maintenance time and minimized the need for expensive remote-handling procedures, resulting in a higher overall availability for the experimenters. In principle, beam losses at injection are almost zero because we are injecting H^- instead of protons. There will be some scattered particles produced as the beam passes through the stripping foil, but these particles will be removed by the halo stripper before the acceleration process begins. Both the booster and the main ring have an aperture four times larger than the expected maximum beam size, which occurs early in the acceleration cycle. This leaves ample space for closed-orbit distortion and any emittance growth that might occur because of nonlinear resonances or transverse space-charge effects.

Careful attention has been given to potential instabilities of the beam. The tunes have been chosen to avoid low-order, single-particle resonances. Feedforward and feedback systems will be necessary to damp longitudinal oscillations associated with the accelerating cavities. The microwave instabilities, which have wavelengths shorter than the bunch length, cannot be controlled with dampers and feedback. They can be eliminated, however, by Landau damping, which makes use of the natural tune spread and energy spread of the beam if the transverse and longitudinal impedance of the beam environment is small. Calculations show that if we make no effort to increase the tune and energy spreads, we must keep the longitudinal impedance less than 11.5 ohms for the booster and 14.5 ohms for the main ring. The transverse impedance must be less than 7 M Ω /m for the booster and 10 M Ω /m for the main ring. Through the use of modern construction techniques, the impedances of the beam chamber and beam monitors, kicker magnets and rf cavities can be reduced to very low levels. The so-called space charge impedance appears to require the use of an active transverse damper, but this requires further analysis. Modern machines such as DORIS at DESY have achieved very low impedances, and proposed accelerators like SSC must find ways to reduce the impedance even further. It is difficult to be sure in advance that all sources of instabilities have been examined and accounted for. We have examined many of the known instabilities, but analysis must continue up to the time of final design. In any case, we have planned space in the lattices for transverse superdampers.

Should the beam become unstable for any reason, such as power surges, magnet failure, etc., the one-turn, full-aperture beam-abort system will protect the accelerators from damage and activation of machine components.

4.1.3 Short Description of the LAMPF II Complex

Figure 4.1.4-1 is a plan view of the accelerators and experimental areas. LAMPF produces 120 macropulses per second of either H^+ or H^- particles. LAMPF II will take every second macropulse (60 Hz), leaving plenty of intensity for the PSR-WNR (Proton Storage Ring/Weapon Neutron Research) Facility. The 60-Hz injection pulses for the booster will be 1 ms long, chopped to 50-MHz microstructure, with 3×10^8 protons per micropulse. During 762 turns, 60 out of 66 rf buckets in the booster will be filled, leaving a 6-bucket gap for the extraction kicker ramp. After the 1-ms injection the dipole magnetic field will rise sinusoidally at a frequency of 42.55 Hz for 11.75 ms. Extraction occurs in one turn, and the magnetic field in the booster falls to its initial value in 3.92 ms. The extracted pulse train from the booster can be injected into the main ring or sent to the neutrino production target in Area N. It takes four pulse trains from the booster to fill the main ring (50 ms). The next 16 pulse trains are sent to the neutrino target area while the main ring raises the proton energy from 6 GeV to 45 GeV (87 ms), slow extracts for 168 ms (50% duty factor), and decreases its magnetic field to the initial value (29 ms). The slow-extracted beam from the main ring can be split in the switchyard and sent to Area A and Area H simultaneously. The 3-Hz cycle then starts over again. One possible variation on this cycle might be to accelerate polarized protons in the booster and use the main ring as a storage ring. When the main ring is full, 6-GeV polarized protons could be sent to Area A and Area H.

Fig. 4.1.4-1 Site plan for LAMPF II.

Tables 4.1.4-III through 4.1.4-V contain numbers describing the rings and the acceleration cycle.

4.1.4 Supplementary Documentation and Evolving Design

The design of LAMPF II has been evolving for over two years and is still under revision. Supporting detailed calculations are printed in supplementary documents, LAMPF II Technical Notes, which are kept up to date as necessary and are available upon request.

4.1.5 Need for Accelerator Research and Development

Although substantial progress has been made, much work remains to be done. Beam-dynamics issues requiring more study include space charge, collective effects, and active damping of instabilities. Our prototype effort should be expanded to include a rapid-cycling magnet, vacuum chamber, and power supply, and also kickers and pulsed power supplies. We must prepare for placing contracts for hardware in the commercial sector. Cost optimization must be attempted on all components and conventional facilities. We expect a significant R&D program in advance of authorization for construction.

TABLE 4.1.4-III

GENERAL CHARACTERISTICS OF LAMPF II ACCELERATORS

	<u>LAMPF I</u>	<u>Booster</u>	<u>Main Ring</u>
Accelerated particles	H ⁻	protons	protons
Injected energy	0.75 MeV	0.797 GeV	6.0 GeV
Final energy	800 MeV	6.0 GeV	45.0 GeV
Ring circumference	---	330.88 m	1323.53 m
Harmonic number	---	66	264
Repetition rate	120 Hz	60 Hz	3 Hz
RMS bunch width at maximum energy	0.1 ns	1.5 ns	1.7 ns
Pulse separation at injection	---	19.88 ns	16.88 ns
Emittance--horizontal at inj.	20 π mm-mrad	7.5 π mm-mrad	1.6 π mm-mrad
Emittance--vertical at injection	20 π mm-mrad	7.5 π mm-mrad	1.6 π mm-mrad
Particles/pulse	6 $\times$ 10 ¹³	1.8 $\times$ 10 ¹³	7.1 $\times$ 10 ¹³
Particles/s	6 $\times$ 10 ¹⁵	1.1 $\times$ 10 ¹⁵	2.1 $\times$ 10 ¹⁴
Average current	1 mA	170 μ A	34 μ A
Injection time	---	1 ms	50 ms
B-field rise frequency	---	42.55 Hz	5.71 Hz
Extraction time	---	1.1 μ s	166.67 ms
B-field reset frequency	---	127.5 Hz	17.14 Hz
Circ. freq. at injection	---	0.762 MHz	0.224 MHz
Circ. freq. at extraction	---	0.898 MHz	0.224 MHz
rf at injection	200 MHz	50.3125 MHz	59.25 MHz
rf at extraction	800 MHz	59.250 MHz	59.79 MHz
Average beam power	0.8 MW	1.02 MW	1.53 MW

TABLE 4.1.4-IV

MAGNET LATTICE

	<u>Booster</u>	<u>Main Ring (Racetrack)</u>
Dipole type	combined function	combined function
Focusing order	FODO	ODOF O FODO
Number of cells	32	24 (bending cells)
Superperiodicity	8	2
Horizontal/vertical betatron tune	7.17/6.30	7.45/6.45
Horizontal/vertical phase advance/cell	80.7°/71.0°	90°/79°
Number of dipoles	64	96
Radius of curvature	20.372 m	117.84 m
Length of dipoles	2.0 m	9.25 m
Dipole field at injection	0.239 T	0.280 T
Dipole field at extraction	1.123 T	1.30 T
Focus/defocus, field index in dipoles	-59.0/54.9	-85.5/80.5
Number of quadrupoles	0	24
Gradients at injection	---	<2.87 T/m
Maximum gradients ^a	---	<13.33 T/m
Typical lengths	---	1.0 m
Number of sextupoles	32	62
Maximum gradients	15 T/m ²	-11.1 T/m ²
Chromaticity horizontal/vertical	Corrected to zero for both	
Typical lengths	0.3 m	0.25 m
Number of long straights	32	2
Length of long straights	~4.0 m	158.3 m

^aThis is the gradient in the separate quads; the gradient in the dipoles is <1.35 T/m.

TABLE 4.1.4-V
ACCELERATION SYSTEM

	<u>Booster</u>	<u>Main Ring</u>
Harmonic number	66 ^a	264 ^a
rf range (inject/extract)	50.3/59.3 MHz	59.3/59.8 MHz
Avg. acceleration rate	442.8 GeV/s	445.7 GeV/s
Max. energy gain/turn	0.80 MeV	3.10 MeV
Maximum rf power to beam	2.00 MW	7.95 MW
Average beam current in bunch	2.5 A	2.5 A
Maximum voltage	1.31 MV/turn	5.09 MV/turn
Voltage per cavity	81.8 kV	125.0 kV
Maximum synchrotron phase	37.8°	38°
Maximum synchrotron tune	0.048	0.020
Number of cavities	16	40
Length of cavities	1.17 m	1.17 m

^aSix rf buckets in the booster and 24 in the main ring will be left empty to accommodate fast extraction from the booster and injection in the main ring.

4.2 BOOSTER

4.2.1 Lattice

4.2.1.1 Specifications. The booster circumference is constrained to be a submultiple of the main-ring circumference. A cost optimum was found at quarter size (circumference = 330.8 m).

We required the booster design to have a transition energy either above the final energy or imaginary, to avoid crossing transition. The major existing machines all cross transition, and there have been numerous studies of this problem. At intensities of 10^{13} /pulse or less, it has been shown that it is possible to cross transition with losses that are too small to measure. However, as the intensity is raised, the CERN machines have losses that show up when attempting to make the phase space larger in preparation for transition crossing. The Brookhaven AGS has difficulty making a good slow spill under the same conditions that minimize transition losses. Although it appears possible to create a transition-jump system to avoid losses, we found it possible to avoid transition altogether. We propose a booster lattice constructed from combined-function dipoles with the drift between magnets varied in a way that raises the transition energy (Gu 83). The machine is made up of eight superperiods. A layout of one superperiod of the proposed lattice is presented in Fig. 4.2.1-1.

The working point was chosen to give a large Laslett tune shift and is sufficiently close to a third-integer tune that slow extraction will be possible in the future. The design has the horizontal and vertical tunes differing by one, which eliminates coupling resonances below those driven by skew sextupoles. The final tune design is 7.17 horizontal and 6.30 vertical. The tune diagram is shown in Fig. 4.2.1-2. The working point is at the upper right-hand part of the shaded area. This necktie-shaped region represents the range of single-particle tunes under the influence of space-charge forces. A parabolic charge distribution in transverse xy space is assumed. This diagram shows all sum and difference resonances up to and including order 5. The nearest sum resonance is $\nu_x + 4\nu_y = 32$, which is a fifth-order resonance with sufficiently narrow width that it will not result in beam blowup (see Sec. 4.2.5.1).

Fig. 4.2.1-1. Layout of 1/8 of LAMPF II booster. F1 through F4 and FS are drifts, HFI and HDI are horizontally focusing and defocusing combined-function dipoles, respectively, and SF and SD are sextupoles; rf stations are located in the drifts labeled F2.

Fig. 4.2.1-2. LAMPF II booster tune diagram showing all resonances through fifth order. Difference resonances are shown as dotted lines and sum resonances as solid lines. The shaded region is occupied by the beam at injection time.

To achieve rapid cycling, our booster design has 16 rf cavities. The cavity length is 1.17 m and its design is based on the Fermilab booster system. The ring has 16 accelerating straight sections of 2.7 m each and 32 straight sections of 4.1 or 4.6 m.

Rapid cycling also dictates low dipole fields, chosen to be 0.239 to 1.125 T. A curved, laminated magnet design with parallel ends keeps costs down. A combined-function magnet also reduces total magnet system costs and reduces ring size while permitting a high superperiodicity lattice.

The magnet aperture is set by tune requirements and safety factors as follows. First, we determine the phase space that yields a Laslett tune shift of -0.20 in the vertical plane at the worst time in the cycle, approximately 0.5 ms after the start of acceleration. We then assume the same phase space in the horizontal plane. To determine the minimum clear aperture, we require an admittance of four times this space plus an allowance for horizontal dispersion and an additional allowance of ± 1 cm for closed-orbit distortion in the horizontal plane and ± 0.5 cm in the vertical plane. Finally, we add an allowance of ± 0.65 cm in the vertical plane for vacuum-chamber wall thickness and clearance to the magnet poles. This large aperture will reduce spill, facilitate tuning, and permit future operation at higher currents. Another option is to operate the machine with a space-charge tune shift of 0.12 and still have a safety factor of two in acceptance.

4.2.1.2 Details of lattice. The lattice is arranged in eight superperiods. Each superperiod consists of four FODO cells. The drift lengths have been adjusted to raise the transition kinetic energy to 8.44 GeV ($\gamma_t = 10.0$).

The matched amplitude and dispersion functions for the booster are shown in Fig. 4.2.1-3. The dispersion crosses zero near the end of two moderate drifts. A detailed layout of one superperiod is shown in Fig. 4.2.1-1. This figure shows the location of the rf cavities and sextupoles. The fast-extraction system and fast-abort system will make use of the two drifts labeled F3; the injection system will be built into two adjacent drifts labeled F1. The chromaticities are given by

$$\xi_x = (\Delta v_x / v_x) / (\Delta p / p) = -1.26 \quad (4.2.1-1)$$

$$\xi_y = (\Delta v_y / v_y) / (\Delta p / p) = -1.37 \quad (4.2.1-2)$$

4.2.1.3 Tracking studies. The booster lattice was developed with the program DINGBAT, written by R. Servranckx of Saskatchewan University (Se 84). This program treats each element in second order identically to the program TRANSPORT. In the tracking mode all higher order terms, due to cross-coupling of elements, are kept to the accuracy of the computer. Although the second-order treatment does not give a symplectic transformation for thick-element

Fig. 4.2.1-3. Matched amplitude functions $\beta_x(s)$, $\beta_y(s)$, and dispersion function $\eta_x(s)$ for one superperiod of LAMPF II booster. The layout above the plots shows lengths of magnetic elements. The plot on the right shows the beam amplitudes assuming an emittance of 7.5π mm-mrad and momentum spread $dp/p = \pm 0.3\%$.

combined-function dipoles, the computational error was made negligible by subdividing the magnets. The computational error thus introduced was not important for up to several hundred turns around the lattice.

Field-error effects may be classified as systematic (such as arise from saturation, remnant fields, eddy currents, and design errors) and random (from construction variation and alignment errors). Because the problem is substantially nonlinear, the tracking study was done with both errors introduced simultaneously in DINGBAT as thin multipoles. The systematic field errors added were sextupole and decapole components with midplane symmetry (upright orientation). The random field errors were thin quadrupole, sextupole, octupole, and decapole components with random orientation at the middle and ends of each dipole. A systematic sextupole component can make a large change in the natural chromaticity as was observed in the Fermilab main ring (St 75). Indeed, a sextupole error field dB/B_0 of 0.2% at 5 cm from the center of the magnet can change the chromaticity by a factor of three.

In order to have a well-defined problem for tracking purposes, we have arbitrarily adjusted these correction sextupoles to give a chromaticity of zero. Realistic operating conditions would lead to a slightly negative chromaticity in both planes. The systematic decapole error was chosen to have the same sign as the sextupole. This choice of sign leads to the maximum effect on the beam.

Using the 192 multipoles with systematic and random components as defined above, we did a search to determine the maximum component of each that can be tolerated without unstable behavior of the beam. Such a search cannot give precise limits, but generally we found that increasing an error by a factor of 2 from the tolerance given in Table 4.2.4-I will lead to instability. Note that non-linearity is important in that the maximum tolerable field errors for systematic or random errors individually would have been five times larger. We compared the results obtained in tracking with the more traditional approach involving individual resonances. It is important to keep in mind that each field error affects not only the lowest order resonance allowed by symmetry, but also all higher order resonances of similar symmetry. The results of tracking 499 turns in the booster are presented in Fig. 4.2.1-4. For this figure, error

TABLE 4.2.1-I

MAXIMUM TOLERABLE FIELD ERROR, LAMPF II BOOSTER DIPOLES

<u>Component</u>	<u>Field Error dB/B_0 @5 cm (%)</u>	<u>Type</u>	<u>Orientation (degrees)</u>
Sextupole	0.12	systematic	vertical
Decapole	0.06	systematic	vertical
Quadrupole	0.03	random	random
Sextupole	0.03	random	random
Octupole	0.03	random	random
Decapole	0.03	random	random

Fig. 4.2.1-4. Tracking 499 turns through LAMPF II booster with maximum tolerable field errors from Table 4.2.1-I. Three particles at the maximum acceptance (40π mm-mrad) are tracked. Particle 1 has $dp/p = 0$, particle 2, -0.5% , particle 3, $+0.5\%$.

components as shown in Table 4.2.1-I were included. The three particles tracked are at the surface of a 40π mm-mrad phase space. Particle 1 is at the central momentum, particles 2 and 3 are at -0.5 and $+0.5\%$, respectively.

In Fig. 4.2.1-5 we show tracking results for one particle whose momentum varies $\pm 0.5\%$ while undergoing synchrotron oscillations with a synchrotron tune of 0.04. The same set of errors was used for both Figs. 4.2.1-4 and 4.2.1-5. The tune and matched beta functions are functions of momentum. In extreme cases

Fig. 4.2.1-5. Tracking 499 turns through LAMPF II booster with maximum tolerable field errors from Table 4.2.1-I. Synchrotron oscillations from 16 uniformly spaced cavities per turn force the particle momentum to oscillate from -0.5 to $+0.5\%$ with a synchrotron frequency of 0.04. The particle is at the maximum acceptance (40π mm-mrad) in both planes.

these can vary so rapidly with momentum that important synchrotron-betatron effects can result from this dependence. The tune and matched beta functions as a function of momentum for the LAMPF II booster are presented in Fig. 4.2.1-6 using error fields as defined in Table 4.2.1-I. Note that over the central 1% dp/p , these quantities vary slowly with momentum.

4.2.1.4 Trim dipoles and quads. Flexibility of tune adjustment is necessary to correct errors and to provide for unforeseen operating modes in the future. Trim quadrupoles to adjust the tune will be a part of the LAMPF II booster, but the design requirements for trim adjustments other than sextupoles are not yet specified in detail.

4.2.1.5 Utilization of the long straight sections. There are 16 drifts of 4.6 m and 16 drifts of 4.1 m in the booster lattice. In addition, there are 16 drifts of 2.7 m for the rf cavities. The 4.6-m drifts will be used for the sextupoles for chromaticity correction, for the collimation system, and for the injection system. The 4.1-m drifts have low dispersion and are useful for the fast-extraction and slow-abort systems. It is clear that most of the 4.1-m drifts will be available for other uses in the future. Such uses include possible slow extraction and internal targets.

4.2.1.6 Closed-orbit distortions and magnet tolerances. Lattice imperfections arising from misalignments and field errors lead to orbit distortions and apparent increase in beam size. Limits on practical apertures thereby imply design tolerances. The following two subsections summarize the theory and the results of our evaluation.

Fig. 4.2.1-6. Tune and matched beta functions vs momentum for LAMPF II booster with maximum tolerable field errors given by Table 4.2.1-I. Horizontal functions are shown with solid lines, vertical functions are dashed.

4.2.1.7 Summary of theory. Magnet errors and misalignments give rise to closed-orbit distortions (COD). The standard treatment (Co 58) uses the equation of motion for a forced harmonic oscillator

$$\frac{d^2 \eta_z}{d\phi^2} + v_z^2 \eta_z = g(\phi) \quad , \quad (4.2.1-3)$$

where the subscript z can mean a horizontal or vertical coordinate. The quantity η_z is related to the physical displacement by $z = \beta_z^{1/2} \eta_z$ and $\phi = \int ds / v_z \beta_z$. The quantity g is the azimuthal pattern of the perturbation and can be expressed by

$$g(\phi) = v_z^2 f(\phi) = v_z^2 \beta_z^{3/2} \frac{\Delta B_{\bar{z}}(\phi)}{B\rho} \quad , \quad (4.2.1-4)$$

where $\Delta B_{\bar{z}}$ is an error field. If z means x , then $\bar{z}$ means y and vice versa. The solution to Eq. (4.2.1-3) is written as

$$\eta_z(\phi) = \frac{v_z}{2|\sin\pi v_z|} \int_{\phi}^{\phi+2\pi} f(\psi) \cos v_z(\pi + \phi - \psi) d\psi \quad . \quad (4.2.1-5)$$

We assume randomly distributed errors and calculate the rms displacements $\Delta z(s)$ using the following expression,

$$\frac{\Delta z(s)}{\beta_z^{1/2}(s)} = \frac{1}{2\sqrt{2} |\sin\pi v_z|} \left[\sum \beta_{zi} \delta z_i'^2 \right]^{1/2} \quad . \quad (4.2.1-6)$$

The summation is over all magnets and the error term $\delta z_i' = \Delta(B\ell)/(B\rho)$, where ℓ is the length of the magnet i . To 98% confidence the maximum COD would be

$$z_M(s) = 3.4 \times \left(1 + \frac{|\sin\pi v_z|}{3} \right) \left[\frac{\beta(s)_z}{\bar{\beta}_z} \right]^{1/2} \Delta z(s) \quad (4.2.1-7)$$

for an elliptic vacuum chamber, where $\bar{\beta}_z$ is the average value of the amplitude function.

The main contributions to the COD are from (1) horizontal or vertical displacement of quadrupoles, (2) longitudinal displacements of dipoles, (3) magnitude errors in the dipole field or length, and (4) rotations around the longitudinal axis of the dipole (roll). In addition, stray fields in the straight sections can affect both the x and y COD.

4.2.1.8 Alignment and calibration tolerances. We simulated the effects due to misalignments and field and length errors in the 64 combined-function booster dipoles. The dipole field at injection is 0.239 T, the bend angle $\theta = 5.625^\circ$, the length $\ell = 2.0$ m, the radius of curvature $\rho = 20.3718$ m, and the field indices $n_F = -59.0$ and $n_D = 54.0$. Equations (4.2.1-6) and (4.2.1-7) were evaluated for the booster. The tunes used were $\nu_x = 7.17$ and $\nu_y = 6.30$; the average beta is R/ν , where R is the mean radius of the machine (52.65 m). Injection beam energy is 797 MeV ($B\rho = 4.869$ T-m). The dipoles were assumed to have field errors $\Delta B/B = \pm 10^{-4}$ rms. Alignment errors were assumed to be ± 0.1 mm transverse and ± 0.2 mm longitudinal. The maximum COD values (z_M) were evaluated using Eq. (4.2.1-7). In the HF dipoles we obtain the maximum values $x_M = 19.1$ mm and $y_M = 3.2$ mm; similarly in the HD dipoles we obtain $x_M = 5.6$ mm and $y_M = 9.3$ mm.

Our conclusion from these studies is that it will be straightforward to get a coasting beam inside the aperture with routinely available alignment tolerances. More acceptance should be achieved with an active closed-orbit correction system or with more precise alignment.

4.2.2 Injection System

4.2.2.1 Properties of the LAMPF beam. The LAMPF high-intensity H^- beam emittance is assumed to be 0.5π mm-mrad (it will be known definitely in mid-'85 with PSR operation). The beam energy spread is typically 0.5×10^{-3} rms $\Delta p/p$ and the microbunch width is about 0.1 ns (30° of 805-MHz phase). If necessary, it is possible to collimate the beam energy and transverse profiles in the booster injection beam line by stripping. The beam will be chopped at 50 MHz to provide clean booster injection with 100% capture efficiency (traditional debunching and adiabatic capture leads to injection efficiency of only 75%). The injection rf voltage of the LAMPF II booster is 0.6 MV/turn. The booster momentum acceptance (0.3%) and the phase acceptance at 50 MHz, are sufficiently large compared to the LAMPF beam parameters that a bucket rotator appears unnecessary, although one has been left in the design provisionally; experience with PSR will be used to make a decision on its necessity.

4.2.2.2 Injection scheme. H^- will be injected through a pulsed-magnet achromatic bump and stripper in a straight section as shown in Fig. 4.2.2-1. The central magnet is 1 m long with 5 kG field. The fraction of H^- that will be stripped by this magnet is less than 10^{-3} . The first and third magnets will be 0.25 m long with 10-kG field. LAMPF measurements (Va 80) will be used to determine a stripper thickness ($\sim 200 \mu\text{g}/\text{cm}^2$) to limit H^0 from partial stripping to 0.1% and transmitted H^- even lower. The transmitted H^- will be stripped to H^0 in the third injection magnet.

Fig. 4.2.2-1. Injection system consists of an achromatic beam bump.

A clean dumping system has been incorporated for the unwanted neutral beam. The bump-magnet angles aim the H^0 across the H^+ orbit in the next lattice magnet (HF); in the next straight section sufficient separation has occurred between the H^0 and the central orbit to permit stripping and extraction of the beam through a simple septum magnet into a conservatively rated 2- μ A beam dump.

To reach the limiting tune shift of -0.20 it is necessary to blow up the phase space in the transverse direction from 0.5π mm-mrad to 10.0π mm-mrad. In the longitudinal direction we want an even larger increase. We must do this in six dimensions without large correlations. Our basic idea is to put the beam into the machine with a three-dimensional offset in x , y , and time. The tune in each of the three planes is different ($\nu_x = 7.17$, $\nu_y = 6.30$, synchrotron tune < 0.05). After many turns of injection with the same offset, we will have generated a six-dimensional matched ellipsoidal shell. In the transverse plane the phase space of the injected beam is significant, and a rather thick shell is generated in this four-dimensional sector. In the longitudinal plane the LAMPF beam is almost a δ function, so we will sweep the phase of the synchrotron from zero to the maximum desired offset. An adequate form for the phase change is

$$\Delta\phi = \phi_{\max}(t/t_{\max})^{0.4},$$

where t is the time from start of injection and $t_{\max}$ is the time at the end of injection; $\phi_{\max}$ is the maximum offset desired. In transverse phase space it is possible that no sweeping is necessary. However, a smoother distribution can be obtained if sweeping is provided. We propose sinusoidal sweeping in the vertical plane, which is straightforward. Horizontal sweeping will be left as an option. Because of the coupling of dispersion to transverse phase space, the horizontal distribution will be smooth even if no sweeping is provided.

Our injection scheme assures that the orientation of the injected beam is matched to the booster. The final phase-space area occupied will be determined by the maximum programmed offsets in x , y , and phase. These parameters are only slightly coupled in the coasting beam; hence good control will be possible.

In six dimensions the generated beam is contained inside an ellipsoid. However, the pattern observed in x - y space at any given location is rectangular. This happens because the transverse dimensions sweep out a Lissajous figure, that of two harmonic oscillators with different frequency. A more complicated

sweeping scheme can eliminate filling of the "corners" of phase space. Such injection schemes will be studied in the future.

The injection stripper foil need cover only about half of the circulating beam area, so a given beam particle will cross the stripper only half the time. The accumulation of beam during injection is linear in time, hence the number of hits on the stripper is one-fourth of the number of turns of injection. For the 762-turn injection process from LAMPF, this means the beam must cross the stripper an average of 191 times.

In 191 intercepts of the $200\text{-}\mu\text{g}/\text{cm}^2$ carbon foil, the beam will be multiple-scattered $\sim 0.24\text{ mrad}$ ($1/\sqrt{2}$ the value in free space because the particles are contained in a focusing system and half of the increase in phase space goes into size). This scattering is smaller than the divergence of a matched $0.5\pi\text{ mm-mrad}$ beam, hence is unimportant. Foil heating in the injection process will be 8 W approximately uniformly in an area of 2 cm^2 . This is much less severe than in the PSR, hence we will be able to use stripper foils developed for PSR. Straggling is negligible, since the mean energy loss per particle is $\sim 0.02\text{ MeV}$, and the rf system makes up these losses.

We have simulated the injection process as described. The resulting phase-space distributions are shown in Figs. 4.2.2-2 and 4.2.2-3. Note that in the 1500 injected particles (2 per turn), there are a few approaching four times the nominal phase-space limit, indicating the necessity for the large safety factor of our beam apertures. Halo scraping will be used to remove such particles.

Fig. 4.2.2-2. Transverse phase space of injected beam.

Fig. 4.2.2-3. Beam cross section and longitudinal phase space after injection.

4.2.3 Acceleration

4.2.3.1 Introduction. This machine has a circumference of 330.83 m, operates at 60 Hz, and accelerates 1.75×10^{13} protons per pulse. In this section we consider rf capture, acceleration, and beam loading.

The 60-Hz magnet repetition rate includes 1 ms for injection, 11.75 ms for acceleration, and 3.92 ms for reset. The time dependence of the ring-magnet field is shown in Fig. 4.2.3-1. This waveform dictates an energy gain/turn given by

$$\Delta E_S(t) = \frac{C\pi f_a}{c} (p^+ - p^-) \sin(2\pi f_a t) , \quad (4.2.3-1)$$

where C is the circumference, f_a is the magnet ramp frequency $f_a = 42.55$ Hz, and p^-/p^+ is the initial/final momentum. The requirement of rapid acceleration to 6 GeV means a peak energy gain of almost 800 keV per turn.

A single rf system delivers an energy gain/turn for the synchronous particle given by

$$\Delta E_S = eV_0 \sin\phi_S(t) , \quad (4.2.3-2)$$

where ϕ_S is the synchronous phase angle and V_0 is the net cavity voltage. We

Fig. 4.2.3-1. Time dependence of booster magnetic field.

take the conservative approach $\phi_s < 38^\circ$, so the rf system must supply a peak $V_0 = 1.3$ MV; this implies 16 cavities of 81-kV peak per cavity.

In the absence of beam loading one has to program V_0 and ϕ_s to satisfy Eq. (4.2.3-2) and to ensure that the bucket area does not decrease and that the small-amplitude synchrotron tune be reasonable (< 0.1). With strong beam-loading the cavities must be detuned off resonance by an angle ψ in order to maintain the required energy gain/turn. The angle is given by

$$\tan\psi = \frac{I_b}{I_0} \cos\phi_s, \quad (4.2.3-3)$$

where I_b is twice the first harmonic of the beam current (at the rf frequency). We remark that $I_b = 1.64 I_{av}$ for a parabolic bunch of full width $\pm 80^\circ$ (out of a maximum possible $\pm 180^\circ$). The current $I_0 = V_0/R_{sh}$, where R_{sh} is the total cavity shunt resistance. The ratio of the power delivered to the beam divided by the power dissipated in the cavities is denoted by P_R and is given by

$$P_R = \frac{I_b}{I_0} \sin\phi_s. \quad (4.2.3-4)$$

By combining Eqs. (4.2.3-2) and (4.2.3-4) we obtain an alternate relation for ϕ_s

$$\sin\phi_s = \left[\frac{\Delta E_s P_R}{R_{sh} I_b} \right]^{1/2}. \quad (4.2.3-5)$$

We wish to keep the beam-loading efficiency $P_R < 1.0$, so for a given R_{sh} another constraint equation for ϕ_s is obtained. Needless to say, Eq. (4.2.3-5) does influence the design R_{sh} for the cavities.

4.2.3.2 rf capture. The LAMPF micropulses will be injected for one macro-pulse (≤ 1.0 ms). The micropulses will be prebunched to a 19.7-ns spacing and synchronously injected into each booster bucket. A hole of 6 buckets will be left for extraction. The rf frequency is 50.3125 MHz and the harmonic number $h = f_{rf}/f_{rev} = 66$ for the booster. An rf capture voltage of 600 kV/turn provides synchrotron motion and bucket area. The injected particles occupy $\pm 85^\circ$ of rf phase with bunch area 0.045 eV and maximum $\Delta p/p = \pm 0.26\%$. The ratio of bucket to bunch area is ~ 2.8 . Figure 4.2.3-2 shows the stationary bucket and inner bunch contour for the booster with 1.75×10^{13} protons in the machine.

Fig. 4.2.3-2. Booster rf bucket at injection time. The region occupied by the beam is cross hatched.

The reduction in bucket area due to longitudinal space charge is small ($\sim 2.0\%$). The small-amplitude synchrotron tune ν_{so} is 0.038.

4.2.3.3 Acceleration. The voltage V_0 and phase ϕ_s were chosen to supply the required $\Delta E_s(t)$ as specified by Eq. (4.2.3-1). Figure 4.2.3-3 shows the time dependence of V_0 , ϕ_s , and v_{so} , where we define explicitly

$$v_{so} = \frac{1}{\beta_s} \left[\frac{\eta h e V_0 \cos \phi_s}{2\pi E_s} \right]^{1/2}. \quad (4.2.3-6)$$

In Eq. (4.2.3-6) E_s is the total energy of the synchronous particle, $\beta_s = p_{sc}/E_s$ and $\eta = 1/\gamma^2 - 1/\gamma_t^2$, where γ_t is the transition gamma for the machine $\gamma_t = 10$. The program was chosen to conserve the bucket area until V_0 slightly exceeded 1.3 MV, then to fix it at 1.31 MV; meanwhile, ϕ_s was chosen to satisfy Eq. (4.2.3-2). After 6.0 ms V_0 is reduced linearly to its final value of 1 MV. The small-amplitude synchrotron tune v_{so} reaches a peak value of 0.048.

Figure 4.2.3-4 shows the time dependence of the rms bunch length θ calculated using the rf program shown in Fig. 4.2.3-3. An envelope equation was written for the time dependence of θ using the prescription of Sacherer (Sa 70). The equation included both external and space-charge focusing forces, and was numerically integrated. Rapid bunching occurs (due to rapid acceleration) and the final value for θ is near 16.5° .

Fig. 4.2.3-3. Booster rf program; V_0 is the cavity voltage; ϕ_s is the synchronous phase angle, and v_{so} is the small-angle synchrotron tune.

Fig. 4.2.3-4. Plot of rms bunch length during acceleration according to the rf program shown in Fig. 4.2.3-3.

4.2.3.4 Beam loading. For the present scenario we choose $P_R < 1.0$; that is, the ratio of the power delivered to the beam to that dissipated in the cavities must not exceed 1.0. This ratio reaches 0.96 only at 7.00 ms into the acceleration cycle if we choose a shunt resistance for each cavity of 25 k Ω [see Eq. (4.2.3-5)], and if we use 16 cavities. Figure 4.2.3-5 shows the time dependence of P_R , ψ , and P_T , where P_T is the total power delivered to the beam and

Fig. 4.2.3-5. Beam loading in the LAMPF II booster; P_R is the ratio of power delivered to the beam to that dissipated in the cavities, ψ is the detuning angle, and P_{TOTAL} is the power delivered by the rf source.

cavities and ψ and P_R are calculated using Eqs. (4.2.3-3) and (4.2.3-4), respectively. The curves in Fig. 4.2.3-5 are slightly pessimistic, since they assume $I_b = 2.0I_{av}$ uniformly; actually, I_b ranges from $1.64I_{av}$ to $1.9I_{av}$ over the acceleration cycle.

The beam-loading treatment assumes that an rf accelerating cavity appears as a parallel RLC circuit driven independently by two VHF transmitters (the generator and the beam). The detuning off resonance is effected by changing the resonant frequency of the cavity by an amount Δf

$$\Delta f = \frac{f_{rf}}{2Q} \tan\psi, \quad (4.2.3-7)$$

where Q is the loaded cavity quality factor and ψ is given by, for example, Eq. (4.2.3-3). One notes that the detuning must start as the machine is being injected with beam and then proceed as indicated in Fig. 4.2.3-5.

4.2.4 Booster Extraction and Abort

4.2.4.1 General requirements. Both the fast-extraction and fast-abort system should be able to operate at the maximum kinetic energy of 6 GeV. They will be similar: each consists of a fast ferrite kicker bending horizontally outward, followed by an extraction septum. These extraction elements are located in 4.1-m drift spaces in the lattice. The septum is located about 102° downstream from the kicker in horizontal betatron phase. The kicker risetimes have to be fast (~ 120 ns) and reach a maximum field of 0.03 T. The septa can each be 3.5 m in length and reach a peak field of 1.0 T at 6 GeV; then the extracted beam will be centered 0.3 m away from the beam center line at the entrance to the next dipole.

The kicker and septum-magnet apertures on the central orbit of the ring will have adequate aperture to pass the maximum expected phase space of $\epsilon = 40\pi$ mm-mrad, which includes tails of the transverse distributions. The maximum beam size in the kickers will be about ± 3.5 cm horizontally by ± 2.5 cm vertically. The septum in each septum magnet can start just over 2.0 cm off-axis horizontally at the upstream end and terminate 3.0 cm off-axis at the downstream end.

4.2.4.2 Fast extraction. The layout of the elements in part of the fast-extraction superperiod is shown at the top of Fig. 4.2.4-1. The kicker and septum locations are indicated. The septum can run at, for example, $B = 1.0$ T. At 6 GeV we expect the "clean beam" emittance to damp to $\sim 2.0\pi$ mm-mrad. In Fig. 4.2.4-1 we further show the monoenergetic horizontal beam envelope through the extraction section. Also shown is the envelope of the beam, which experiences a $+2.5$ -mrad outward deflection at the kicker center; this beam would easily clear the iron septum (not shown). At 6 GeV the kick can be effected with a bend strength $B\ell = 0.0573$ T-m (for example, 164 G for a 3.5-m kicker).

4.2.4.3 Fast-abort system. The components of the fast abort system must be programmed to track the accelerator magnets. The apertures should be able to pass a 40π mm-mrad kicked beam. Figure 4.2.4-2 shows the corresponding beam

Fig. 4.2.4-1. Booster fast-extraction system. The monoenergetic envelopes of the 6-GeV beam in the extraction region are shown cross hatched. The lattice and extraction elements are indicated.

Fig. 4.2.4-2. Booster full-aperture fast-abort system. The monoenergetic envelopes of the 6-GeV beam are shown cross hatched. The lattice and extraction elements are indicated.

envelopes for this emittance. In this case a +4.5-mrad kick is needed to abort the beam. The kicker (assumed 3.5 m long) would have to reach a field of 0.03 T at 6 GeV. The kicked beam will exceed the horizontal good-field aperture of the

downstream D dipoles ($x < \pm 30$ mm); this is all right, since the beam is to be discarded anyway. We add that the succeeding magnets that are part of the abort system must be programmed to be able to steer the beam to a dump. Design of this transport channel is not expected to pose any serious difficulties.

4.2.5 Stability Analysis

On one hand, the high intensity of the LAMPF II booster (1.77×10^{13} protons per pulse) makes stability an important issue. On the other hand, rapid cycling (12.75 ms from injection through extraction) means that many instabilities will not have time to affect the beam emittance. We have examined known sources of instabilities in circular accelerators and identified those that could cause trouble. It is our conclusion at this stage of the analysis, that the booster will be stable against transverse oscillations with a damper and that it will be stable against longitudinal oscillations if a feedback system is added to the rf cavities to stabilize the Robinson instability.

Instabilities can be divided into those that affect single-particle motion and those that depend on the coherent action of all the particles in the beam. The single-particle instabilities are driven by imperfections in the magnetic lattice, like a forced harmonic oscillator. The coherent instabilities are characterized by transfer of beam energy to transverse or longitudinal vibrational modes of the beam through interaction with the beam chamber and other metal objects near the beam. Instabilities with wavelength longer than the bunch length in the bunched beam can usually be controlled by dampers and other feedback systems. Instabilities with wavelength shorter than the bunch length can be avoided by providing a source of incoherence, such as momentum or tune spread. Section 4.2.5.1 treats single-particle instabilities and Sec. 4.2.5.2 deals with coherent instabilities.

4.2.5.1 Single-particle instabilities. The transverse/betatron vibrations of the particles about the closed orbit of the machine are characterized by the horizontal and vertical tunes, ν_x and ν_y , which give the number of oscillations per revolution of a particle going around the ring. The general condition for transverse resonance is

$$n_x \nu_x + n_y \nu_y = p \quad , \quad (4.2.5-1)$$

where n_x , n_y , and p are integers. The order of the resonance is $|n_x| + |n_y|$. The driving force for the resonance is any imperfection of the magnetic guiding field or rf cavities. The force seen by the particle is periodic and can be written as a Fourier series,

$$F(t) = \sum_{p=-\infty}^{\infty} F_p \exp(ip\omega_0 t) \quad , \quad (4.2.5-2)$$

where ω_0 is the circulation frequency. The integer p in this summation is the source of the integer p in the resonance condition above. If the ring has superperiodicity n ($n = 8$ for the booster) and the same imperfection or driving term appears in each cell, then the Fourier coefficient F_{kn} (k is another integer) will be large. This makes the superperiodicity resonances $p = kn$ of special concern.

From practical experience the resonances of major concern are superperiodicity resonances of order less than fifth. These were shown in Fig. 4.2.1-1 above. The working point $(v_x, v_y) = (7.17, 6.30)$ is far from all superperiodicity resonances except the third-order difference resonance $2v_x - v_y = 8$. If we assume a beam of uniform density in transverse cross section, then the worst Laslett tunes shifts are $\delta v_x = -0.07$ and $\delta v_y = -0.20$. These tune shifts occur shortly into the acceleration cycle and become smaller as the particle energy increases. The shaded area in Fig. 4.2.1-1 is an estimate of the maximum tune spread of the beam for realistic density distributions. Because of synchrotron motion, particles in the beam will cross the fourth-order resonances $v_x + 3v_y = 26$ and $4v_y = 25$, which are not superperiodicity resonances. The growth times for these resonances are much larger than the 12.75-ms injection plus acceleration time, as we will show below.

The resonance phenomenon in accelerators is complicated by the fact that the betatron frequency becomes amplitude dependent near resonance. If the betatron frequency shifts toward the resonant frequency, the particles are quickly lost. If the frequency shift is away from the resonant frequency, then the particles reach a stable amplitude, although this amplitude may still be larger than the size of the beam pipe. It is common practice to calculate resonance widths. Particles with a frequency difference

$$e\omega_0 = |p - n_x v_x - n_y v_y| \omega_0 \quad (4.2.5-3)$$

smaller than the resonance width are considered unstable.

Although the resonance width estimates the region of instability, it does not determine the beam amplitude growth near the resonance. The best way to investigate this effect is through computer simulations. We have done this to some extent in the tracking studies discussed in Sec. 4.2.1, but further studies are needed.

Resonance widths depend on the periodic nonlinear driving forces, which in turn depend on the size of the multipole errors in dipoles and elsewhere in the ring. These errors can be systematic or random. The systematic errors enhance the superperiod resonances, whereas the random errors apply to all resonances.

Another way to access the importance of these resonances is to examine growth times relative to the cycle time of the accelerator. It is convenient to calculate the number of turns required for a 10% increase in beam size. If this is larger than the $\sim 10\,000$ turns for the beam to reach extraction, then there is no problem.

In view of all these factors, we have done the following calculations. We have examined all resonances up to fifth order in the vicinity of the working point. If the resonance is a superperiodicity resonance, we have assumed that it is driven by a 0.1% systematic error (measured at 5 cm horizontally from the center line) in the dipoles of multipolarity appropriate to the order of the resonance, and we have calculated the resonance width using the formulas found in Guignard's standard treatment (Gu 70), (Gu 77). If the resonance is not a superperiodicity resonance, we have assumed a 0.01% random, multipole error in all magnets at 5 cm. In those cases where the resonance lies within the tune-spread area, we have calculated the number of turns for a 10% blowup in beam size, again using formulas from Guignard for repeated crossing of a single resonance. The results are summarized in Table 4.2.5-I. From the table we see that no resonance is within a resonance width of the working point and that those resonances within the tune spread have growth times greater than 10 000 turns.

In addition to the betatron resonances, one must also consider synchro-beta-tron resonances, which are defined by the resonance condition

$$n_x v_x + n_y v_y + n_s v_s = p \quad , \quad (4.2.5-4)$$

where n_x , n_y , n_s , and p are integers; v_s is the synchrotron tune; v_x is the horizontal betatron tune; and v_y is the vertical betatron tune. The effects of coupling between longitudinal and transverse oscillations in particle beams were first observed in 1971 (Cr 71). More recently satellite resonances due to this coupling have limited the current on DORIS at DESY (Pi 80).

TABLE 4.2.5-I

WIDTHS AND GROWTH TIMES FOR RESONANCES NEAR THE WORKING POINT^a

n_x	n_y	p	e^b	widths δe	Growth Times (turns)
1	3	26	0.07	0.3×10^{-3}	275×10^3
0	4	25	0.20	0.23×10^{-3}	370×10^3
0	5	31	0.50	0.4×10^{-4}	2×10^6 ^c
1	4	32	0.37	0.6×10^{-4}	----
2	3	33	0.24	0.6×10^{-4}	----
3	2	34	0.11	0.6×10^{-4}	----
4	1	35	0.02	0.73×10^{-4}	----

(a) $xv_x = 7.17$, $v_y = 6.30$

(b) $e = n_x v_x + n_y v_y$

(c) The present computer program is not set up to do fifth-order cases when both n_x and n_y are nonzero. Clearly the growth times for these cases must be comparable to the $n_x = 0$, $n_y = 5$ case, since the widths are approximately the same for all fifth-order resonances.

The importance of these resonances depends on the strength of the coupling and the sign of n_s . There are three sources of coupling: dispersion, chromaticity, and the discrete nature of the rf-cavity distribution. Dispersion is usually the most important source of coupling. Piwinski and Wrulich (Pi 76) have shown that if the dispersion functions $\eta(s)$ and $\eta'(s)$ do not vanish at the cavities, the phase advance of the horizontal betatron motion can become complex, leading to unstable betatron motion when

$$v_x \pm v_s = Nk, \quad (4.2.5-5)$$

where N is the number of superperiods and k is an integer. For the booster this condition can never be satisfied, because $v_x = 7.17$, $N = 8$, and v_s is never larger than 0.04.

On the other hand, because the synchrotron motion is nonlinear and in the presence of chromaticity-correcting sextupoles, the betatron motion can be nonlinear, one must be concerned about the growth rates of higher-order resonances given by Eq. (4.2.5-4). The growth rates per revolution of these resonances have been estimated for the $n_y = 0$, $n_x = 1$ case by Suzuki (Su 82) for various values of n_s and appear to be reduced in strength by a factor of

$$\sim \frac{\phi_{\max}^{(n_s-1)}}{2^{(n_s-1)}(n_s^2-1)(n_s-1)!}$$

with respect to the $n_s = 1$ resonance, where $\phi_{\max}$ is the maximum synchrotron amplitude. The lowest resonance of concern to us is the $n_s = 4$ for $v_s \sim 0.04$. By using Suzuki's theory we have been able to show that the growth time for the $n_s = 4$ resonance is very large compared to the 12.75-ms cycle time because of the high order ($n_s > 4$) and because the dispersion functions $\eta(s)$ and $\eta'(s)$ at the rf cavities are small.

The betatron-synchrotron coupling caused by chromaticity is driven by magnetic imperfections. The general condition given in Eq. (4.2.5-4) applies. Because of space-charge tune spread, the synchrotron sidebands of the fourth-order resonances $v_x + 3v_y = 26$ and $4v_y = 25$ are of some concern. The effect of these resonances can be reduced by correcting the chromaticity to zero and keeping the magnetic imperfections to less than one part in 10^{-4} on orbit.

The momentum oscillations caused by the discrete cavity distribution, which caused the resonance problems in NINA (5 cavities and a $v_x = 5.25$), should have negligible effect for the booster (16 cavities and $v_x = 7.17$).

Perhaps the best way to assure oneself that the synchrobetatron sidebands have no effect is to do computer simulations. We have a program called ARCHSIM that will do this and some simulations have been done using 2 rf cavities instead of 16. The beam appears stable throughout the acceleration cycle.

4.2.5.2 Coherent instabilities. A beam of charged particles circulating in a synchrotron interacts with its metallic environment (beam chamber, rf cavities, beam-position monitors, inactivated kicker magnets, etc.) to create electromagnetic wake fields that can act back on the beam. Coherent effects arise that can distort the charge distribution or produce unstable oscillations of the charge distribution. The thresholds for these effects, which are dependent on the beam intensity, have been examined. By using state-of-the-art construction techniques for the beam chamber, rf cavities, etc., we find that coherent instabilities will not prevent us from obtaining the beam current we desire.

In the last 15 years accelerator theory has provided a framework in which the qualitative and quantitative aspects of these coherent instabilities can be explained. The strength of the coupling between the beam and the accelerator environment is characterized by two frequency-dependent impedances $Z_{\parallel}(\omega)$ and $Z_{\perp}(\omega)$ (Wa 80), (Ru 81), and (Ru 83). Recent observations of coherent effects in accelerators has shown stability thresholds consistent with a broad-band resonator model of the impedances. It is still difficult to calculate impedances ab initio and therefore to estimate instability thresholds for new accelerators.

The approach taken here is to use the threshold criteria to put upper limits on the impedance functions and then show that with careful construction it should be possible to keep the impedances below these limits. For our purposes coherent instabilities will be divided into five classes: single bunch (longitudinal and transverse), coupled bunch (longitudinal and transverse), and the Robinson instability, which is caused by beam loading of the rf cavities. The Robinson instability is really a longitudinal, single-bunch instability, but we are treating it separately because of its severity.

The stability limits for longitudinal and transverse single-bunch instabilities are (Ru 83), (Gu 77)

$$\left| \frac{Z_{\parallel}(\omega)}{n} \right| < \frac{3.75 \times 10^9 \beta^2 \gamma |\eta|}{I_p} \left(\frac{\Delta p}{p} \right)^2 \quad (4.2.5-6)$$

and

$$|Z_{\perp}(\omega)| < \frac{2.95 \times 10^9 \beta^3 \gamma \Delta v_{\perp}}{I_p \bar{\beta}_{\perp}}$$

where

$$Z_{\perp} = iR \left[\frac{Z_0}{\beta^2 \gamma^2} \left(\frac{1}{a^2} - \frac{1}{b^2} \right) - (1 + i) \frac{\sqrt{2RZ_0}}{\sqrt{\beta(n-Q)\sigma}} \frac{1}{b^3} \right], \quad (4.2.5-7)$$

where $\omega = n\omega_0$ with ω_0 the angular revolution frequency, I_p is the peak current in the bunch ($I_p \sim N_b e \beta c / \sqrt{2\pi} \sigma_\ell$); σ_ℓ is the rms bunch length; E is the particle's total energy ($T + mc^2$) in electron volts; $\Delta p/p$ is the longitudinal momentum spread (half width at half max); $\Delta v_\perp$ is the spread of the transverse betatron tunes; η is the quantity $(1/\gamma_{tr}^2 - 1/\gamma^2)$; $\beta_\perp$ is the average betatron function; and e is the magnitude of the proton charge (1.6022×10^{-19} coulombs). The tune spread is given by

$$\Delta v_\perp \cong |(n - v_y)\eta + v_y \xi| \frac{\Delta p}{p} + \Delta v_{oc} \quad , \quad (4.2.5-8)$$

where v_y is the vertical tune, ξ is the corresponding chromaticity, Δv_{oc} is the tune spread, which can be introduced by a zeroth-harmonic octupole distribution, and n is every integer greater than the tune. The $Z_\parallel/n$ limit does not depend on n , but the $Z_\perp$ does. Our main concern is with microwave instabilities that have wavelengths less than the rms bunch length; therefore we will evaluate n at

$$n = \frac{\omega}{\omega_0} = \frac{2\pi R}{2\sigma_\ell} \quad , \quad (4.2.5-9)$$

where σ_ℓ is the rms bunch length. Coherent oscillations with wavelengths longer than the bunch length can presumably be damped electrically (see Sec. 4.4.3).

Although many of the parameters in the above formulas vary during the acceleration cycle, it is easily seen that the extreme values occur near injection and extraction. Values of these parameters and the resulting limits on $Z_\parallel/n$ and $Z_\perp$ for the booster are displayed in Table 4.2.5-II.

These requirements are not much more severe than those on existing synchrotrons. LAMPF II obtains high intensities by rapid cycling rather than high circulating currents. Note also from Eq. (4.2.5-6) that longitudinal instability necessarily occurs at the transition energy ($\eta = 0$) in synchrotrons that cross transition. The LAMPF II booster always operates below transition.

In older synchrotrons (AGS, CERN-PS, CERN-ISR), which made no special effort to reduce impedances during construction, one finds that $|Z_\parallel/n| \approx 10-100 \Omega$ has been obtained, and the resulting coherent instabilities significantly limit beam intensity. In more recent synchrotrons designed for minimal impedance (PEP, PETRA) the longitudinal impedance $|Z_\parallel/n|$ is $\sim 1 \Omega$, and hence microwave instabilities are less of a limitation. For LAMPF II a complete inventory of impedance sources will have to be made before construction to minimize this source of instabilities (Ru 79), (Ng 83), and (Gu 77). In addition, it may be necessary to increase the tune spread and momentum spread to decrease the coherence of the beam.

TABLE 4.2.5-II

IMPEDANCE LIMITS FOR LAMPF II BOOSTER

Parameters ^a	At Injection	At Extraction
E	1.735 GeV	6.938 GeV
$\Delta p/p$ (rms)	0.164%	0.09%
σ_z (rms)	0.711 m	0.273 m
$ \eta $	0.283	0.0083
I_p	6.69 A	20.5 A
n	277	612
$\Delta v_{\perp}$	0.133	0.0219
$ Z_{\parallel}/n <$	690 Ω	11.5 Ω
$ Z_{\perp} <$	19.2 M Ω /m	6.8 M Ω /m
$b^2 Z_{\perp} (2R) <$	264 Ω	93 Ω

^aOther parameters: $\gamma_{tr} = 10$, $N_b = 2.95 \times 10^{11}$, $v_x = 7.17$, $\xi = -1.0$, $\Delta v_{oct} = 0.02$, circumference = 330.883 m, $\beta_x = 10.7$ m, and $b = 0.038$ m.

There are eight obvious sources of impedance: (1) space charge and inductive-wall interaction, (2) resistive-wall image currents, (3) higher order rf-cavity modes, (4) resonances in kicker magnets, (5) beam monitors, (6) bellows, (7) steps in the beam-chamber wall, and (8) vacuum ports. The first two sources generate rather broad-band impedances, and therefore it is important to examine their low-frequency behavior as well as their high-frequency behavior. Usually low-frequency instabilities can be controlled with a damper, but if the growth rates associated with low-frequency coherent instabilities were to become too large, one might not be able to construct a damper to handle the instability. We have made a rough evaluation of both the longitudinal and transverse impedances from these two sources; the results are summarized in Table 4.2.5-III. This table shows that the the total longitudinal impedance from space charge and resistive-wall effects is less than the longitudinal impedance limit, and therefore the beam will be longitudinally stable at low frequencies. This is not true for the transverse motion. Although the resistive-wall impedance is less than the transverse impedance limit, the space-charge impedance is too large. The total transverse impedance is close enough to the limit that a low-power transverse damper should have no difficulty in stabilizing the beam (see Sec. 4.4.3).

Table 4.2.5-IV shows the corresponding impedances at the beginning of the microwave instability region (wavelength $\sim$ bunch length). Once again, the longitudinal impedance is much lower than the stability limit. It would appear from this table that transverse stability is violated by a factor of 2. Clearly, further analysis of this problem is necessary, but there are many ways to make improvements in stability (Zo 83). One way to improve the microwave

TABLE 4.2.5-III

LOW-FREQUENCY IMPEDANCES FOR THE BOOSTER^a

	<u>At Injection</u>	<u>At Extraction</u>
T (GeV)	0.797	6.0
$ Z_{ }/n _{sc}\Omega$	173.0	8.91
$ Z_{ }/n _{rw}\Omega$	0.36	0.36
$ Z_{ }/n _{tot}\Omega$	174.0	9.06
$ Z_{ }/n _{lim}\Omega$	690	11.5
$ Z_{\perp} _{sc}M\Omega/m$	39.8	9.87
$ Z_{\perp} _{rw}M\Omega/m$	0.18	0.18
$ Z_{\perp} _{tot}M\Omega/m$	39.9	9.95
$ Z_{\perp} _{lim}M\Omega/m$	4.534	8.14

^a $n = \omega/\omega_0 = 8$; $f_0 = \omega_0/2\pi \sim 0.9$ MHz.

TABLE 4.2.5-IV

HIGH-FREQUENCY IMPEDANCES FOR THE BOOSTER

	<u>At Injection</u>	<u>At Extraction</u>
n	277	612
T GeV	0.797	6.0
$ Z_{ }/n _{sc}\Omega$	173.0	8.91
$ Z_{ }/n _{rw}$	0.32	0.02
$ Z_{ }/n _{tot}$	173.3	8.93
$ Z_{ }/n _{lim}$	690	11.5
$ Z_{\perp} _{sc}M\Omega/m$	39.8	9.87
$ Z_{\perp} _{rw}$	0.002	0.002
$ Z_{\perp} _{tot}$	39.8	9.88
$ Z_{\perp} _{lim}$	19.2	6.76

stability is to increase the radius of the beam. The transverse space-charge impedance depends on $(1/a^2 - 1/b^2)$, where a is the beam radius and b is the chamber radius. For example, if the beam radius were increased by a factor of 1.5, the transverse space-charge impedance would be reduced by a factor of 2.7 at injection and 2.3 at extraction, which would do the job. The stability limits can be increased by changing the chromaticity and/or the proposed octupole tune shift. For example, changing the chromaticity from $\xi = -1.0$ to $\xi = -3.0$ at extraction will change the stability limit to $10.73 \text{ M}\Omega/\text{m}$, which is adequate.

The other sources of impedance mentioned above can be controlled by modern construction techniques. Low-impedance button beam monitors can be used because the booster current is large. Bellows and vacuum ports can be shielded. Abrupt steps in the chamber wall can be avoided by inserting tapered sections. At frequencies below 100 MHz the impedance of kicker magnets can be reduced by the introduction of image-current carriers that do not affect the magnetic flux (Gi 83). Because of the need to control the Robinson instability and other transient effects, we will have to build a rather sophisticated feedforward and feedback control system for the rf power. The impedance of the cavities will be determined by the control system. The transverse coupled-bunch instabilities, which are driven by transverse rf-cavity modes and by the low-frequency part of the resistive-wall impedance, have typical growth times of several hundred to several thousand turns. Because the booster is rapid cycling, most of these modes can be ignored. The faster modes can be handled by the transverse damper.

The last major instability of concern is the so-called Robinson instability. As the beam passes through the rf cavity it will induce a voltage in the cavity that will cancel part of the voltage from the rf power supply. As a result the beam will not be accelerated as it should unless the cavity impedance is modified (detuned) in such a way as to compensate for this beam-loading effect. For stability, the detuning phase angle must satisfy Robinson's criterion

$$2V_s |\cos\phi_s| > I_b R_{sh} |\sin 2\psi| > 0, \quad (4.2.5-10)$$

where the detuning angle ψ is given in terms of the change in cavity resonant frequency by

$$\tan\psi = \frac{\Omega - \omega}{\alpha} = \frac{2Q(\Omega - \omega)}{\Omega}, \quad (4.2.5-11)$$

or else the beam will be unstable against small perturbations in the rf power supply or the beam current (Gr 82). In the above equations $eV_s \sin\phi_s$ is the change of particle energy per turn required to keep up with the acceleration program. The quantity I_b is the largest component in the Fourier decomposition of the beam current; it is roughly twice the average circulating current, but depends somewhat on the width of the bunches. The shunt impedance R_{sh} , the quality factor Q , and the resonant frequency $\Omega = 1/\sqrt{LC}$ are parameters character-

izing the rf cavity. Griffin has shown that the Robinson stability criterion is equivalent to requiring that the ratio of the power delivered to the beam to the power dissipated by the rf cavity,

$$P_R = \frac{(V_s I_b \sin \phi_s)/2}{(V_s^2/2R_{sh})}, \quad (4.2.5-12)$$

must be less than one or else the moving rf-bucket area vanishes and beam intensity is lost until P_R becomes less than one again (Gr 82). If for safety's sake one requires that P_R be 0.5, then there is a tremendous amount of rf power that must be supplied by the rf system. At the time of maximum acceleration the beam power is 2.0 MW. A ratio of 0.5 implies a total rf power of 6.0 MW. This is achievable with the proposed rf amplifier discussed in Sec. 4.2.7 below.

Fermilab has been able to achieve power ratios close to unity by using carefully designed feedback systems (Gr 75). The CERN-PS has a rather complicated rf feedback system that can still exhibit instabilities under certain conditions (Pe 75). Analysis of the feedback system for the Daresbury SRS high-current electron storage ring suggests that with proper design it might be possible to achieve power ratios of nearly two (Be 82). Our own analysis (Wa 84), based on the analyses of Griffin (Gr 82) and Bell et al. (Be 82) suggests that obtaining a power ratio of one is technically possible. Perturbations of the system lead to stable but slowly damped oscillatory motion. To date we have studied only simple RC feedback loops with one corner frequency. We plan to study more complicated systems later.

4.2.6 Magnets and Magnet Power Supplies

4.2.6.1 General requirements. The booster is a 60-Hz, 6-GeV combined-function, rapid-cycling synchrotron. The dipole magnets have to be designed to operate with low ac losses and provide adequate magnetic-field quality in the gap over a range of 0.239-1.1234 T. The dipoles will have laminated cores of relatively high-silicon steel. The cross section of the booster dipoles is shown in Fig. 4.6.2-1.

There will be trim quadrupoles, sextupoles and perhaps octupoles; their construction is fairly standard and will not be discussed. There will also be injection, fast-extraction and fast-abort kicker magnets as well as septum magnets. The parameters for these kickers and septum magnets were given in Secs. 4.2.2 and 4.2.4. A preliminary design of a fast-extraction magnet system is given below. At this time there has been no design work on the injection kicker and the septum magnets, but it appears that corresponding magnets used for the Los Alamos PSR and the ZGS at Argonne (Fo 79) would be adequate for our purposes.

The power supply for the dipoles is an innovative design, replacing the "standard" dc-biased sine-wave excitation (which would be easier to implement) with a shaped waveform giving the three advantages of an extended flatbottom for injection; a slower ramp-up, thus reducing peak rf power needed; and a faster ramp-down, which increases the useful time per cycle.

Fig. 4.6.2-1 Simplified cross section of LAMPF II booster dipoles. Dimensions are given in cm.

4.2.6.2 Dipole magnets. The dipole magnets for the booster need a specified gradient magnetic field over an aperture 4.42 cm high by 11.08 cm wide for the focusing dipoles (HF), and 6.9 cm high by 6.06 cm wide for the defocusing dipoles (HD). The gradients are 3.258 T/m and 3.035 T/m, respectively. The field rises at 42.55 Hz and falls at 127.5 Hz. Thus, the core must be laminated 3% silicon steel to keep eddy-current and hysteresis losses low. For construction simplicity, the dipoles will have parallel ends and be built to the radius of the central trajectory (as were the PSR magnets (Ha 83)). With an inorganic interlamination insulation the 36F145(M15) steel costs about \$64/cwt, and the core losses will be about 1.7 W/kg. The coil must also be designed to minimize eddy-current losses (keeping conductor cross sections small and out of high-leakage-flux regions).

The dipole effective length for the proposed lattice is 2 m (5.625° bend angle). The ring has 64 of these dipoles. The cross section is shown in Fig. 4.2.6-1. Each dipole will weigh about 8 tons and is made to a radius of 20.372 m; this is something that we know how to do with some precision from construction of the PSR (Fo 79). The usual way to design the pole-tip shimming for field uniformity is to use MIRT (one of the POISSON group codes). Our PSR experience ($dBy/By < 10^{-4}$) indicates that MIRT overshims poles. Because the booster is very intolerant of sextupole errors, corrections applied to the pole-tip shape need to be confirmed by measurements on a prototype.

Coil construction will be conventional, directly water-cooled copper conductors with conductors paralleled to keep the inductance within power-supply limits. Individual conductors will be small ($7 \text{ mm}^2 \times 3.5 \text{ mm i.d.}$) to keep eddy-current losses down, and the coil is kept away from fringe fields of more than 0.5 T for the same reason. Conductors are paralleled to match the power-supply current limit (5 kA), then magnets are put in a series until the voltage limit (8.2 kV) is reached. Table 4.2.7-I gives a list of the operating conditions of a turn in the highest leakage field. Insulation of the coil will be adequate to withstand 10 kV to the core without using sheet mica, which has the potential of causing gas-retention problems under radiation.

TABLE 4.2.6-I

DESIGNED OPERATING CONDITIONS FOR THE "HOTTEST" TURN
OF A 30-TURN MAGNET COIL

Length of one turn	5.0 m
Peak/RMS current	1000/613 A
Resistance at 2020°C	2.27 mΩ
Peak resistive voltage/turn	1.92 V
Outside dimensions	7 mm sq
Round hole i.d.	3.5 mm
Current density	16 A/mm ²
I ² R Power	1.0 kW
Max. eddy-current power	1.12 kW
Temperature differential	12°C
Cooling-water flow rate	2.6 L/S
Pressure differential	5.6 atm.
Flow velocity	4.45 m/s
Temperature drop between copper and water	1.5°C

4.2.6.3 Kicker-magnet design. This section describes the kicker-magnet systems for beam extraction and fast abort. The heart of these systems is a pulsed-magnet module consisting of a low-inductance ferrite magnet, an energy-storage capacitor or pulse-forming network, and a thyatron switch. Each module is a self-contained kicker magnet 0.875 m long and capable of reaching a peak magnetic induction of 0.03 T in 120 ns. A single module produces a field integral of 0.026 T-m. Larger values may be obtained by cascading modules.

The size of the individual magnet module is set by the largest value of the inductance that is practical to drive. The 120-ns turn on time to full field determines the voltage the pulser must supply across the magnet terminals. In order to keep the size of the pulser reasonable and to ensure reliable operation, we want to keep the charging voltage as low as is practical. For the magnet module described, the charging voltage is less than 30 kV, which corresponds to a current of 1.2 kA in the magnet.

The proposed-low-inductance pulsed-magnet module design is shown in Fig. 4.2.6-2. It consists of a box formed from two rectangular conductors and two field-shaping ferrite pole faces. The ferrite is necessary to provide the required field uniformity between the conductors. The magnet aperture is 80 mm (horizontal) by 60 mm (vertical) and 3.5 m long.

The dominant parameter in the design of a pulser to drive a module is the turn-on time from zero to full field strength. For a given turn-on time T and current I , the voltage that must be supplied across a magnet's terminals must exceed LI/T , where L is the magnet's inductance. To assure high reliability in the pulse power system, we want to keep the charging voltage under 30 kV. To meet both of these constraints, we have limited the inductance by limiting the magnet length. Each module is 0.875 m long, has an inductance of approximately 1.75 μ H, and requires 1.2 kA to produce a 0.03 T field.

Fig. 4.2.6-2. Simplified schematic of kicker-magnet module.

4.2.6.4 Extraction septum. A possible septum magnet for the fast-extraction and fast-abort systems is shown in Fig. 4.2.6-3. This septum is best run with the same waveform as the ring magnets, since stray fields will be smaller than in the dc case.

4.2.6.5 Power supplies. More efficient beam capture during injection and a significant reduction in rf power requirements during beam acceleration can be achieved for the accelerators of LAMPF II with ring-magnet currents that have a flatbottom followed by a dual-frequency cosine wave. As illustrated by the solid curve in Fig. 4.2.3-1, a constant injection field is followed by a rising cosine field of frequency f_1 during acceleration. Fast extraction takes place at the crest of this wave. Thereafter the magnets are reset with a cosine wave of frequency f_2 . At a time t_3 the cycle repeats. The efficiency, defined by the ratio of injection plus acceleration times to the time of one cycle, is $(t_2 - t_0)/(t_3 - t_0)$. It is typically 77% as compared to the 50% efficiency of a conventional dc-biased cosine wave: an improvement of 54%. A power supply with these features is discussed in (Pr 83).

Figure 4.2.6-4 shows a simplified circuit that can generate a magnet current having a flatbottom, a ramp-up cosine wave, a flattop, and a ramp-down cosine wave. Figure 4.2.6-5 is a simplified schematic of the bending magnet arrangement. A separate but similar circuit will drive the quadrupoles, using the same master timer.

Fig. 4.2.6-3. Schematic representation of an ac septum magnet for the LAMPF II booster.

Fig. 4.2.6-4. Schematic of magnet power-supply circuit.

Figure 4.2.6-6 shows oscilloscope traces of current and voltage traces obtained at ANL during a feasibility demonstration, with magnet currents around 150 A.

Fig. 4.2.6-5. Simplified schematic of bending magnet arrangement.

Dual Frequency Operation

Dual Frequency Operation with Flat Bottom and Flat Top

Fig. 4.2.6-6. Oscilloscope traces of prototype power-supply voltages and current. The right figure shows the waveforms with clamps activated to make a flattop and flatbottom; the left figure shows the results for the dual frequency waveform only.

4.2.7 rf System

4.2.7.1 Requirements. In Sec. 4.2.3 we discussed the performance requirements (maximum voltage, maximum power, tune shift, etc.) for the booster. During the 763-turn injection, the cavities must be continuously detuned to compensate for increasing beam loading. This shift occurs during the 1-ms injection time and is a small fraction of the ~20% tune shift that occurs during the 11.75-ms acceleration phase of the cycle. Beam loading during injection should be no special problem. Since extraction occurs in one turn, beam-loading effects are nonexistent or easily handled by retuning the cavity to the generator frequency at extraction.

During injection the voltage per turn must be large enough to ensure capture of the injected beam. A conservative choice is 0.6 MV/turn based on a ratio of 3 for rf-bucket area to beam area in longitudinal phase space. No power is delivered to the beam during injection.

During acceleration the power delivered to the beam rises to a maximum at the middle of the acceleration phase and falls to zero again at the end. The single-proton energy gain per turn at maximum power is calculated to be $\Delta E_s = 0.798$ MeV/turn. If we take the conservative approach by keeping the synchronous angle $\phi_s < 38^\circ$, then the peak gap voltage per turn must be $V = \Delta E_s / (e \sin \phi_s) = 1.31$ MV/turn. This power and voltage are large for a machine of this size, but are achievable by using a large number of cavities. Prototyping discussed below leads us to believe that 81 kV per rf gap and up to 330 kW per cavity are realizable. This implies the need for at least 16 cavities. To be conservative, we will discuss 20 cavities. The spare accelerating capacity could be useful in case it is necessary to kill some coherent instability. One or two cavities could be operated at a frequency different from that of the other cavities. Experience at the FNAL booster has demonstrated that this works. The lattice design contains more than enough space for twenty 1.17-m-long cavities.

As the protons increase velocity, the average current increases from 2.20 to 2.50 A. The additional detuning needed should be no problem. The transients caused by the six empty buckets (needed for the rise of the extraction-kicker magnetic field) cause less than a 5% change in the gap voltage and can be ignored or handled by a fast feedback system like that used on the FNAL booster and main ring.

In summary, we need about 16 cavities, and each cavity must develop 81 kV, deliver 125 kW maximum power to the beam, and be able to change resonant frequency from 50 to 60 MHz.

4.2.7.2 Simple quarter-wave cavity concept. The cavity design chosen is based on a simple coaxial quarter-wave resonator. Figure 4.2.7-1 shows the full-scale booster-cavity design. The tuner is the ring of ferrite and the solenoidal magnetic field coils at the right side of the cavity. The bias field changes the permeability of the ferrite in the space between the coaxial cylinders, and hence changes the effective inductance L of the cavity; the resonant frequency, of course, is $1/\sqrt{LC}$. Table 4.2.7-I gives the cavity design parameters for the booster. Increasing the cavity length would reduce the required bias field variation, but a longer cavity increases the amount of ferrite needed and brings the ferrite closer to the region of high electric field, making it more difficult to install the ceramic gap insulator and coupling loop. The length chosen is near optimum for the booster.

Fig. 4.2.7-1. The 1.17-m long booster cavity. The ferrite rings are 0.32 m i.d. by 0.60 m o.d. and are 0.0254 m thick. The cooling is accomplished via conductive heat transfer to beryllium-oxide discs located between each ring. Beryllium oxide has the property of being an excellent thermal conductor, almost as good as aluminum, though electrically nonconductive. The permittivity of beryllium oxide is high with a value of 10, which reduces losses. The tuner compartment is isolated from the beam-line vacuum by a ceramic insulator. An external magnet supplies the perpendicular bias field.

TABLE 4.2.7-I

DESIGN PARAMETERS FOR THE COAXIAL CAVITY

<u>Booster cavity characteristics</u>	<u>Value</u>
Cavity outer diameter	0.35 m
Cavity inner conductor diameter	0.16 m
Ferrite region outer diameter	0.60 m
Ferrite region inner diameter	0.32 m
Length of entire cavity	1.17 m
Operating frequency	50-60 MHz
Main gap voltage	81 kV
Ferrite permeability variation	3.0-1.4
Ferrite permittivity	14.0
TDK Y-5 ferrite Curie temperature	180°C
Ferrite operating temperature	57°C
Shunt impedance (without ferrite losses)	25,000 Ω
Transmission line impedance of coaxial cavity	40 Ω
Cavity Q (without ferrite losses)	20,000
Operating cavity Q (Q reduced for stability)	1000

The LAMPF II cavity design applies the external magnetic bias field perpendicular to the rf magnetic fields in the cavity and operates the ferrite in the frequency region below the gyromagnetic resonance. Also, by selecting the right ferrite one can obtain a high-Q resonant system, tunable in frequency over 20% and having very low losses in the ferrite material. This is shown dramatically in Fig. 4.2.7-2, which compares ferrite material biased parallel to the rf magnetic field with the same material biased perpendicular.

Fig. 4.2.7-2. Comparison of nickel-zinc ferrite biased perpendicular and parallel to the rf field. This test cavity has a different geometry from the high power test cavity. These are very low-power tests, but clearly demonstrate the improvement in the magnetic Q of the ferrite with perpendicular bias.

Since accelerator application of perpendicular biasing is relatively new, a short discussion of the theory is appropriate here. If a magnetic dipole in a ferrite is subjected to a dc magnetic field H and a small rf magnetic field h perpendicular to the dc field, the rotating vector $(H + h)$ describes a cone. The magnetic moment M precesses around the cone. Gyromagnetic resonance occurs when the radio frequency is synchronized with the precession. It can be shown that this requires

$$H_{\text{res}} = f/2.8 \quad , \quad (4.2.7-1)$$

where H_{res} is the dc magnetic field in oersteds and f is the radio frequency in MHz. The permeability of the ferrite is a tensor of the form

$$\bar{\mu} = \mu_0 \begin{bmatrix} \mu_x - jk & 0 \\ jk & \mu_y & 0 \\ 0 & 0 & 1 \end{bmatrix} , \quad (4.2.7-2)$$

where the dc magnetic field is applied in the z direction and the material is

saturated in the z direction. The permeability of the ferrite has both a dispersive and a dissipative component and can be written as

$$\mu = \mu' - j\mu'' \quad (4.2.7-3)$$

Fig. 4.2.7-3 shows an idealized plot of permeability vs applied dc magnetic field. The dissipative losses are proportional to μ'' . The greatest loss occurs when $H = H_{\text{res}}$. The quantity μ_e is the effective μ for the dispersive part of the permeability μ' . It varies depending on the relation between the direction of wave propagation in a cavity and the dc magnetic field. The cavity frequency is optimum when the dc bias is between H_2 and H_1 . This is because the value of μ'' is lowest. Materials with low values (1.5-2.5 Oe) of (spin-wave linewidth) must be selected. Thus, a large variation in μ_e can be reached under a very low loss condition.

It should be strongly stated that perpendicular biasing is practical only for accelerator applications with ferrites having sufficiently low saturation magnetization. For example, the magnetic Q of nickel-zinc increases dramatically when perpendicularly biased, but the required magnetic field is impractical (3500-5000 Oe). Fortunately, newer materials, such as magnesium-manganese and microwave garnets exist that saturate at 600 Oe and lower. Thus, it becomes practical to operate above saturation (below the gyromagnetic resonance in frequency) with the dc bias field perpendicular to the rf magnetic field. Under these conditions the ferrites have the following properties:

- very high magnetic Q values, $Q = 1000 \dots 10000$,
- very high electric Q values, $Q = 100 \dots 1000$,
- little dependence of the material properties with respect to rf amplitude, tuning speed, and other operating conditions.

Fig. 4.2.7-3. Plot of permeability vs applied dc magnetic field.

4.2.7.3 Ferrite high-power density tests. A ferrite test cavity was constructed to explore the effect realistic power densities within perpendicularly biased ferrite have upon the magnetic Q of the material. The model cavity is shown in Fig. 4.2.7-4. It is a quarter-wave coaxial cavity with ferrite loading on the low-voltage end. The solenoid provides an axial dc magnetic field, whereas the rf magnetic fields of the cavity are in the θ direction.

An example of the results obtained for candidate materials can be seen in the measurements of (TDK-Y5) garnet. Figure 4.2.7-5 shows the Q_{magnetic} of the ferrite as a function of applied external magnetic field. Figure 4.2.7-5 also shows the power density in the ferrite as a function of applied external magnetic field. Figure 4.2.7-6 shows how the Q_{magnetic} varies with power density. The Q_{magnetic} was determined from a transmission-line model of the test cavity. The nominal rf power into the test cavity was 1 KW.

Fig. 4.2.7-4. The ferrite test cavity. The magnetic field was measured to be uniform over the region occupied by the ferrite test rings.

Fig. 4.2.7-5. Results for TDK Y-5 ferrite with Q_{magnetic} , μ_{ferrite} , and power density in the ferrite plotted vs the applied external perpendicular magnetic bias field. These measurements were conducted at realistic ferrite power densities.

Fig. 4.2.7-6. Results for the ferrites plotted as Q_{magnetic} versus power density. These data were taken near the knee of the ferrite Q curve and indicates that there is only a modest decline in Q as a function of power density even in a relatively high loss region.

To determine the power density in the ferrite as a function of frequency, the data from Fig. 4.2.7-5 was weighted by the frequency/voltage function for the booster. Figure 4.2.7-7 gives the spectral distribution of power loss in the ferrite. From this we see that power loss is peaked at the highest frequency (60 MHz), even though the ferrite has the highest loss at low μ . The

Fig. 4.2.7-7. Spectral distribution of power loss in the ferrite tuner. Recall that the ferrite has its greatest losses at a μ of 3.0, but when weighted by the frequency/voltage function for the booster, the net power loss is small when compared to that at the highest frequency (60 MHz).

TABLE 4.2.7-II

FERRITE POWER-DISSIPATION LOSSES OF THE SIX-TOROID CAVITY^a

Ferrite	Saturation Magnetization		Ferrite Dissipation		
	Electric Q	$4\pi M_s$ (Oe)	Mag. mW/cm ³	Elect. mW/cm ³	Total mW/cm ³
TDK G102 Ni-Zn spinel	> 500	5100	169	194	363
TDK G26 Mg-Mn-Al	64	840	118	1515	1633
TDK G139 Mg-Mn-Al spinel	500	1300	81	194	275
TDK Y300 Al garnet	> 1000	300	370	97	467
TDK Y5 Al garnet	> 1000	600	82	97	179
TDK Y1 Al garnet	> 1000	1200	44	97	141

^aTDK Y1 Al and TDK Y5 Al garnets have outstanding properties and require minimal bias fields. Weighting by the frequency/voltage function has been included.

results of this spectral weighting are summarized in Table 4.2.7-II for the various ferrite materials evaluated.

4.2.7.4 Attenuation of spurious cavity modes. As seen by the cavity, the beam consists of a series of sharp spikes of current. These excite a vast spectrum of undesirable lossy modes in the cavity, which need to be suppressed. A suppression technique has been developed at BNL using slotting in the outside wall of the cavity to prevent current from flowing in directions that would result in undesired modes. The fields that cause such modes pass through the slots and into an outer cavity wall made of resistive material where they are damped. BNL has obtained 30-dB suppression of undesired modes via this slotting technique (G1 83).

4.2.7.5 Producing the necessary bias field. It is assumed that the dc bias field H needed to obtain a certain ferrite permeability μ_{mat} can be calculated from the material's saturation magnetization $4\pi \cdot M_s$:

$$\mu_{\text{mat}} = 1 + 4\pi \cdot M_s / H \quad (4.2.7-3)$$

For these one obtains

$$H_{\text{dc}} = 12...48 \cdot 10^3 \text{ A/m} \quad \text{for} \quad \mu_{\text{mat}} = 3.00...1.4 \quad (4.2.7-4)$$

An estimate of the required magnet parameter is given in Table 4.2.7-III. The requirements are higher than the bias requirements for the FNAL booster. For example, each FNAL booster cavity is biased by a 60-kW (30V/2000A) bias supply. With the bias windings mounted outside the rf resonator, one of the main problems occurring in most parallel-biased cavity designs, decoupling of the bias circuit and the rf fields, is avoided. Instead, the problem is whether

TABLE 4.2.7-III

MAXIMUM ELECTRIC DESIGN PARAMETERS FOR THE PERPENDICULAR BIAS SYSTEM

<u>Parameters for Bias Magnet</u>	<u>Value</u>
Number of windings	10
Maximum dc bias field	2200 Oe
Maximum dc current	2000 A
Maximum self inductance	47 μ H
Maximum tuning rate (booster)	20 KHz
Maximum dc voltage	170 V

the fast varying dc bias field can be passed through the resonator walls without unacceptable losses or dissipation by eddy currents. Problems of this kind can be avoided by constructing the cavities outer conductor of stainless steel with a thickness less than the skin depth at the maximum bias field frequency, or by slotting the sides of the cavity. At 20 kHz, the skin depth in stainless steel is about 2.3 mm. With 0.5-mm stainless cavity walls plated on the inside with 0.02-mm copper, eddy current effects should be negligible.

It is possible to calculate how much the magnetic field inside the ferrite will be reduced by polarization of the ferrite and inhomogeneous field of a finite-length magnet. The results of this calculation are shown in Fig. 4.2.7-8. A good field in the ferrite can be obtained by adding laminated

Fig. 4.2.7-8. A POISSON calculation of the magnetic field inside the ferrite-filled tuner region. The field has a uniformity of approximately 15%.

iron pole pieces on each of the ferrite tuner regions. This bias field can be obtained by means of a water-cooled high-current solenoid.

4.2.7.6 Cooling of perpendicular-biased ferrites. The large magnetic and electric quality factors available in microwave ferrites with perpendicular bias reduce the rf losses drastically. However, this advantage is mostly dissipated by raising the gap voltage per resonator to the point at which ferrite losses again become significant.

Due to the relatively low thermal conductivity of ferrites, power losses inside the material lead to high temperature gradients. The $1/r^2$ distribution of the magnetic and electric power densities in a coaxial cavity normally produce high radial temperature gradients. Excessive temperature gradients can cause mechanical stresses that break the ferrite elements. High absolute temperatures can result in thermal runaway within the ferrite material. If the temperature approaches the Curie temperature for the material, the permeability of the ferrite changes and a system failure results. Mean power densities of about 0.2 W/cm^3 are considered to be the maximum values for stable operation. The most common and effective cooling technique is the use of water-cooled copper discs between the ferrite rings; this method must be excluded for a perpendicular-biased system because of the eddy currents induced in the cooling discs, and hence we must use some alternative technique. The ideal solution would be to have a material between the ferrite discs that has good thermal conductivity and is an excellent electrical insulator. Such a material is beryllium oxide, which has almost the same thermal conductivity as aluminum but with an ϵ of 10. A calculation of the ferrite temperature as a function of location within the ferrite discs indicates that a maximum temperature of 57°C will result at power density of 0.25 W/cm^3 . From these estimates one concludes that we are able to operate the ferrite well below the Curie temperatures given in Table 4.2.7-1.

Figure 4.2.7-9 shows a SUPERFISH electric-field calculation of the booster-cavity geometry shown previously in Fig. 4.2.7-1. Notice that there are no

Fig. 4.2.7-9. A SUPERFISH electric-field calculation of the booster-cavity geometry shown previously in Fig. 4.2.7-1. By locating the ferrite in the manner shown, the desired tuning range and ferrite power densities can be obtained.

regions of excessively high electric-field line density that would tend to overheat the ferrite.

4.2.7.7 Design of the accelerating gap. The design of the ceramic gap seal is considered to be one of the most severe problems, and may turn out to be the main limitation on attainable gap voltage. The coaxial cavity design provides space both radially and longitudinally for the gap geometry. The risk of sparkovers can probably be kept low in the vicinity of the gap by keeping the ceramic seal out of the high field region.

4.2.7.8 Beam loading in the LAMPF II cavities. Beam loading becomes significant when the power absorbed by the beam becomes comparable with the power dissipated in the cavity and, therefore, becomes a problem in the LAMPF II machines. One solution is to raise the power dissipation in the cavity walls to a level comparable to the power passed to the beam. The increased losses need not appear in the ferrites, but could be concentrated in the walls of the cavity, which are easier to cool. Additional losses can be obtained by making the cavity walls, especially the inner conductor, out of stainless steel or soft iron. However, because of its poor electric efficiency, this method should be avoided if possible. Another alternative would be to shunt half of the rf power through an external load.

The extent to which feedback loop control methods can be used instead of, or in combination with, additional loading of the cavity is discussed elsewhere.

4.2.7.9 rf-power amplifiers and coupling into the cavity. The rf amplifier supplies both the power to the beam, $P_b = I_b V \sin \phi_s$, and the power dissipated in the cavity, $P_c = V^2 / (2R_{\text{shunt}})$. With feedback control for stability, it should be possible to operate in a mode where $P_c \sim P_b$. Since the maximum $P_b \sim 125$ kW for the booster, we need at least 250 kW from the amplifier. If feedback is not completely successful, we need more power than this. The power requirements of the main ring (~ 260 kW) are higher than for the booster. The same amplifier will be used in both rings, making construction and maintenance costs lower.

Another requirement on the amplifier is a 20% frequency swing from 50 MHz to 60 MHz. Although many amplifiers have been designed with significantly wider frequency bandwidth, the known designs are limited in power to approximately half that of our requirements. We have designed and are building an amplifier

Fig. 4.2.7-10. Block diagram of LAMPF II rf amplifier featuring simplified IPA and driver circuitry, FPA mounted on cavity and transmission-line transformers.

that will deliver a peak power of 330 kW in the frequency range required. A block diagram of the amplifier is shown in Fig. 4.2.7-10. The 200-W solid-state amplifier is a purchased unit. In the following discussions the final power amplifier will be called FPA; the driver power amplifier will be called the Driver; and the intermediate power amplifier will be called IPA. The FPA uses a 4CW 150,000 tetrode. This tube is very similar to the widely used 4CW 100,000E except that it has improved water-cooling characteristics that extend its power capabilities. The output circuit for the FPA is the accelerating cavity itself.

The unique portion of the rf amplifier is found in the circuitry between the 200-W solid-state amplifier up through the FPA input circuit. The requirement of this circuitry is to obtain the needed power gain over the needed frequency bandwidth in as simple a manner as possible. Any amplifying device has a constant gain-bandwidth product. Therefore gain may be traded for bandwidth and vice versa, but when extreme sacrifice of gain for bandwidth is pursued, one usually sacrifices efficiency as well. The Driver and the IPA are operated in this low-gain, low-efficiency, high-bandwidth condition; however, their required power output is less than 2% of the FPA output, and hence the low efficiency of these stages has negligible effect on the overall efficiency of the rf amplifier. The alternative to this simple two-tube IPA-Driver design would be the use of many small tubes in distributed or parallel tube circuits. This latter technique may be the only viable solution where significantly wider bandwidths are required, but our two-tube design represents a much simpler, more reliable and lower tube-cost approach. Furthermore, the IPA and Driver circuits are located remotely from the FPA; that is, they are outside of the beam tunnel. This feature greatly simplifies maintenance problems.

The second unique feature of these stages is the use of transmission-line transformers to match the inherent high-output impedances to their succeeding low-impedance inputs. The advantage of using transmission-line transformers is that they can be tailored to enhance the gain near the frequency extremes, where other parts of the circuit exhibit reduced gain characteristics. Furthermore, transmission-line transformers are significantly less prone to unwanted perturbations in their amplitude vs gain characteristics than other types of transformers.

As previously noted, the FPA output circuit is integrated into the accelerating cavity. By its nature, this cavity must be operated in a very narrow-frequency bandwidth condition and its resonant frequency is made to track the synchronous beam bunch frequency. Therefore, it is not necessary to sacrifice gain for bandwidth in the FPA.

An equivalent circuit model was made to determine the geometry needed to capacitively couple the energy from the rf amplifier into the cavity. The simple model of the amplifier/accelerating cavity is shown in Fig. 4.2.7-11. Ideal performance is defined as the presentation of $700\ \Omega$ real to the rf power-amplifier-tube active region, over the full frequency range from 50 to 60 MHz. The result yields a coupling capacitor value of approximately 20 pf. The needed coupling plate is certainly of a practical geometry for a coupling capacitor. The load admittance presented to the amplifier-output active region appears to confirm strongly the practicality of the proposed capacitive coupling scheme. The imaginary part of the admittance remains insignificant over the frequency range, and the real part remains almost constant. Taking the reciprocal, the load impedance varies from $691\ \Omega$ real at 50 MHz to $713\ \Omega$ real at

Fig. 4.2.7-11. A simple equivalent model to determine the parameters for the capacitive coupling of the amplifier into the cavity. The coupling capacitor needed is only 20 pf (C_4).

60 MHz. This departure from the desired constant value of 700 ohms is so small that it can be considered insignificant.

4.2.8 Booster Vacuum System

4.2.8.1 General. The challenge in the booster vacuum system is to provide a magnet chamber with low eddy-current losses, but with electrical properties that control beam instabilities. This challenge is met by a unique combination of ceramic and thick-film technologies.

4.2.8.2 Vacuum chambers.

The booster magnets will have ceramic vacuum chambers of slip-cast, high-density, high-purity alumina. This provides an excellent vacuum surface, readily bakeable, for the desired 10^{-7} Torr level (Ne 84) with high radiation tolerance in case of beam spill*. The dipole chambers will have short straight sections joined to approximate the 20.4 m radius of curvature of the magnets. Joints will be glass-sealed (400 C° to 1000 C° in air) with the end joints made by brazed-on thin metal flanges which can be edge-welded.

The main reason for the ceramic chamber is to avoid eddy-current heating and eddy-current magnetic field distortion caused by the rapid-cycling magnetic fields. This presents a problem, because one needs some conductivity in the beam chamber to avoid the buildup of static charge. Furthermore, one needs a low-impedance path for the high-frequency image currents necessary for beam stabilization.

These requirements are, to some extent, conflicting, but there are many possibilities available to treat the chosen chamber material. Very thin metallic coatings can be applied inside or outside by sputtering, vacuum deposition or electroless plating. Thick-film technology provides conductive layers from a few to a hundred microns thick, with resistivities from that of silver ($1.5\text{ m}\Omega/\text{sq}/\text{mil}$ thickness) to $10^8\text{ }\Omega/\text{sq}/\text{mil}$; metallo-organics provide sprayable metallic layers 0.05 to 0.2 microns thick. Such films can be applied

*In LAMPF Exp. 381 alumina tubing from Rutherford's SNS was irradiated using 800-MeV protons. Doses up to 1.6×10^{13} protons/cm² (about $7\text{ W}/\text{cm}^2$) did not produce failure.

in various configurations, longitudinal or circumferential strips, for example. Glassy dielectric layers can also be provided. The current proposal is to make a magnet vacuum chamber as shown in Fig. 4.2.8-1 for the dipoles. Some calculated properties of the structure are listed in Table 4.2.8-I.

An interesting extension of the vacuum chamber technology is to create a coil pattern on the surface of the ceramic tube which can be used in the measurement of the magnetic field.

Samples of joints and various surface coatings are in preparation and will be tested in a 60 Hz magnet for measurement of eddy-current losses and field distortion. The magnet vacuum-chamber is the focus of considerable study and experimental work.

4.2.8.3 Pumping system.

The alternatives for pumping the long skinny vacuum chamber of LAMPF II are a distributed pump integral with the chamber, and a series of pumps between components. Though distributed pumps work well (SPEAR reports 2×10^{-10} Torr without baking), these ion pumps will not work in the booster ac magnetic field. Though titanium sublimation pumps would work, the cost of the added complication to a ceramic vacuum chamber would be considerable.

Fig. 4.2.8-1. Booster dipole-magnet vacuum chamber.

TABLE 4.2.8-I

GEOMETRY AND FREQUENCY ($n = 100$) = 76 MHz

Parameter, Units	Symbol	Inside Layer	First Layer	Second Layer
Skin depth @ 76 MHz, m	δ	2×10^{-6}	10^{-5}	10^{-5}
Cut-off freq. $\delta = t$, H_z	fco	2×10^{10}	4×10^7	4×10^7
rf resist., Ω	R_{rf}	600	3.2 long	3.7×10^4 Ω/m circum.
Eddy-current power, W/m	P	11	10	19
Interlayer cap., nf/m	C		4.9	360

We are therefore proposing discrete triode or star-electrode ion pumps, with provision for rapid roughing to 10^{-4} Torr by a few transportable packages, including cryopumps and turbo-molecular pumps. The volume of one superperiod is around 500 ℓ . The target vacuum is 10^{-7} Torr, at which triode ion pumps exhibit over 100% of their rated pumping speed.

In this molecular flow regime, actual pumping speed is limited by the conductance of the pipe, which for air is given by the Knudsen formula

$$C_{air} = 28.7(K') \frac{A^2 B^2}{L(A+B)+2.66AB} \ell/s, \quad (4.2.8-1)$$

where A and B are the pipe height and width in cm (6.2×11), L is the pipe length in cm (one tenth of a dipole is 20), and K' is Clausing's factor (0.982).

For a pump situated between two adjacent dipoles, C_{air} would be about 200 ℓ/s , and a 400 ℓ/s ion pump would be an appropriate size to pump both ways. Using a 1000 hour outgassing rate of 10^{-10} T- ℓ/s -cm², the two adjacent dipoles require 70-80 ℓ/s capacity each.

Clearly, extra pumping capacity has to be associated with devices such as accelerating cavities, and the out-gassing characteristics of all internal surfaces need to be known. More precise pump sizing can be done based on that information. Estimates of beam-induced desorption suggest it will not be a problem; thermal desorption can be reduced by bake-out; and the effects of multipactoring remain to be studied.

4.2.8.4 Requirements of low-impedance vacuum chamber. There are a number of requirements for a low-impedance vacuum system beyond the need for a special vacuum pipe in the magnets. Everywhere in the system special precautions are required. Among these are (1) finger stock inside all bellows so that the beam

sees only a smooth-walled, high-conductivity beam pipe; (2) grids over pumping ports to carry the image current of the beam; (3) smooth transitions in beam-chamber dimensions (that is, size transitions slow compared with the diameter of the pipe); (4) special vacuum chambers inside all kickers and pulsed magnets; (5) beam diagnostic instrumentation designed to present a low impedance to the beam; (6) high-frequency capacitors connecting the longitudinal stripes on the ceramic vacuum chamber with the metal vacuum transitions at one end of each magnet; and (7) an electrical connection between the interior conductive coating of the ceramic chambers and the metal transition pieces.

As at PEP and PETRA, all vacuum-system components will have to meet a specification for impedance. A step in the quality control will be a measurement of the impedance of each component before installation (We 79, Ng 83).

4.3 MAIN-RING DESIGN

4.3.1 Lattice Design

4.3.1.1 Specifications. The main ring is designed to be the highest energy synchrotron that fits on the present LAMPF site. It appears that a racetrack-shaped ring with ~ 1.3 -km circumference has the largest arc diameter that can be accommodated. The racetrack straight sections are efficient for slow extraction; 160 m is a comfortable length for this purpose. The long straight sections are also useful for accommodating rf cavities.

The apertures of the main-ring magnets are conservatively designed. We have allowed sufficient space for a phase space of 8π mm-mrad and dp/p of $\pm 0.3\%$ (the injection phase space is 2π mm-mrad and $\pm 0.15\%$). An additional allowance of ± 1 cm horizontal and ± 0.5 cm vertical has been provided for closed-orbit distortion. The safety factor on the apertures thus obtained will be useful for minimizing beam spills.

We want to avoid crossing transition because this is a potential source of beam loss. With injection energy of 6 GeV, it is possible to design the main ring with transition energy of 5.52 GeV. This, of course, means choosing a relatively low set of tunes ($\nu_x = 7.45$ and $\nu_y = 6.45$) because the transition energy is proportional to the focusing strengths and hence to the tunes. Low tunes and a large circumference lead to uncomfortably high dispersion. To minimize this problem at injection and extraction, one can make the straight sections dispersionless by requiring that the betatron phase advances through the arcs of the racetrack be integral multiples of 2π . In order to minimize dispersion in the arcs, it is necessary to include dispersion-suppressor cells at the end of each arc. Another advantage of making the straight sections dispersionless is that the synchrotron-betatron coupling in the rf cavities is reduced. We chose a tune near $1/2$ to facilitate slow extraction and to work in the region between 0.33 and 0.5, which is free of sum resonances of order less than 5. The tune diagram is shown in Fig. 4.3.1-1.

A machine operating above transition should have positive chromaticity (an unnatural sign) in order to avoid the head-tail instability, which means chromaticity-adjustment sextupoles are required. The arcs are second-order achromats that minimize geometric aberrations caused by the chromaticity-adjusting sextupoles. This choice also makes the machine rather insensitive to sextupole errors in the dipoles.

With the radius of the arc fixed, the energy of the machine is determined by the highest practical dipole field in a rapid-cycling machine. We took this to be 1.3 T. This field permits a maximum energy of 45 GeV, which is acceptable to the users.

The resulting design is similar to that of stretcher rings designed for GeV electron beams. The main components (long straight sections, second-order achromatic arcs, and dispersion suppressors) are the same. The very-high-beta section for extraction (betatron function ~ 350 m at the center) is different. A quarter layout of the main ring is presented in Fig. 4.3.1-2. We show in Fig. 4.3.1-3(a) a plot of the betatron and dispersion functions through a machine superperiod. In Fig. 4.3.1-3(b) we show a plot of the beam sizes expected through the region at 6 GeV.

Fig. 4.3.1-1. Tune diagram for the LAMPF II main ring, showing all resonances through fifth order. Sum resonances are shown as solid lines; difference resonances are dotted. The shaded region is occupied by the beam at injection time.

4.3.1.2 Details of the arc cells. All the arc cells are symmetric about their centers. This design puts beam waists in desirable locations and allows straightforward choices for sextupole locations. It also minimizes the number of parameters required to match the beam between different sections. The arc cells have a phase advance of 90° in the horizontal plane and 79.2° in the vertical plane. This small departure in vertical focusing from the standard second-order achromat design produces no significant effects and allows a choice of tunes that differ by one between the horizontal and vertical planes. Each cell consists of four combined-function magnets (HD, HF, HF, HD), as shown in Fig. 4.3.1-3(a). The field index of the HD (HF) magnet is 80.5 (-85.6). In each arc there are 8 identical cells of the type described above. The total phase advance of these 16 cells is 4π rad in the horizontal plane and 3.5π rad in the vertical.

4.3.1.3 Dispersion-suppressor cells. Two dispersion-suppressor cells are located at each end of the arc and are similar to the arc cells except that each bending magnet is half length. The horizontal and vertical phase advances are the same as in the normal cells. This requires that the strength of the quadrupole component of the combined-function magnets be larger because of the reduced length. The field-index of the defocusing dipole (HD2) is 115.8 and the field index for the focusing dipole (HF2) is -122.4. The amplitude functions are nearly perfectly matched to the other arc cells. The sextupole in these cells must correct not only the chromaticity of each cell but also the chromaticity of the adjacent long straight sections. The strengths of the sextupoles required

1/4 SECTION LAYOUT
LAMPF-II MAIN RING
45 GeV

Fig. 4.3.1-2. Layout of 1/4 of the LAMPF II main ring. Magnets labeled Q are quadrupoles, and those labeled HF and HD are horizontally focusing and horizontally defocusing dipoles, respectively. Sextupoles (not shown) are located at the center of each HF-HF pair and each HD-HD pair.

is 0.021 T for the SD magnet and -0.028 T for the SF magnet at 5 cm from the center of the sextupole. These strengths are readily achievable.

4.3.1.4 Long straight sections. Each long straight section is also symmetric about the center. Each half section consists of a short matching cell and a symmetric high-beta cell. The matching cell contains a quadrupole doublet. The high-beta cell consists of two quadrupole doublets symmetric about the center; the phase advance through this cell is $\pi/2$ rad. The focusing strengths were chosen to make the horizontal-amplitude function small and the vertical-amplitude function large after the first doublet. The result, as seen in Fig. 4.3.1-3, is an amplitude function at the center of a straight section that is 350 m horizontal and 15 m vertical.

Fig. 4.3.1-3. (a) Behavior of indicated functions through a main-ring superperiod. (b) First-order beam sizes expected at 6 GeV.

4.3.1.5 Tracking studies. We have performed tracking studies on the main ring similar to those on the booster. A systematic sextupole field error was put into each dipole with sextupole strength of 0.1% of the dipole field strength at 5.5 cm horizontally out from the center line of the dipole. The result was a large change in the chromaticity that changed from zero to -20. Even this large chromaticity can be corrected by trim sextupoles, but the phase-space plots become quite "fuzzy" after a few hundred turns. This result indicates that we must be careful to keep the sextupole error in the dipole down to much less than 0.1%.

4.3.1.6 Trim quadrupoles and dipoles. The present concept is to use a triple-function trim magnet with separate dipole, quadrupole, and sextupole windings in a suitable location, to be optimized in subsequent studies.

4.3.1.7 Utilization of straight sections. Forty cavities are required. These can be located in any of the drifts except the central high-beta region. There are also four useful drifts of 9.5 m each with moderate dispersion in the dispersion suppressors. One long straight section is used for a slow-extraction system and a fast-abort system. The other straight section will handle injection system.

4.3.1.8 Closed-orbit distortions and alignment tolerances. The theory of closed-orbit distortions (COD) was given in Sec. 4.2.1.6 and will not be repeated here. For the main ring we considered errors due to misaligned quadrupoles and to misalignment and field and length errors in the dipoles.

The main ring is composed of two bend sections that use 88 combined-function dipoles and two long straight sections, each containing 12 quadrupoles. The dipole field at injection is 0.194 T, the radii of curvature $\rho = 111.95$ m. The tunes used were $\nu_x = 7.45$ and $\nu_y = 6.45$, and the mean machine radius used was $R = 210.6$ m. The kinetic energy is 6 GeV ($B\rho = 22.93$ T-m). Equations (4.2.1-6) and (4.2.1-7) were evaluated using the main-ring lattice functions; the results are given in Table 4.3.1-I for misaligned quadrupoles. If we assume transverse displacement errors of 0.1-mm rms, then the misaligned quadrupoles will contribute at most 3.3 mm to both the horizontal and vertical COD in the dipoles.

We next considered the errors in the combined-function dipoles. Equation (4.2.1-6) was again evaluated. We assumed rms transverse placement errors of 0.1 mm, an rms field error $\Delta B/B = 10^{-4}$, and rms length/axial-placement error of 0.05 mm. The results are also presented in Table 4.3.1-I. With 98% confidence it appears that the horizontal COD will not exceed 10 mm in the arcs (dipoles) and that the maximum vertical COD should be about 5 mm in the dipoles.

Our conclusion from these studies is that it will be straightforward to get a coasting beam inside the aperture with routinely available alignment tolerances. Further study is required to decide whether full acceptance will be achieved with an active closed-orbit correction system or with more precise alignment.

4.3.2 Injection

The main-ring injection will be performed in a long straight section (LSS). Figure 4.3.2-1 shows the behavior of the betatron functions β_x and β_y in the 160-m LSS. The system is mirror symmetric about the center and focusing is obtained using 12 quadrupoles. The elements depicted at the top of Fig. 4.3.2-1 include the quadrupoles as well as the basic injection hardware.

The extracted beam from the booster will be transported to the LSS and be injected using a horizontal septum and fast kicker. These elements are centered at 39.37 and 77.41 m relative to the center of the LSS, respectively.

TABLE 4.3.1-I

MAXIMUM CLOSED-ORBIT DISTORTIONS

Errors	Location	x plane (mm)	y plane (mm)
Misaligned quadrupoles	st. section	19.0	5.0
" "	HF dipoles	3.2	1.8
" "	HD dipoles	1.7	3.3
Dipoles	st. section	60.0	7.7
"	HF dipoles	10.0	2.8
"	HD dipoles	5.3	5.2

Fig. 4.3.2-1. Amplitude functions in long straight section of main ring. The injection elements are shown located in the LSS.

Fig. 4.3.2-2. Centroid of the 6-GeV injected beam between the injection septum and kicker.

Figure 4.3.2-2 shows the injected beam path in the LSS section between the septum and kicker. At the center of the septum the coordinates are $x_s = -50$ mm and $x'_s = 1.92$ mrad. The kicker will be 2.5 m in length; at 6 GeV the field

required is about 0.024 T in order to deflect the beam by 2.65 mrad. Since four booster batches are required to fill the main ring, four gaps in the bunch structure are required. The kicker will be turned off during the booster acceleration periods of 16.7 ms. With approximately six missing bunches in each gap, kicker rise and fall times of ~ 120 ns are required. The septum angle will be large enough to bring the injection line in without disturbing the rf-cavity layout. A field of 1.0 T over a 3.0 m length affords a 0.13-rad bend at 6 GeV.

A transport line will match the 6-GeV proton beam from the booster to the main ring. The transverse match required is eight dimensional, since both betatron and dispersion matching is required. The longitudinal match will be done in the booster itself by reducing the final rf voltage to produce the longer bunch required for the main-ring injection.

The match points are appropriately the upstream end of the booster extraction kicker (a) and the downstream end of the main-ring injection kicker (b). Then at (a) we have for the Twiss parameters $\beta_x = 22.3$ m, $\alpha_x = 2.677$, $\beta_y = 4.22$ m, $\alpha_y = 0.303$, the horizontal dispersion $\eta_x = -0.074$ m, and its derivative $\eta'_x = -0.135$. At location (b) we require $\beta_x = 14.182$ m, $\alpha_x = 0$, $\beta_y = 64.285$ m, and $\alpha_y = 0$; the dispersions and their derivatives are zero at this point. The extraction from the booster introduces a dispersion that must be considered in the transport line to the MR injection septum. The horizontal dispersion cancellation must not only remove the booster η_x and η'_x but also overcorrect for the ensuing horizontal dispersion introduced by the MR injection septum and injection kicker.

4.3.3 Main-Ring Acceleration

4.3.3.1 Requirements. This machine has a circumference of 1323.3 m, operates at 3 Hz, and accelerates 7×10^{13} protons per pulse. We consider injection, acceleration, and beam loading. The main-ring rf system must be capable of handling the transients caused by a partially filled circumference (1/4 full to full) during the injection process.

The 3-Hz-magnet repetition rate includes 50 ms for injection, 87.5 ms for acceleration, 166.67 ms for slow extraction (50% duty factor), and 29.17 ms for reset. The magnet waveform, shown in Fig. 4.3.3-1,* dictates an energy gain/turn given by

$$\Delta E_s(t) = \frac{C\pi f_a}{c} (p^+ - p^-) \sin(2\pi f_a t) \quad , \quad (4.3.3-1)$$

where C is the circumference, f_a is the magnet ramp frequency = 5.714 Hz, and

*All beam-dynamics calculations have been done with the sinusoidal waveform. However, recent work indicates that a power supply for a trapezoidal waveform with equal rise and fall times may be cheaper. Essentially the same rf requirements obtain for both systems. The cost estimate is based on the trapezoid system.

Fig. 4.3.3-1. Time dependence of main-ring magnetic field.

p^-/p^+ is the initial/final momentum. The requirement of rapid acceleration to 45 GeV means a peak energy gain of 3.09 MV/turn.

The rf cavities deliver an energy gain/turn for the synchronous particle given by

$$\Delta E_s = eV_0 \sin \phi_s(t) , \quad (4.3.3-2)$$

where ϕ_s is the synchronous phase angle and V_0 is the net cavity voltage. We take the conservative approach $\phi_s < 38^\circ$, so the rf system must supply a peak $V_0 \sim 5$ MV; this implies, for example, 40 cavities for 125-kV peak/cavity.

In the absence of beam loading, V_0 and ϕ_s are programmed to satisfy Eq. (4.3.3-2) and to ensure that the bucket area does not decrease and that the small-amplitude synchrotron tune be reasonable (< 0.1). With strong beam loading the cavities must be detuned off resonance at angle ψ in order to maintain the required energy gain/turn. The angle is given by

$$\tan \psi = \frac{I_b}{I_0} \cos \phi_s , \quad (4.3.3-3)$$

where I_b is approximately 1.8 times the first harmonic of the beam current (at the rf frequency) and $I_0 = V_0/R_{sh}$, where R_{sh} is the total cavity shunt

resistance. The ratio of the power delivered to the beam divided by the power dissipated in the cavities is denoted by P_R and is given by

$$P_R = \frac{I_b}{I_0} \sin \phi_s . \quad (4.3.3-4)$$

By combining Eqs. (4.3.3-2) and (4.3.3-4) we obtain an alternate relation for ϕ_s , which further constrains beam-loading efficiency P_R and cavity shunt resistance R_{sh} ,

$$\sin \phi_s = \left[\frac{\Delta E_s P_R}{R_{sh} I_b} \right]^{1/2} . \quad (4.3.3-5)$$

4.3.3.2 Injection. Four successive batches of 1.75×10^{13} protons from the 60-Hz booster are stacked around the main ring in a period of 50 ms. The main-ring harmonic number $h = 264$ and for the booster $h = 66$, so the injection is synchronous. The batches from the booster must be less than 66 rf buckets in length to allow for the rise and fall times of the main-ring injection and booster extraction kickers; we assume 6 empty buckets between each batch. Thus the 7×10^{13} protons are contained in 240 bunches; this corresponds to $\sim 3 \times 10^{11}$ protons in each bunch.

The rf capture voltage for the 50-ms injection phase is chosen to be 1 MV/turn; the bucket area is 0.9 eVs and amply contains the injected beam as shown in Fig. 4.3.3-2. The cross-hatched contour shown occupies an area of 0.073 eVs with rf phase occupying $\pm 40^\circ$ about the synchronous phase (180°). This bunch length (± 1.875 ns) is longer than one would normally obtain from the booster [Co 84]. Of course, one could increase the capture voltage to get a better match to the booster output, but this causes large values of the momentum spread dp/p . The large dispersion in the arcs ($\eta = 8$ m) requires that we keep $dp/p < \pm 0.15\%$. Therefore it is better to decrease the booster voltage at extraction to increase the bunch length.

The longitudinal space-charge forces are strongly related to the mismatch. At 6 GeV these forces are attractive (focusing) in the booster and repulsive (defocusing) in the main ring. Equivalently, the booster operates below transition energy ($\gamma_t = 10$), whereas the main ring operates above transition energy ($\gamma_t = 6.0322$). The longitudinal space-charge force is proportional to a factor given by

Fig. 4.3.3-2. Main-ring rf bucket at injection time. The region occupied by the beam is shown cross hatched.

$$f_L = \left[\frac{g_0}{2\gamma^2} - \frac{\beta}{Z_0} \left| \frac{Z_L}{n} \right| \right] N, \quad (4.3.3-6)$$

where $\gamma = E/m_p c^2$, $\beta = v/c$, the impedance of free space $Z_0 = 377$ ohms, and $g_0 = 1 + 2\ln(b/a)$, b/a is the ratio of beam pipe to beam diameter (~ 2 at injection), the inductive wall impedance term is given by Z_L/n , and N is the number of protons in a bunch.

Equation (4.3.3-6) indicates $f_L > 0$ at low energies (γ small), but becomes negative above several GeV for representative values for Z_L/n (10-20 ohms). Figure 4.3.3-3 shows a plot of the rms bunch length during acceleration in the 60-Hz booster for two values of Z_L/n (0 and 10 ohms). In this plot the rf voltage has been reduced linearly from 1.3 to 0.28 MV for $6 < t < 12$ ms in order to lengthen the bunch prior to injection into the main ring.

4.3.3.3 Acceleration. The voltage V_0 and phase ϕ_s were chosen to supply the required energy gain $\Delta E_s(t)$ as specified by Eq. (4.3.3-1). Figure 4.3.3-4 shows the time dependence of V_0 , ϕ_s , and A_{bucket} , where we define explicitly

$$A_{\text{bucket}} = \frac{8R\alpha(\Gamma)}{c} \left[\frac{2EeV_0}{\pi h^3 \eta} \right]^{1/2}. \quad (4.3.3-7)$$

In Eq. (4.3.3-7) E is the total energy of the synchronous particle, R is the machine radius ($C/2\pi$), $\alpha(\Gamma)$ is the moving bucket factor, and $\eta = 1/\gamma^2 - 1/\gamma_t^2$, where the transition gamma for the machine $\gamma_t = 6.0322$. The program was chosen

Fig. 4.3.3-3. Time dependence of booster rms bunch length. The bunch curves are explained in the text.

Fig. 4.3.3-4. Main-ring rf program; ϕ_s is the synchronous phase angle and A_{bucket} is the rf bucket area.

to conserve the bucket area until V_0 reached 5.0 MV, then to fix it at 5.0 MV; meanwhile ϕ_s was chosen to satisfy Eq. (4.3.3-2). One notes the bucket area drops to 0.54 eVs during the cycle. This should still be ample to contain the bunch ($\epsilon_L \sim 0.07$ eVs). The small-amplitude synchrotron tune ν_{so} (not shown) varies from 0.0075 at injection to a peak value of 0.019.

Figure 4.3.3-5 shows the time dependence of the rms bunch length θ calculated using the rf program shown in Fig. 4.3.3-4. An envelope equation was written for the time dependence of θ using the prescription of Sacherer [Sa 70]. The equation is written as [Le 71]

$$\frac{d}{dt} \left(\frac{1}{r} \frac{d\theta}{dt} \right) + s\theta + \frac{\Omega f_L}{\theta} - \left(\frac{S}{\pi} \right)^2 \frac{r}{\theta^3} = 0, \quad (4.3.3-8)$$

where $r = -h^2 \eta / (m R^2 \gamma)$, $s = -e V_0 \cos \phi_s / (2\pi h)$, S is the rms longitudinal emittance of the bunch ($S \sim 0.07/6$) eVs, and m is the proton mass. The term Ω is equal to $3r_p m c^2 h / R$, where r_p is the classical proton radius $r_p = 1.54 \times 10^{-18}$ m and f_L is given by Eq. (4.3.3-6). We assume $|Z_L/n| = 20$ ohms for this calculation. Equation (4.3.3-8) was numerically integrated. The final value for θ is 18.3° (or 0.85 ns in time units).

4.3.3.4 Beam loading. For the present scenario we choose $P_R \leq 2.0$; that is, the ratio of the power delivered to the beam to that dissipated in the cavities must not exceed 2.0. This ratio reaches 2.0 halfway into the acceleration cycle if we choose a shunt resistance for each cavity of 80.6 k Ω [see Eq. (4.3.3-5)] and if we use 40 cavities. Figure 4.3.3-6 shows the time dependence of P_R , ψ , and P_T , where P_T is the total power delivered to the beam and cavities and ψ and P_R are calculated using Eqs. (4.3.3-3) and (4.3.3-4), respectively. The detuning angles are large, reaching $\psi = 80^\circ$.

Other options are clearly feasible. For example, use of 50 cavities each with $R_{sh} = 47.7$ k Ω , results in a maximum $P_R = 1.33$ and maximum $P_T = 14.35$ MW; at midpoint ψ drops below 60° in this case.

Fig. 4.3.3-5. The rms bunch length in main ring during acceleration.

Fig. 4.3.3-6. Main-ring beam-loading parameters P_R is the ratio of the power delivered to the beam divided by that dissipated in the cavities, ψ is the detuning angle, and P_{TOTAL} is the power supplied by the rf source.

The beam-loading treatment assumes that an rf-accelerating cavity appears as a parallel RLC circuit driven independently by the rf system and the beam. The detuning off-resonance is effected by changing the resonant frequency of the cavity by an amount Δf , where

$$\Delta f = \frac{f_{rf}}{2Q} \tan \psi, \quad (4.3.3-9)$$

where Q is the loaded-cavity quality factor and ψ is given by, for example, Eq. (4.3.3-3). One notes the detuning must start as the machine is being injected with beam and then proceed as indicated in Fig. 4.3.3-6. Other compensation may be necessary for a partially filled ring.

4.3.4 Extraction and Abort

Both fast abort and slow extraction will be performed in a single long straight section (LSS). Figure 4.3.4-1 shows the behavior of the betatron functions β_x and β_y in the 160-m LSS. The system is mirror symmetric about the center, and focusing is obtained using 12 quadrupoles. The elements depicted in Fig. 4.3.4-1 include the quadrupoles as well as the basic abort and extraction hardware. The fast-abort system is composed of two kickers and a septum. The slow-extraction system consists of perturbing elements (not shown), an electrostatic wire septum (ES), and an iron septum (slow-extraction Lambertson).

Fig. 4.3.4-1. Amplitude functions in extraction LSS. Also shown are positions of fast-abort and slow-extraction elements.

4.3.4.1 Slow-extraction mode. Slow extraction will occur in the horizontal plane using half-integer resonant extraction. A pair of quadrupoles will be used to drive the tune from 7.45 to near 7.50. These quadrupoles are located an odd multiple of π apart on the 15th harmonic and are oppositely powered; this ensures no 0th or 30th harmonic contributions. A similar pair of octupoles introduce an amplitude-dependent tune spread, thus splitting the $x - x'$ phase space into stable and unstable regions. Particles occupying the unstable regions stream out along the separatrix where, for example, they arrive at the ES in a region of nonzero electric field (they have jumped the wires). These particles receive an outward deflection in the ES and arrive displaced at the slow-extraction Lambertson where they can be bent out of the machine.

It was decided to locate the four perturbing elements in centers of appropriate horizontal sextupoles (SF) where $\beta_x = 59.154$ m and $\beta_y = 16.99$ m. This puts them in the regular arcs. Relative to the center of the septum ES, the quadrupoles are at $s = 262.46$ m and $s = 919.72$ m; the octupoles are at $s = 438.91$ m and $s = 1096.16$ m. This puts the quadrupoles in the centers of the 3rd regular bending cell of both arcs and the octupoles in the centers of the 7th regular bending cell of both arcs. This position also puts the quadrupoles and octupoles in phase on the 15th harmonic.

The extraction dynamics were studied assuming the perturbing elements are thin lenses. Initially the quadrupole powers (P_Q) were increased with the octupoles assumed off. The machine went unstable for $P_{Q1} = -P_{Q2} = 0.0054$ m⁻¹. Then we reduced P_{Q1} to 0.00525 m⁻¹ and energized the octupoles to $P_{O1} = -P_{O2} = 3.0$ m⁻³. Figure 4.3.4-2 shows the $x - x'$ phase space at the ES for three values of emittance $\epsilon_{geom} = 0.04, 0.02,$ and 0.01π mm-mrad at 45 GeV. The contours

Fig. 4.3.4-2. The x - x' phase space at the ES for three values of emittance. Three particles are tracked for 100 turns; the outer particle is extracted after jumping the septum at x_s .

represent the coordinates of three particles tracked for 100 turns and plotted for each turn. The outer particle is unstable and is extracted by jumping the ES wires (shown at x_s). We use the data in Fig. 4.3.4-2 to estimate the step size and extraction efficiency for 0.1-mm-thick wires. They both increase with x_s ; for $x_s = 60$ mm we obtain 99.3% efficiency and a step size of 21.0 mm at x_s ; similarly for $x_s = 70$ mm we find an efficiency of greater than 99.5% and a step size of 31 mm at x_s (the inefficiency is the fraction of beam that strikes the wires). The calculation is for a monoenergetic beam with zero phase space, so that the inefficiency will probably be larger than calculated by this procedure.

Particles jumping the ES wires receive about 0.43-mr kick outward—at 45 GeV this is obtained with a field of 5.0 MV/m over a length of 4.0 m. Figure 4.3.4-3 shows the trajectories of three particles in the LSS between the ES and slow-extraction Lambertson.

The unstable trajectories can dictate the horizontal aperture in the accelerator. Figure 4.3.4-4 shows the behavior of two particles tracked through 1/2 of the LSS, the dispersion suppressor, and into the arcs. The beginning coordinates are 60 and 80 mm. The trajectories show the maximum horizontal sizes to be expected on the last turn through the accelerator. The displacements are ± 35 mm in the arcs. The septum can probably be moved outward to, for example, $x_s = 70$ mm, thus allowing displacements to 100 mm at the ES without exceeding the horizontal aperture in the bending sections of the machine. However, one would have to be careful about the unstable beam striking other hardware, for example, the injection septum, on the next-to-last turn.

The ES should be 4.0 m in length and be capable of 5.0 MV/m across a 30-mm gap. The Lambertson should be 10.0 m in length and attain a peak field of 1.0 T. As before, the perturbing fields should be obtained from auxiliary

Fig. 4.3.4-3. Trajectories of three particles in the LSS between the ES and the slow-extraction Lambertson. The two trajectories starting above $x_s = 60$ mm are kicked out.

Fig. 4.3.4-4. These trajectories show the maximum horizontal sizes on the last turn through the accelerator. They determine the aperture.

multipole windings in the sextupoles. At 45 GeV the required $B'l$ is 0.8 T for the quadrupoles, and $B'''l$ is 460 T/m² for the octupoles.

4.3.4.2 Fast abort. The fast-abort components are indicated at the top of Fig. 4.3.4-1. Two fast kickers are energized simultaneously to impart a coherent horizontal betatron oscillation; the displaced beam is then brought out by the fast-abort septum. The centers of these elements are located at $s = 1253.98$ m, 1265.62 m, and 1304.26 m relative to the center of the ES, respectively. These components must be programmed to track the accelerator magnets. The apertures should be able to pass the entire beam including tails. This means a kicked beam with $\epsilon = 8.0\pi$ mm-mrad should be passed. Figure 4.3.4-5 shows the corresponding beam envelopes for both an unkicked and kicked beam. We impart positive kicks of 0.6 and 1.0 mrad, respectively, using the two kickers. At 45 GeV, K_1 ($l = 3$ m) needs a peak field of 0.031 T; the second kicker K_2 should be 5 m in length and reach a similar field. The fast-abort septum needs to be 10 m long and reach 1.0 T.

4.3.5 Stability Analysis

The analysis for the main ring closely parallels that for the booster. None of the formulas or qualitative discussion will be repeated. Only tables of results and interpretation will be given. Section 4.3.5.1 deals with single-particle instabilities and Sec. 4.3.5.2 discusses coherent instabilities.

4.3.5.1 Single-particle instabilities. The tunes for the main ring are $\nu_x = 7.45$ and $\nu_y = 6.45$ in order to be near the half-integer resonance for resonant slow extraction. The racetrack configuration has superperiodicity of two. All the sum and difference resonances of order less than fifth are shown in Fig. 4.3.1-1. The transverse space-charge tune-shift region is schematically indicated in Fig. 4.3.1-1. This region does not cross any superperiodicity resonances of order less than fifth. The horizontal tune was chosen so that the working point lies below the diagonal difference resonance. The lattice design is flexible enough to change to another working point should it be necessary.

Fig. 4.3.4-5. Beam envelopes in the full-aperture fast-abort system.

TABLE 4.3.5-II

IMPEDANCE LIMITS FOR THE LAMPF II MAIN RING

Parameters ^a	At Injection	At Extraction
E	6.938 GeV	45.938 GeV
$\Delta p/p$ (rms)	0.09%	0.04%
σ_{ρ} (rms)	0.35 m	0.26 m
$ \eta $	0.0095	0.0274
I_p	15.9 A	21.9 A
n	1897	2564
$\Delta v_{\perp}$	0.030	0.033
$ Z_{\parallel}/n <$	14.5 Ω	79.0 Ω
$ Z_{\perp} <$	10.02 M Ω /m	55.06 M Ω /m
$b^2 Z_{\perp} /(2R) <$	14.87 Ω	81.74 Ω

^aOther parameters: $\gamma_{tr} = 6.0$, $N_b = 2.95 \times 10^{11}$, $v_x = 7.45$, $\xi = 1.0$, $\Delta v_{oct} = 0.02$, circumference = 1323.53 m, $\bar{\beta}_x = 12.5$ m, $\bar{b} = 0.025$ m.

TABLE 4.3.5-III

LOW-FREQUENCY IMPEDANCES FOR THE MAIN RING^a

	At Injection	At Extraction
T (GeV)	6.0	45.0
$ Z_{\parallel}/n _{sc} \Omega$	3.31	5.64
$ Z_{\parallel}/n _{rw} \Omega$	1.81	1.82
$ Z_{\parallel}/n _{tot} \Omega$	4.06	5.56
$ Z_{\parallel}/n _{lim} \Omega$	14.5	79.0
$ Z_{\perp} _{sc} \text{ M}\Omega/\text{m}$	11.5	2.05
$ Z_{\perp} _{rw} \text{ M}\Omega/\text{m}$	6.88	6.88
$ Z_{\perp} _{tot} \text{ M}\Omega/\text{m}$	14.5	7.57
$ Z_{\perp} _{lim} \text{ M}\Omega/\text{m}$	9.087	36.16

^a $n = \omega/\omega_0 = 8$, $f_0 = \omega_0/2\pi \sim 0.023$ MHz.

TABLE 4.3.5-IV

HIGH-FREQUENCY IMPEDANCES FOR THE MAIN RING

	<u>At Injection</u>	<u>At Extraction</u>
n	1897	2564
T GeV	6.0	45.0
$ Z_{ }/n _{sc} \Omega$	3.31	5.64
$ Z_{ }/n _{rw} \Omega$	0.035	0.030
$ Z_{ }/n _{tot} \Omega$	3.34	5.62
$ Z_{ }/n _{lim} \Omega$	14.5	79.0
$ Z_{\perp} _{sc} M\Omega/m$	11.5	2.05
$ Z_{\perp} _{rw} M\Omega/m$	0.03	0.02
$ Z_{\perp} _{tot} M\Omega/m$	11.5	2.07
$ Z_{\perp} _{lim} M\Omega/m$	10.0	55.1

(microwave) instability is of concern. As in the case of the booster, the transverse space-charge impedance can be reduced by increasing the emittance of the beam. For example, if the beam radius is increased by a factor of 1.5, the space-charge transverse impedance is reduced by a factor of 2.5, which is more than adequate for stability.

With regard to the coupled-bunch instabilities, comments made for the booster apply to the main ring. The same active damper that will be needed for fast correction of the injection errors will damp transverse coupled-bunch instabilities. Landau damping of the longitudinal coupled-bunch modes will occur if there is no resonator in the system with a shunt impedance less than 100 k Ω and if the synchrotron frequency spread can be made larger than the growth time of the instabilities.

With regard to the Robinson instability, there are two problems. The first is the transients created during injection, when for 50 ms the ring is partially empty. The second is the ratio of the power delivered to the beam divided by the power dissipated by the rf cavity during the acceleration phase.

Calculations show that the voltage droop in the rf cavities in the empty times during injection is more than 40%. This means that when a train of bunches comes around to a cavity, there will be a substantial change in the gap voltage as the first bunch enters the cavity. If it is not possible to detune the cavity fast enough to ensure that the first few bunches see the proper gap voltage, some bunch lengthening may occur in these bunches. The rf-cavity voltage and phase must be adjusted on a submicrosecond time scale. A better solution is to preprogram (feedforward) the cavities before the transient arrives. In fact there may be no problem at all. If one can leave the cavities in the proper detuned state during the empty part of the turn, then when the batch returns to the cavity, the beam loading brings the cavity back to the proper voltage automatically. Unfortunately, if the empty cavity is not on resonance this will cause some power loss. There may be other reasons not to leave the cavity detuned.

We have not carried out detailed calculations of the beam-loading problem for the main ring, but it is clear that the problem will be no worse for the main ring than it is for the booster, since the number of particles per bunch is the same.

4.3.6 Main-Ring Magnets and Power Supplies

4.3.6.1 General. The main ring bending is accomplished with combined-function magnets operating at 3 Hz. The field gradient and lower frequency make some detail design changes necessary in components, but in principle the magnets and power supplies follow the concepts set out for the booster (see Sec. 4.2.6).

4.3.6.2 Dipole magnets. There are four designs of dipoles as shown in Table 4.3.6-I. It is clearly worthwhile to try to consolidate the number of magnet types while retaining a satisfactory lattice, because the differences in aperture and gradient are quite small. As shown in Fig. 4.3.3-1, the main ring ramps up at 5.7 Hz, has a 167-ms flattop, and ramps down at 17 Hz. Using 64F360 steel (about 1% silicon), the core loss will be about 0.5 w/kg with this 50% duty factor. Because of the gradient, the magnets must be built to follow the beam trajectory accurately; the radius is 117.85 m. The HF, HF2, and HD dipoles will weigh about 20 tons each; the HD2's, 10 tons.

A POISSON-generated half-section of an HF2 dipole is shown in Fig. 4.3.6-1. The poles are smooth hyperbolae, with shims. The suggested coil operating

TABLE 4.3.6-I

DIPOLES FOR LAMPF II MAIN RING

# Required	Type	Length (m)	Field on Axis (T)	Gap (cm)	Gradient T/m
32	HF	9.255	1.3	10.2 W × 5.84 H	0.944
8	HF2	9.255	1.3	10.5 W × 5.36 H	1.35
32	HD	9.255	1.3	8.02 W × 6.64 H	-0.888
16	HD2	4.628	1.3	6.78 W × 6.70 H	-1.278
88

Fig. 4.3.6-1. HF2 main-ring dipole (top half).

TABLE 4.3.6-II

COIL OPERATING CONDITIONS

Length of one turn	20 m
Peak/RMS current	1400/1000 A
Resistance at 20 C	4.4 m Ω
Resistive voltage/turn	6.44 V
Outside dimensions	10 mm ²
Round hole i.d.	5 mm
I ² R power	5.26 kW
Max. eddy-current power	0.8 kW
Temperature rise	23°C
Cooling water flow rate	3 l/min
Cooling water velocity	3 m/s
Pressure differential	5.5 atm
Temp. diff. between copper and water	1.05°C

conditions are given in Table 4.3.6-II. The eddy-current heating is much less than in the booster, without any special coil construction. One possible arrangement is four conductors in parallel, 5.6 kA maximum current, and 2.87 kV maximum per magnet.

4.3.6.3 Quadrupoles. The quadrupole requirements are in Table 4.3.6-III. These could be consolidated to two designs, one 20-cm bore $\times$ 5.1 T/m, the other 10 cm bore $\times$ 13.4-T/m. All the quadrupoles are 1 m long. Neither design is particularly demanding, and construction similar to that shown in Fig. 4.3.6-3 is envisaged. The design will follow that used at SPEAR and PEP at SLAC (Fig. 4.3.6-3) and multipole harmonics will be "tuned out" using techniques established with Los Alamos' PSR quadrupoles. Typical results are given in Table 4.3.6-IV.

4.3.6.4 Kickers. The main-ring full-aperture fast-abort system will be similar to that of the booster. (Initially there will be no fast-extraction system.) A sufficient number of kicker modules will be used to provide the required kick angle. The kick must track the main-ring field, and the availability of a large number of modules makes possible adjusting the kick by choosing the correct number of modules to be fired. Only a dc power supply for the thyratrons will be required.

TABLE 4.3.6-III

QUADRUPOLES FOR LAMPF II MAIN RING

Type	# Required	Bore diam (cm)	Gradient T/m
QF1	8	19.66	4.91
QD1	8	13.08	-5.1
QBF	8	9.14	13.4
QBD	4	8.8	-11.3

Fig. 4.3.6-3. Quadrupoles for LAMPF II main ring.

TABLE 4.3.6-IV

RANGE OF QUADRUPOLE HARMONICS

Current	420A (%)	490A (%)
n = 3	0.004 - 0.048	0.007 - 0.041
n = 4	0.012 - 0.043	0.018 - 0.043
n = 5	0.026 - 0.077	0.008 - 0.085
n = 6	0.059 - 0.142	0.058 - 0.113

The injection kicker for the main ring is designed in a manner similar to the fast-abort kicker by assembling the appropriate number of modules to provide the required deflection. The injection-kicker pulse must turn off quickly, unlike the abort-kicker pulse. The injection kicker does not track the ring magnets and has a fixed-voltage power supply. The pulse-forming network is a straightforward design and exploits the mutual inductive coupling between the adjacent inductors to reduce the ripple in the pulse supplied to the magnet module. Similar networks have been constructed and no problems are anticipated.

4.3.6.5 Power supplies. The concept for the main-ring power supplies follows that for the booster. A shaped waveform is generated, with a flatbottom for injection, a rise at 5.7 Hz for acceleration, an 167-ms flattop for extraction and a ramp-down at 17 Hz. Using fast extraction, the flattop would be eliminated and the repetition rate would double to 6 Hz.

To keep the peak voltage of the ring at a reasonable level, the magnets would have 12 series supplies. Four supplies feed the matching sections, each with six dipoles (4 of which are the half-length HD2's), whereas the remaining 64 magnets are in 8 groups of 8. This results in a power-supply rating of 5.6 kA, ± 12 kV for the conventional-magnet scenario.

It has recently become apparent that there could be considerable cost savings from using a rotating machine as an inertial energy-storage device. Two possibilities are being studied: the surplus ZGS motor-generator set (two 81.8-MVA generators driven by a 15 500-hp motor) and unused nuclear-power-station generators. We presently favor the ZGS motor-generator system. This system would deliver a waveform with linear ramps rather than sinusoids. Because of time constraints on writing this proposal the beam-dynamics work has been done using a sinusoidal waveform. The cost estimate uses the Argonne ZGS motor-generator set. Further work during the conceptual design stage will clarify this issue.

4.3.6.6 Superconducting main ring. Consideration is being given to making the magnets of the main ring with superconducting coils. Following suggestions from R. R. Wilson, the coils would use cabled superconductor in a nonmetallic duct containing liquid helium. At the field level of these magnets, an iron core would be retained (superferric) to gain the magnetization of the iron and to control the gap field shape. The coil cooling will cryopump the magnet vacuum chamber. Any new advances in ac superconducting cable technology would make this approach all the more attractive.

4.3.7 rf System

There are many possibilities for the main-ring cavities. These include using the booster cavities with fewer ferrite rings installed and the Fermilab main-ring cavities. In either case, a version of the booster amplifier with two rf stages rather than three is appropriate because of the smaller tuning range required. Our main-ring R&D program will address these issues after the more difficult booster system has been successfully prototyped.

4.3.8 Vacuum System

The main-ring vacuum system will have a low impedance as is the case in the booster. Although the repetition rate is lower, the field is higher and the length of each section is longer. The conductivity of the surface must be high-

er, since the stability criteria are harder to satisfy. Our conclusion is that a metal pipe cannot be used and that a plated ceramic pipe is required as in the booster.

4.4 AUXILIARY SYSTEMS

4.4.1 Controls

This section addresses the scope of the controls network and comments on special requirements. Although the LAMPF II accelerators are unique in combination, the separate elements present few unique or new controls problems. The LAMPF linac already runs a high power beam with high pulse repetition rate and complicated pulse-to-pulse beam distribution. The booster synchrotron timing requirements are generally parallel to the LAMPF Proton Storage Ring. The circulating beam currents are also comparable to the PSR, and much of the instrumentation would be similar. We therefore conclude that the experience necessary to implement the LAMPF II control system is readily available. The technical details, especially in the computer system, would be defined at a later date using the state of the art. We expect that the falling cost of computing units enable a safe scaling of current project data system costs to the demands of LAMPF II.

4.4.1.1 Function of the system.

The primary function of the control system is to provide operator interface to the equipment. The modern control system excels in human-engineered features. Selective display of data and concentration of controls improve operator efficiency. Automation of routine tasks such as monitoring improve machine reliability.

The beam-ready system can be implemented with software as in the PSR or with commercially-available programmable hardware. Either approach can reduce their labor required to keep the beam-ready system accurately reflecting the beamline hardware. The timing-gate system is another autonomous function that benefits from computer programmability and the emergence of commercial equipment.

We expect to build a system with adequate redundancy, for example, with a backup of the central control computer (if that architecture were chosen). The second computer would usually be available for machine- and control-system development.

4.4.1.2 Issues in choosing a computer system. The main issues in designing a computerized control system are the size and complexity of the demands on the system, communication network configurations, compatibility between the main large computers and the computers controlling pieces of equipment, equipment interfaces, software configurations, and cost.

There is a considerable advantage in using computers of the same architecture from the microprocessor-chip level to the large, powerful machines and in requiring that this architecture provide a large (32-bit) address space. Both these considerations will result in considerable savings in the cost of writing and testing software. Another consideration in the choice of a computer family is the availability of useful control and debugging software and staff familiar with its use. These considerations can easily overrule any cost advantage of a particular vendor.

Most control systems place computers at two or more levels in the system. At the highest level, the main function of the computers is operator interaction and support. At the next level are the equipment computers that provide the monitoring functions of the system and in which are coded the specific procedures for each piece of equipment. The number and size of computers at this level is determined by the number as well as the amount of activity each generates. On general principles it is better to build as much control into the lower level computers as possible, thus relieving the higher level computers from communications tasks.

The choice of communications network is between a message-based networking system and a bus-to-bus connection, such as shared memory. The former is receiving much attention and development with products like Ethernet, but it uses more overhead. The bus-to-bus connection, as in Fastbus, might be provided by connecting the various computers in some way to a modern 32-bit backplane. The communications medium standard is presently coaxial cable. Optical fiber, providing wider bandwidth, is being commercially developed and may soon be an attractive alternative.

The final choice to be made is how to connect particular pieces of equipment to the computers of the system. Currently, CAMAC is the clear choice in the field and is likely to remain so for some years. However, another possibility being explored at CERN for LEP is the use of a small microprocessor system for each piece of equipment with an ASCII network connection to the equipment computers. The equipment engineers, not the controls staff, then program these microprocessors.

4.4.1.3 Number of channels. The size of the control system depends on the number of channels. A number of recent projects, including the SLAC SLC, CERN PS controls upgrade, and the Fermilab Tevatron used roughly 30 000 channels. We assume LAMPF II will require 10 000 channels each for the booster, main ring, and experimental areas. An additional 10% of the channels will be high-cost channels. The low-data-rate channels are suitable for turning equipment on and off and monitoring relatively-slowly changing voltages, currents, etc. The high-data-rate channels are needed for some diagnostics like beam-position monitors and for timing and synchronization pulses.

4.4.1.4 An illustrative design using today's technology. For the purpose of costing a system, the present PSR control-system design will be used with some extensions to meet the LAMPF II control requirements. This system is described in (C1 83). Figure 4.4.1-1 shows a possible computer control system. The console is driven by one of the two upper-level machines. The second machine acts as a development machine, but can be connected quickly as the spare if necessary. Both computers share a common, redundant disk and tape system. As we said above, the intercomputer network standard by LAMPF II remains to be defined. The remote equipment interface standard is assumed to be CAMAC, at least for purposes of costing.

4.4.2 Beam instrumentation

This section briefly reviews the various types of beam instrumentation necessary for LAMPF II. As we found with controls, few new techniques appear to be needed. Most of the instrumentation has been developed for the LAMPF linac and storage ring. We will, however, need to pay special attention to the impedance of the monitors as seen by the beam.

Fig. 4.4.1-1. A possible computer architecture for the LAMPF II system.

Beam properties to be measured are the particle distributions in transverse and longitudinal phase space, synchronous properties in the rings, and beam current and beam losses. Subsystems have various requirements on speed, accuracy, detail, bandwidth, etc. Specialized subsystems may participate in active control loops.

4.4.2.1 Transverse measurements. Typical profile instruments used at LAMPF are moving wire scanners, fixed multiwire detectors, and moving slits. These devices can produce accurate measurements for beam optics. The dynamic range extends up to a few milliamps average beam current (not usable in ring machines). For measurements within PSR, a non-intercepting quadrupole moment detector is being developed. This uses the four-quadrant beam position monitor (BPM) with different signal processing. The BPM system would be extensively used in the rings since it is non-intercepting, fast, and has a suitable dynamic range. With the addition of a pulsed deflector, the BPM's allow measurement of ring betatron tune. The BPM system can also provide longitudinal and intensity measurements.

4.4.2.2 Longitudinal measurements. A beam-driven rf pickup system with adequate bandwidth to indicate microbunch length as well as mean phase is necessary and is being developed in Line D and the PSR. The technology is state-of-the-art but not unprecedented. A system for closed-loop energy and phase control will be used in the PSR and is close to what is required for the booster.

4.4.2.3 Current monitors. At LAMPF, an accurate (0.1%) beam transformer system is used which would also be usable in the LAMPF II transport lines. In the rings, wall-current monitors will be used. The rf pickup systems can also supply beam current information, albeit with less precise gain.

4.4.2.4 Beam loss protection. Current monitors are useful in beam loss and beam protection systems. Radiation detectors of various kinds are also an essential and independent method of protection. Fast shutoff and fast abort systems have beam and radiation detector inputs.

4.4.3 Active Dampers

4.4.3.1 General comments. Unless the impedances $Z_{\parallel}$ and $Z_{\perp}$ can be kept small enough (or the momentum and tune spread large enough), it may be necessary to damp unstable oscillations with wavelengths longer than or equal to the bunch length. The most important of these unstable modes are the so-called longitudinal and transverse resistive wall instabilities (Ne 65), (La 65), (Co 66), or (Sa 72). An active damper system may also be needed to reduce the amplitude of the injected beam in case injection is not perfectly on the design orbit.

If the number of bunches n_b in the beam were small (<5), an ordinary damper would probably be adequate, but since n_b is 60 for the booster and 240 for the main ring, it will probably be necessary to consider a so-called superdamper of the type developed at FNAL. An ordinary transverse damper measures the average amplitude of the beam at a location where the betatron function is large (near a focusing quadrupole) and applies a transverse, sinusoidal voltage kick, 90° in betatron phase downstream, of a size proportional to the beam-displacement amplitude. The phase of the sinusoidal kick is $(n\omega_0 t + \phi)$, where $\omega_0 = \beta c/R$ is the circulation frequency, ϕ is a predetermined delay that can be a function of time, and n is usually 1 or 2 times the betatron tune. The system requires only a relatively narrow band amplifier and simple stripline detectors and deflectors.

The superdamper damps instabilities on a bunch-by-bunch basis. The displacement of an individual bunch from the axis is detected and this signal is processed through an energy-dependent, variable-delay line and a wideband, distributed amplifier to a deflector, which applies a restoring kick almost a full revolution later (An 77).

There are many accelerators that have installed damper systems. These include the MURA 50-MeV electron accelerator, Argonne's ZGS, Rutherford's 8-GeV proton synchrotron NIMROD, CERN's ISR and SPS, FNAL's booster and main ring, SPEAR, and PETRA. This list is not exhaustive. There is very little doubt that a transverse damper can be developed for the LAMPF II rings.

Longitudinal dampers are less common. The horizontal transverse damper can influence the longitudinal dipole modes through betatron-synchrotron coupling. Also, damping has been applied to quadrupole modes. Introducing frequency modulation into the rf cavities at the fundamental circulation frequency ω_0 helps, too. Another effective measure is to run one of the rf cavities at a slightly different harmonic number. It is also helpful to work with full buckets, which use the nonlinear change in the synchrotron frequency near the separatrix to provide Landau damping. This probably is not an option for LAMPF II because we want to use the rf system as a buncher near the end of the acceleration cycle. The growth rate of the longitudinal resistive-wall instability is

directly proportional to the energy and would be the largest at the end of the cycle when the beam is highly bunched. Based on the stability calculations in Secs. 4.2.5 and 4.3.5, the longitudinal resistive-wall instability should not be a problem, and we are making no plans for an active longitudinal damper. The next section describes the transverse superdamper in more detail.

4.4.3.2 Requirements for a transverse damper. The damping time τ_d must exceed the anticipated growth time τ_g of the instability. For example, a worst-case estimate of the growth time (ignoring damping due to tune spread) for the lowest mode transverse resistive-wall instability is given by (Co 66),

$$\tau_g = \frac{2\pi^2 b^3 v}{(2\pi R) r_0 c N} \left[(n-v) \frac{2Z_0 \sigma (2\pi R)}{\pi} \right]^{1/2} \frac{\gamma}{\sqrt{\beta}}, \quad (4.4.3-1)$$

where the definitions and values of the constants and the τ_g just after injection are given in Table 4.4.3-I. The growth time is worst just after injection. Because γ , β , and N change with time into the cycle, the growth of the instability is more nearly linear than exponential with time. A calculation of the growth factor

$$F(t) = \exp \left[\int_0^t \frac{dt'}{\tau_g(t')} \right] \quad (4.4.3-2)$$

showed that a coherent disturbance at injection would grow only by a factor 4.6 by the end of extraction in the booster. The growth factor for the main ring is

TABLE 4.4.3-I
EVALUATION OF THE GROWTH TIME τ_g

	Symbol	Booster	Main Ring
Avg. half-height of chamber	b	0.0311 m	0.0237 m
Machine circumference	$2\pi R$	330.88 m	1323.32 m
Proton classical radius	r_0	1.534×10^{-18} m	
Velocity of light	c	2.9979×10^8 m/s	
Vertical betatron tune	v	6.30	6.45
Next largest integer	n	7	7
Number of protons in ring	N	1.77×10^{13}	7.07×10^{13}
Impedance of free space	Z_0	376.61 Ω	376.61 Ω
DC wall conductivity	σ	6.25×10^7	6.25×10^7
Energy/(mc ²) at max τ_g	γ	1.735	6.938
Velocity/c at max τ_g	β	0.8414	0.9908
Max. growth at rate	τ_g	7.5 ms	0.3 ms

much larger. The growth times for the two rings in Table 4.4.3-I differ by a factor of 24.3. Furthermore, the cycle times are a factor of 20 different, giving more time for growth in the main ring. It would appear from this calculation that the booster might not need an active damper if Landau damping due to space-charge effects and chromaticity is large enough. The main ring will almost certainly need an active superdamper.

Critical quantities in the design of a damper are the system gain constant, the deflector power, and the bandwidth. The system gain constant V' is the deflection voltage on the stripline deflector plates per unit of beam displacement. It can be shown (Hi 80) that the system damping constant is approximately

$$V' = \frac{2\beta^2 E \tau_r d}{e l B \tau_d \sin(2\pi v)} \quad , \quad (4.4.3-3)$$

where the various quantities are defined in Table 4.4.3-II below. If for safety we require that $\tau_d \sim \tau_g/5$, then we can evaluate V' as shown in Table 4.4.3-II. Using these system gain constants and assuming that the amplitude growth in any revolution will be <0.5 cm, then the position detector and deflection amplifiers together must produce a voltage of at least 47.0 V for the booster and 6.6 kV for the main ring. For a 50- Ω system this corresponds to a power of 45 W for the booster and 881 kW for the main ring. The distributed amplifier for the booster is easy to design and fabricate. The amplifier for the main ring is very difficult. Before the final design of the main-ring damper, one must find a more realistic way of estimating the growth time τ_g that includes the effect of Landau damping. It is very likely that this will reduce the required power by a factor of 40. Even then it would be more economical to build a ~ 15 -kW ordinary damper system and supplement it with a ~ 6 -kW superdamper.

TABLE 4.4.3-II

EVALUATION OF THE SYSTEM GAIN CONSTANT V'

	Symbol	Booster	Main Ring
Machine circulation time	τ_r	1.31 μ s	4.43 μ s
Velocity/c at injection	β	0.841	0.9908
Total energy at injection	E	1.780 GeV	6.938 GeV
Electrode plate separation	d	0.06 m	0.05 m
Electrode plate length ^a	l	1.0 m	1.0 m
Betatron amplitude at detector	B	22 m	60 m
Vertical tune	v	6.31	6.31
Required damping time	τ_d	1.5 ms	0.06 ms
System gain constant	V'	9.4 kV/m	1.33 MV/m

^aElectrode plate lengths are an important parameter in the system. Detector lengths may differ markedly from the deflector length.

The bandwidth of the amplifier for the bunch-by-bunch damping must span the range between the pulse repetition rate at the detector and the nominal rise time of the individual pulses. The pulse repetition rate for the booster is 50.3 MHz and for the main ring, 59.7 MHz. Each pulse is approximately 5 ns wide, and hence the rise time is approximately 2 ns. The inverse of this rise time gives an upper limit to the needed bandwidth of approximately 500 MHz. This is a moderately wide bandwidth, but well within the state of the art. At this stage of analysis it appears that a transverse damper can be constructed for both the booster and the main ring without prohibitive cost.

4.5 SHIELDING AND ACTIVATION

4.5.1 Magnitude and Location of Losses

There are a number of possible spill sources in the LAMPF II accelerators. Our accelerator calculations have not progressed to the point where it is possible to give precise estimates of the expected losses. In most cases, we have given upper limits that will be controlled by use of spill monitors, as is the experience at LAMPF. A list of the expected sources is given in Table 4.5.1-I.

In existing machines approximately 30% of the beam is lost in the early stages of acceleration. Typically this occurs when the space-charge tune shift and synchrotron frequency are at a maximum and the bucket area is simultaneously near a minimum. In the LAMPF II machines, the large apertures and bucket areas should serve to prevent a large part of these losses. There is no reliable way to estimate the spills at this time. We have assumed that a factor of 10 improvement will be obtained; hence our estimate of 3% spill. We expect that all of this loss will be contained by the collimator system.

TABLE 4.5.1-I
EXPECTED SPILL SOURCES IN LAMPF II

<u>Location</u>	<u>Fraction (%)</u>	<u>Energy (GeV)</u>	<u>Protons/s</u>	<u>Type</u>
<u>Booster</u>				
Injection cleanup	1	0.8	1×10^{13}	continuous
H ⁰ dump	1	0.8	1×10^{13}	continuous
Collimators	3	1.2	3×10^{13}	continuous
Uncontrolled spill	1	6.0	1×10^{13}	intermittent
Abort dump	1	1.6	2×10^{13}	intermittent
Fast extraction	1	6.0	1×10^{13}	continuous
<u>Main Ring</u>				
Injection cleanup	1	6.0	2×10^{12}	continuous
Collimators	3	9.0	6×10^{12}	continuous
Uncontrolled spill	2	45.0	4×10^{12}	intermittent
Abort dump	3	45.0	6×10^{12}	intermittent
Slow extraction	1	45.0	2×10^{12}	continuous

The uncontrolled spills and aborts will be limited by spill monitors. Hardware interlocks will automatically reduce the injected-beam repetition rate to the point where larger spills are impossible. Experience at LAMPF indicates that tuning to prevent losses is feasible with this logic.

The slow-extraction system of the main ring will intercept less than 1% of the beam. Most of this spill will be localized to the magnetic extraction septum. We have conservatively estimated that the booster extraction system will spill the same amount of beam. This allows for the future addition of a slow-extraction system for the booster, if desired.

From an examination of the list of losses, it is clear that all but the uncontrolled spills will be at well-defined locations. Thus local shielding of these points is possible. The uncontrolled spills will be kept to a level that does not activate the machine.

4.5.2 Shielding Requirements

The LAMPF II accelerators will be located underground in tunnels. The tunnel elevation is chosen to put the machines far enough underground that buildings on the surface can be occupied by normal employees working full time.

The required shielding is based on a continuous loss of 10^{13} protons/s at any one spot. This will be more than sufficient for the uncontrolled spills, which will most likely be distributed in many places around the ring. There is no problem with a single pulse, as one complete pulse of the booster could be spilled every two seconds without exceeding this limit. Although 10^{13} in a single spill is a large loss of beam, existing machines such as the Brookhaven AGS routinely tolerate accidents of this magnitude without short-term damage to components. Thus the most important problem is protection of people who may be working on the surface. Additional local shielding will be provided around dumps, collimators, and septa to keep the surface dose to the required level.

There is considerable experience with the shielding necessary for 1- to 30-GeV protons. The work at CERN has been particularly well documented (Go 71). The initial shielding requirements for LAMPF II have been based on the CERN data. The requirements for transverse shielding are presented in Table 4.5.2-I.

Longitudinal shielding requirements are more difficult to determine, since the longitudinal dose after thick shielding is dominated by muons. Untested extrapolation from existing data is required to set the longitudinal thickness. Fortunately, the LAMPF II accelerator site layout naturally provides a far

TABLE 4.5.2-I

TRANSVERSE SHIELDING REQUIREMENTS FOR LAMPF II

	Thickness (m)	Thickness (feet)
Steel	3.8	12.5
Compacted tuff	13.0	43.0
Virgin tuff	16.3	53.0

thicker longitudinal shield than transverse shield. Further calculations are needed to define the longitudinal shield more precisely. Additional transverse shielding in the form of earth or concrete can be provided to correct any weak spots at moderate cost.

The site plan for LAMPF II is based on 53' transverse and 250' longitudinal shielding in virgin tuff. The beam-line elevation of 6900 feet of the LAMPF II accelerators was chosen to meet the transverse requirement. The longitudinal shielding requirement may need revision later.

4.6 MODIFICATIONS TO LAMPF I

The present LAMPF linear accelerator will make a very good injector for LAMPF II. Production beam current of 1.0 μA average at 9% duty factor are routine. The LAMPF II booster design requires less than one-eighth of the LAMPF beam, because it is running at 60 Hz as compared with 120 Hz for LAMPF and at an rf frequency of 50.3125 MHz as compared with 201.25 MHz for LAMPF. Furthermore, not all the rf buckets can be filled, because there must be a notch to allow for the rise of the extraction kicker magnet. One-eighth of the present 1.0 mA beam is 125 μA of H^+ . LAMPF II requires 170 μA of H^- . The modifications required are thus quite modest and consist of increasing the duty factor and developing an H^- source comparable in brightness to the present H^+ source.

4.6.1 Duty-Factor Improvements

The LAMPF accelerator was designed for 12% macroscopic duty factor. Present performance of the machine has included long periods of running at 10.5% duty factor and short test periods of all components at 12% duty factor. An increase from 9 to 12% will raise the effective current available to LAMPF II from 125 to 167 μA , which almost meets the 170- μA requirement.

4.6.2 H^- Source

LAMPF is now being modified to produce high-current H^- beams for the PSR. These modifications will be complete by mid-1985, and at that time the accelerator will deliver 100 μA of H^- at 9% duty factor, which is to $\sim 3.45 \times 10^8$ $\text{H}^-/\text{microbunch}$. The booster design goal can be met by a small increase in peak current. A more attractive possibility is to develop a factor of two increase in micropulse charge at 50 MHz, which could eliminate the requirement for a flatbottom on the injection cycle.

4.7 POLARIZED BEAMS

LAMPF II will readily accelerate polarized beams. In the following sections we show calculations indicating that with careful alignment and pulsed quadrupoles, it should be possible to obtain polarized beam at any energy up to 45 GeV. The polarized beam current depends on future improvements to the polarized H^- ion source which are beyond the scope of this document.

4.7.1 Mechanisms for Depolarization

In a proton synchrotron the polarized beam is injected and then accelerated with its polarization parallel to the vertical guide field. Off-axis particles that encounter horizontal field components experience a spin precession of $2\pi G\gamma$ each turn. The quantity $\kappa = G\gamma$ is called the spin tune; G is the anomalous magnetic moment (≈ 1.7928 for protons); and γ is the usual relativistic gamma. If κ is close to ν_0 , where

$$\nu_0 = n_x \nu_x + n_y \nu_y + n_s \nu_s + n \quad , \quad (4.7.1-1)$$

then a small perturbing magnetic field will lead to a substantial deviation of the spin direction and hence depolarization. In Eq. (4.7.1-1) ν_x , ν_y , and ν_s are the horizontal, vertical, and synchrotron tunes, respectively; n_x , n_y , n_s , and n are integers.

The degree of depolarization depends upon the difference between κ and ν_0 . The depolarization resonance width $\epsilon(\kappa)$ is the width of the interval around ν_0 such that a beam of particles with κ at the edge of that interval will have its polarization reduced by $>50\%$. Normally one does not sit at a depolarization resonance, but rather passes through resonances during the acceleration phase. The Froissart-Stora equation (Fr 60) gives the amount of depolarization that occurs for an adiabatic passage of an accelerated beam through a resonance;

$$\frac{P_f}{P_i} = 2 \exp(-\pi^2 |\epsilon|^2 / 2\alpha) - 1 \quad . \quad (4.7.1-2)$$

In Eq. (4.7.1-2) P_f (P_i) is the polarization long after (before) passage through the resonance. The quantity α is the crossing speed and is given by

$$\alpha = \frac{1}{2\pi} (G\Delta\gamma - \Delta\nu_0) \quad , \quad (4.7.1-3)$$

where $\Delta\gamma$ and $\Delta\nu_0$ are the changes in γ and ν_0 per turn during acceleration. Evidently there are two ways to avoid depolarization: either one quickly

crosses a weak resonance $|\epsilon|^2/\alpha \ll 1$ or slowly crosses a strong resonance $|\epsilon|^2/\alpha \gg 1$, in which case the spin will be totally flipped. The depolarization resonance width ϵ is given by (Co 80)

$$\epsilon(\kappa) = \frac{1}{2\pi} \int_0^{2\pi} \xi(\theta) e^{-i\kappa\theta} d\theta, \quad (4.7.1-4)$$

where $\xi(\theta)$ is a function of the spin tune, magnet lattice, and the vertical excursion of the beam. The form of Eq. (4.7.1-4) is a Fourier amplitude, and $\epsilon(\kappa)$ will be significant if $\xi(\theta)$ contains frequencies that coincide with κ .

The strongest depolarizing resonances are of two classes: (1) imperfection resonances that occur for $\kappa = n$ and arise because of errors in magnet alignment or field; (2) intrinsic resonances that occur when $\kappa = mS \pm \nu_y$, where S = superperiodicity of the machine. In calculating ϵ one uses for the vertical excursion (1) the vertical closed-orbit distortion or (2) the size arising from betatron oscillation, respectively, for these two cases. The letters n and m represent positive integers.

4.7.2 Depolarization in the Booster

A program has been written to study depolarization during acceleration; Eq. (4.7.1-2) was used to evaluate the depolarization at each resonance crossing. Equation (4.7.1-3) was evaluated using the expression

$$\Delta\gamma = \frac{\pi R \omega_a}{c \times m_p c^2} (p_e c - p_i c) \sin \omega_a t, \quad (4.7.2-2)$$

where $2\pi R$ is the circumference; ω_a is π/τ_a ; $\tau_a = 11.75$ ms is the time for the acceleration phase; p_i (p_e) is the momentum at injection (extraction); and $m_p c^2$ is the proton rest energy.

Two intrinsic resonances are encountered during acceleration from 0.797 to 6 GeV for the booster with 8 superperiods ($S = 8$). The vertical excursion $y(\theta)$ is $(\beta_y(\theta)\epsilon_y)^{1/2}$, where β_y is the vertical beta function and ϵ_y is the geometric emittance. We assume that $\epsilon_y = 1.0\pi$ mm-mrad at 0.797 GeV and that it damps inversely with momentum. Table 4.7.2-I lists the resonances, momentum, α , $|\epsilon|$, and P_f/P_i for the four resonances. The $\kappa = \nu_y$ resonance results in a calculated depolarization of 40%. Fast-pulsed quadrupoles can be used to introduce a tune shift $\Delta\nu_y$. For $\Delta\nu_y = 0.05$ we can obtain $P_f/P_i = 0.988$ for ($\kappa = \nu_y$).

A similar analysis was performed for the 10 imperfection resonances $4 \leq N \leq 13$; this time the vertical excursion was taken to be the vertical closed orbit. The closed orbit was generated in a Monte Carlo simulation that assumed vertical magnet misalignments of up to $\pm 2\sigma_y$ with $\sigma_y = 0.1$ mm. Figure 4.7.2-1 is a plot of P_f/P_i obtained at each value of n ; some depolarization is predicted

TABLE 4.7.2-I

DATA FOR THE BOOSTER INTRINSIC RESONANCES

Resonance	p GeV/c	α	$ \epsilon $	P_f/P_i
ν_y	3.161	2.255×10^{-4}	5.670×10^{-3}	0.599
$16-\nu_y$	4.989	2.314×10^{-4}	5.268×10^{-4}	0.996

Fig. 4.7.2-1. Spin depolarization from imperfection resonances in the booster.

for $n > 11$. This depolarization can be cured by "harmonic matching," that is, simply introducing vertical kicks that are adjusted to cancel ϵ for each value of n (Ch 83).

4.7.3 Depolarization in the Main Ring

Similar studies were performed for the main ring during acceleration from 9 to 45 GeV. In this case $\tau_a = 87.5$ ms in Eq. (4.7.2-2). The low superperiodicity of the racetrack ($S = 2$) means that 75 intrinsic resonances must be crossed. The values of P_f/P_i were calculated for the intrinsic resonances using a starting input emittance at 6.0 GeV of 0.5π mm-mrad. It was observed that $P_f/P_i > 0.999$ for all $\kappa = 2m - \nu_y$. Table 4.7.3-I lists the resonances, momentum, α , $|\epsilon|$, and P_f/P_i for the three cases with $\kappa = 2m + \nu_y$, where $P_f/P_i < 0.98$. Presumably tune-shifting quadrupoles could reduce these depolarizations to less than 1%.

TABLE 4.7.3-I

INTRINSIC SPIN-DEPOLARIZATION RESONANCES FOR THE MAIN-RING

Resonances	p GeV/c	α	$ \epsilon $	P_f/P_i
$16 + \nu_y$	11.712	6.20×10^{-4}	1.98×10^{-3}	0.980
$18 + \nu_y$	12.762	6.74×10^{-4}	6.01×10^{-3}	0.838
$58 + \nu_y$	33.717	8.73×10^{-4}	4.55×10^{-3}	0.927

The situation for the 73 imperfection resonances to be crossed is shown in Fig. 4.7.3-1 for $14 \leq n \leq 87$. We assume $\sigma_y = 0.2$ mm. Very little depolarization is expected according to this analysis. It must be stated, however, that the imperfection resonances do represent a formidable challenge just by virtue of their large number and close spacing.

Fig. 4.7.3-1. Depolarization in the main ring from imperfection resonances. These results may be optimistic.

4.8 ACCELERATOR RESEARCH AND DEVELOPMENT PLAN

4.8.1 Introduction

In Secs. 4.1 through 4.7 is presented the first iteration of a LAMPF II accelerator design report. Although much progress has been made, many holes remain. The ground rules of this first iteration included no effort at cost optimization. Beam-dynamics issues clearly need further study, particularly in the areas of space charge, collective effects, and active stabilization. We are a long way from being able to proceed confidently with hardware procurement. Although most of the proposed components are "conventional" by modern accelerator standards, few if any firms specialized in the manufacture of these components exist. Those who manufactured accelerator components in the past have largely disappeared. We must therefore work on the problem of technology transfer to our suppliers. All of these problems argue for a substantial up-front investment in research and development for LAMPF II. This research and development effort should include beam dynamics, cost optimization, development of prototype components, and technology transfer. In the sections which follow we will discuss some of the research and development topics which must be addressed. This is meant to be a representative list of topics, not a complete list.

4.8.2 Beam-Dynamics Issues Requiring Further Study

Beam-dynamics issues that affect the cost of the machines should have the highest priority. The dominant unresolved issue is that of the aperture of the magnets. We started with an aperture set by the vertical Laslett tune shift arbitrarily chosen at 0.20 and a large safety factor. An analysis of stability requirements led to a need to increase the beam size to assure transverse stability. Some safety factor still remains. Since the cost of magnets and power supplies is dominated by the aperture required, further study to refine the aperture requirement is needed at the earliest possible date.

A second major beam-dynamics issue affecting cost is beam loading. At the present time this work is incomplete. We must define the limits of stability of the rf system in the presence of errors, noise, and variation of beam loading due to nonuniform bucket loading. The required bandwidth of the rf amplifiers must be understood. The booster preamplifier bandwidth may need to be doubled, compared to the prototype under construction, in order to be able to cope with high-frequency effects (i.e., an additional 10 MHz). Before we modify the design for the preamp, we should have a better defined specification and some results from the full-power ferrite tests.

The vacuum chamber for the rapid-cycling magnets will be difficult to produce. The most difficult stability requirement to meet is the transverse impedance. The dominant term in the transverse impedance does not contain the resistance of the wall. With further study it is possible that a less expensive vacuum chamber could be specified.

An examination of the stability section indicates that the main ring would be easier to stabilize if the injection energy were increased. A larger main-ring acceptance in invariant phase space might be more economically obtained with a higher injection energy as well. A detailed study of the total cost of the accelerators as a function of the booster energy should be undertaken as soon as possible.

We have studied the impact of the accelerator design on performance of the polarized beam. A detailed study of the spin resonance-crossing hardware will be needed.

Once all of the requirements, including those of the polarized beam, are better defined, we must design a correcting magnet system. This system must be able to cope with all expected problems resulting from magnet field errors and misalignment. Considerable flexibility must be built in for programming the necessary corrections as a function of time and for handling unforeseen problems.

Numerous suggestions have been made for superconducting magnets for the main ring. No time has been devoted to following up on these suggestions. Part of the problem is that the power-supply cost and power consumption are tied closely to any decision about superconducting magnets. For the resonant power supply considered until recently, it would be necessary to eliminate the losses in the chokes and capacitors to take full advantage of superconducting magnets. Recent power-supply work has shown that a motor-generator power supply can be obtained economically, thus eliminating the choke and capacitor losses and providing more flexibility to consider superferroc magnets. Further study is required.

It is clear from this discussion that there is a pressing need for more beam-dynamics work immediately and a need for deeper study of effects not yet considered. We are developing a small group of accelerator physicists to work on these problems. This group will be invaluable for tolerance studies, tuneup scenarios, and in solving problems that will come up later during design, construction, turn on, and operation of the accelerator.

4.8.3 Booster Hardware Development Work

We have started one hardware development project, the booster rf system. From this project there has already emerged a major technical innovation in the use of perpendicular-biased ferrites, that eliminates altogether the high loss problems that have plagued previous proton accelerators in this energy range. A full-power amplifier and test stand are under construction, and the design of a full-scale cavity and tuner are nearly complete. Testing of the rf system is scheduled during calendar year 1985. A next step in this effort is to test the cavity with beam. We believe that it will be possible to test the beam-cavity interaction using the Los Alamos Proton Storage Ring (PSR) during 1986.

Several issues remain to be resolved before the rf design can be considered ready for fabrication. First, we must demonstrate that the cavity and ferrite can handle the necessary voltage and power without overheating or electrical breakdown. Then we must assure stability of the entire system under full beam loading. A cost trade-off study of ferrites that require lower bias should be undertaken. And finally, we must find and cure all the high-order modes that could cause beam-stability problems. Since many of these modes are three-dimensional in nature and involve the coupling of the amplifier to the cavity, only a full-scale prototype study will be successful. We are prepared to do all of this work, and adequate funding is available in FY-85. The estimated cost of the booster-cavity rf-development program from 1983 through 1986 inclusive is 1.5M\$. This is an indication of the scope of other required development projects.

Once we have firmed up the booster design, we will start the construction of a prototype combined-function dipole magnet, vacuum chamber, and power supply. Although the magnet is conventional, many technical challenges remain to produce a high-quality field. In particular, the control of assembly and alignment tolerances of the many thin laminations is important. Perhaps a grinding operation after assembly and curing of the epoxy bonding of the laminations would be possible. We also need to gain some experience in producing coils for rapid-cycling magnets. In particular, a study should be made of the trade-offs between small, directly cooled conductors and more finely divided indirectly cooled conductors (Litz cable). A magnet-measuring facility for rapid-cycling magnets should be established, and the precision of our magnet measurement techniques developed at the 10^{-4} level.

The vacuum chamber presents some unique problems. The main idea is to marry the technology of ceramic tubing with that of hybrid circuit technology (the conducting stripes) on a large scale. In our initial discussions with possible vendors, it is clear that no firm is willing to take on the whole job. We will have to find one supplier of tubing, another to make the joints, and a third to provide the coatings. And once we have fabricated a prototype, we will have to set up a facility for impedance measurement. This facility will be useful for all of the vacuum components of both rings and perhaps for other accelerator projects at Los Alamos as well.

The proposed dual-frequency resonant power supply for the booster has not been done elsewhere. During 1984 we had a subcontract with Walter Praeg at Argonne National Laboratory to test the basic idea and provide a proof-of-principle unregulated supply. We must now study the regulation problem and the effects that will occur when several subsections are connected into a ring (audio-frequency delay line modes). We must also consider the cross-talk and line-flicker problems and come up with an optimum solution to the charging supply problem.

Kicker magnets of the type required for the booster fast-extraction and fast-abort systems are not now in use at Los Alamos. Compared with the ferrite kickers in use at other laboratories, the LAMPF II kickers require a larger aperture and much higher repetition rate. Recent advances in thyatron technology indicate that high-power pulsers with 60-Hz capability and long lifetime are possible, but only if we break the kickers into reasonable length (1-m) modules. And from PSR experience, we have learned that good regulation and perhaps even active flattening of the kicker pulse are required if the kickers are not to determine the pulse-to-pulse stability of the beam in the horizontal plane. Finally, we must develop a shielded vacuum pipe passing through the kickers which presents an adequate impedance to the beam and at the same time does not strongly disturb the magnetic field of the kicker. It is important to design a build a prototype kicker module and power supply.

The beam diagnostic instrumentation for LAMPF II will be different from that in use elsewhere because of the need for low-impedance pickups, the high repetition rate, and the dynamic range required (from polarized beam to high current). A development program for beam-diagnostic instrumentation is required soon.

4.8.4 Main-Ring Hardware Development Work

The main-ring rf system is a less difficult problem than the booster, since a much smaller tuning range is required. Because a large number of units are required (40), it will pay to design a separate system rather than taking the option of using the booster cavities in the main ring. The main ring will also have a transient beam-loading problem during the injection period when the circumference is only partially filled. This requires a special feedforward system like that developed at Fermilab and at CERN. The main-ring rf development effort will be less difficult after the booster system is completed. We should therefore postpone main-ring rf development for about two years.

There are two problems to be studied for the main-ring magnets. These are (1) the extra problems in constructing 9.5-m-long laminated magnets (the booster magnets are 2 m long), and (2) the possibility of superconducting (superferric) magnets. If a decision is made to pursue development of superferric magnets, this will be a large undertaking for an organization that has no previous experience with superconducting magnets. Thus development of superferric magnets should be started at a very early date if they are to be used for the main ring.

The power supply for the main ring has undergone rapid changes. At the time of writing this proposal, it appears that the cheapest system consists of refurbishing the Argonne National Laboratory Zero Gradient Synchrotron motor generator set and using it with new electronic controls at LAMPF II. A thorough study of this option and the backup of using a large new generator as an energy storage device are required as soon as possible. It is especially important to assess the control and ripple problems, as slow extraction is extremely sensitive to magnetic-field regulation problems.

The slow-extraction system requires an electrostatic extraction septum. This septum must tolerate the large beam current that impinges on the wires and still must hold off the required voltage. Silicon carbide wires have been shown to be able to survive comparable currents in the LAMPF beam-monitoring harps and wire scanners. We must develop techniques for precision mounting of the brittle silicon-carbide wires. A test in the LAMPF beam would demonstrate performance under realistic conditions. The septum problem is similar to that of stretchers for dc electron machines, as the power density on the wires is nearly the same. Electron machines have an additional problem, because the septum polarity is such that field emission from the wires is possible. The availability of a high current from LAMPF and the need for septum development for LAMPF II might mean that a cooperative effort with Bates Laboratory or with CEBAF could be profitable for all involved.

A number of development problems are common between the booster and the main ring. In particular, the vacuum pipe, beam diagnostic, and magnet measuring problems are nearly identical. With our modular concept, the kickers can be common as well. Nevertheless, significant differences between booster and main-ring components can be expected to appear as we get further into the conceptual design, and development funding for work on the differences is expected.

Up to this point we have not mentioned the large need for computer software development for the control systems of the two machines. A new control system, distinct from the LAMPF computer control system will be required. Software

development will be very costly, and extrapolating from our own experience and that of other laboratories, we find that 50-100 man years' software effort will be required. A substantial part of this effort could be saved if a well-established control software system were obtained commercially or from another laboratory. In the near future we should study the available options. We should take note of the possibility of using the new PSR control system or the controls that are being done for the Brookhaven AGS. There will be special merit in obtaining a system from another circular accelerator in the same energy range as LAMPF II, since applications software might also be transferred.

4.8.5 Innovative Approach to Research and Development Work

It is important to consider innovative funding approaches for the R&D effort. For example, there exists no U.S. firm with recent experience in rapid-cycling magnet construction. It might make sense to fund several vendors to make a prototype magnet, then later choose one for constructing the production magnets on the basis of success with the prototype and the bid for production models. Properly handled, this could accomplish the required technology transfer, improve communications between the laboratory and vendors, and minimize the possibility of poor performance and litigation that might otherwise occur. The recent experience at Los Alamos with the PSR magnets indicates that this is a real problem that must be dealt with by proper management techniques.

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CHAPTER 5.

SCOPE OF THE EXPERIMENTAL AREA FACILITIES

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5. SCOPE OF THE EXPERIMENTAL AREA FACILITIES

5.1 INTRODUCTION

At this time it is possible to describe the scope and general features of the experimental facilities that should be provided for a broadly based exploitation of the LAMPF II accelerators. The requirements are clear, but specifying the details of a thoroughly optimized set of facilities would be premature. General features and major parameters are known and will be presented along with the directions to take and issues to be resolved in future work.

The physics opportunities described earlier require that the LAMPF II experimental areas should provide a number of high-intensity, high-quality beams simultaneously to several experiments. More explicitly, the major requirements are as follows:

- High-intensity beam
- High beam "quality," meaning
 - High brightness
 - High purity
 - High resolution for nuclear physics
- High beam availability
- Multiple ports for high throughput
- Access to beams of all energies and species that can be produced at LAMPF II, including ν , K , π , μ , $\bar{p}$, and polarized protons
- Flexibility to meet the requirements of future physics opportunities

A conservative set of facilities is described that meet these requirements using existing, proven methods and technology. The high-intensity facets are based on experience at LAMPF and the other contemporary meson factories, whereas the secondary beam-channel designs borrow heavily from past and present experience with kaon and $\bar{p}$ beams at other accelerators in the 6- to 30-GeV range. Existing LAMPF facilities are reused whenever feasible. Possibilities for significant enhancements of performance and/or cost reductions are presented along with an R&D plan for developing them. It is important to recognize that the experimental facilities described have not yet been optimized. They are included to show the scope and cost basis for facilities adequate to exploit the physics opportunities presented by LAMPF II. The details will change as physics requirements are refined, as ideas for improvements are studied and developed, and when the tradeoffs and optimization are completed.

5.2 CONFIGURATION AND GENERAL FEATURES OF THE EXPERIMENTAL AREAS

5.2.1 Layout

Three separate extracted-beam areas are envisioned, as shown in the layout of Fig. 4.1.4-1. In addition, space is left in each accelerator tunnel for an internal gas-jet target and associated experimental equipment. Fast-extracted beams from either accelerator can be transported to a target located below grade in Area N to serve the needs of ν and pulsed-muon experiments. However, initially Area N would be configured only for 6-GeV protons. Slow-extracted beam from the 45-GeV main ring would be split and transported to two areas, A and H, simultaneously. The new Area A, which reuses much of the building space and shielding from the present Area A, would house two primary beam targets and the lower energy secondary beams and spectrometers. Area H is a new experimental area for high-energy beams and spectrometers. As presently planned, Area H has two primary beam targets with the target cells located below grade to minimize shielding costs. The contents of each of these areas are discussed in subsequent sections.

5.2.2 Shielding and Remote-Handling Considerations

Shielding requirements have had a major impact on the experimental-area design. It is a major cost item and is closely coupled to maintenance and remote-handling design considerations. The amount of shielding in the vicinity of a thick target, a beam dump, or an area of significant spill is set by a maximum permissible dose to inhabitants or the background requirements of experiments, whichever is lower. Shielding estimates were made assuming a maximum permissible dose rate of 0.04 mSv/h for occupied areas. All shielding for 45-GeV facilities was designed for the full 34 μ A of beam that each area is capable of accepting, although it would normally receive less when operating simultaneously with other facilities.

The shielding thickness transverse to a thick target is primarily determined by the need to attenuate hadrons. Longitudinal shielding requirements are greater and are essentially determined by the requirement to attenuate high-energy muons, and will therefore depend to some extent upon the decay-path length available for production of muons. Preliminary estimates are about 21 hadron attenuation lengths for transverse shielding and 18 000 g/cm² or more for longitudinal shielding against muons (Le 84). The transverse requirements are about the same as presently in use at LAMPF I. Translated into iron shielding, these requirements are 3.7 m transverse and over 23 m longitudinal. From the large amount of forward shielding required, it clearly is desirable to minimize the number of beam stops to shield. Such considerations played a significant role in choosing a configuration of two thick targets in series per primary beam line, rather than the more flexible choice of several primary beam lines, each feeding only one thick target. It was also a consideration in designing for two forward charged-particle secondary beams per target in the low-energy Area (A).

In the present configuration of Area A, the secondary beam lines are 1.5 m above grade (floor level), and so it will be in the reconfiguration for LAMPF II. A cross section of the proposed shielding transverse to the primary beam is shown in Fig. 5.2.2-1. The 1.8-m-thick slab of magnetite concrete beneath the bulk shielding serves to reduce the "ground shine" of neutrons that

Fig. 5.2.2-1. Above-grade shielding configuration.

would otherwise penetrate the lower density tuff below grade. A layer of concrete approximately 1 m thick surrounds the iron on the top and sides and stops the low-energy neutrons that get through the iron.

The iron shielding above the target cells and main beam-line components is arranged in full-length vertical plates or slabs for easy removal by an overhead crane and for quick access by currently available LAMPF remote-handling equipment. The width (of the plates) transverse to the primary beam line is sufficient to provide quick access to any target-cell component by present remote-handling equipment. A section of the shielding along the beam line in a target cell is shown in Fig. 5.2.2-2. The vertical arrangement of plates with "key plugs" to break the otherwise uninterrupted stream path in vertical cracks between large slabs is shown. The lower portions of those shielding plates close to a production target are cooled by internal water passages connected to a water manifold at the top of the slab. Power and water services for target-cell components are brought in from above in stepped vertical channels located to the side of the removable vertical shielding plates.

Fig. 5.2.2-2. Target-cell removable-shielding configuration.

Moderate-sized, critical components that require more frequent replacement or service, such as targets, beam diagnostics, and protection instrumentation, are accessed through individual, fixed vertical penetrations. Their vacuum envelope extends to the top of the penetration. The component with its services and connections can be quickly removed or replaced through the penetration with the overhead crane and without need for extensive use of remote-handling

equipment. Horizontal penetrations are also possible and may be advisable if the overhead space becomes crowded with penetrations. They have the disadvantage of requiring additional retraction mechanisms and clear floor space for access. Nothing beats the flexibility of an overhead crane for handling heavy components.

The shielding and servicing concepts outlined above have evolved from the operating experience with the high-intensity beams at LAMPF and the other meson factories. Most are being implemented and extended in the current upgrade to the A-6 beam-stop area. Further development and refinement of these and other alternatives that promise to further reduce shielding and remote-servicing costs or downtime should receive strong encouragement.

In the above-grade approach shown in Figs. 5.2.2-1 and 5.2.2-2, all shielding is high density and therefore expensive. It is, however, well suited to the needs of low-energy kaon beams, which must be kept short to avoid excessive decay losses. For longer secondary beam lines, considerable savings in shielding costs can be obtained by placing the primary beam and target cells below grade, as shown in the cross section of Fig. 5.2.2-3. Access and remote servicing is through high-density shielding over the target cell and is essentially the same as for the above-grade approach in Area A. Shielding to the side and below the target cell and behind the beam stop is now tuff or soil, which is much cheaper. Some magnetite concrete poured in place or stacked near the top and to the side of the target cell is needed to provide at least 21 hadron attenuation lengths of shielding in all transverse directions. Savings of 50% or more in shielding costs are achieved by this approach, as compared with a comparable above-grade configuration. For this reason the primary beam

Fig. 5.2.2-3. Below-grade shielding configuration.

lines and target cells of the neutrino area (N) and the high-energy area (H) are proposed to be below grade.

5.2.3 Targets

The optimal secondary beam production target would be approximately one interaction length of the highest density material that will have a useful lifetime in the full-intensity primary beam. One interaction length is a good balance between the conflicting goals of having most of the primary protons interact to produce secondaries and having none of the secondaries lost to interactions in the target. Beyond one interaction length, the total flux of forward secondaries escaping the target starts to decrease despite the fact that more are produced in the target. The proposed configuration of the two targets of the order of one interaction length thick in series will ensure that most protons interact to produce secondary beams before reaching the beam stop. High density means shorter targets and reduced depth-of-field effects in the secondary beam-line optics. More fundamentally, high density means high source brightness, which is a most desirable starting point for producing high-quality beams.

Severe heat loads will be encountered in high-density, thick targets at LAMPF II. Calculations using CASIM (Va 75) indicate 80 kW of heat will be deposited in a 5-mm-diam, 150-mm-long (one interaction length) copper target bombarded with the full-intensity 45-GeV beam. Most of the heat arises from electromagnetic showers initiated by π^0 's produced in the target. Only about half of the shower energy is contained in this target. A one-interaction-length target of a refractory metal such as tungsten will absorb most of the shower energy, thus leading to a heat deposition of about 150 kW. Power densities such as these present formidable problems of heat removal and target lifetime. The loads mentioned above should be compared to the approximately 30-kW heat loads presently encountered at LAMPF I in 60-mm-long graphite targets bombarded by 1000 μ A of 800-MeV protons.

Two basic approaches to target cooling are being considered and both appear workable, given careful engineering. In one shown schematically in Fig. 5.2.3-1(a), the stationary finned target is water cooled. In the other, shown in Fig. 5.2.3-1(b), the heat from a rotating target is transferred by radiation to water-cooled chamber walls. A third, more conservative alternative would be a rotating water-cooled target.

Preliminary heat-transfer calculations (Br 84) indicate that a carefully optimized copper target with fins can be adequately cooled at one-half of full intensity. At full intensity, cooling is marginal, since interior temperatures near the exit of a one-interaction-length target are close to the melting point. Shorter targets are more feasible at full intensity, because power densities are lower owing to the fact that electromagnetic showers take several radiation lengths to develop maximum intensity. Higher density targets of this design such as tungsten are even more difficult to cool and may not be practical. The simplicity, small size, and absence of moving parts are the chief attractions of this design.

A rotating cylindrical target of a refractory metal such as tungsten can be cooled by radiation if allowed to operate at temperatures of around 2000°C or greater. For a one-interaction-length target (~100 mm in length) this implies a

(a) STATIONARY TARGET

(b) ROTATING TARGET

Fig. 5.2.3-1. Targets.

diameter of about 200 mm. Bearings that will operate reliably at these elevated

temperatures and in vacuum are one of the main difficulties with this approach. Careful attention must also be paid to the cooling of the chamber walls that receive the radiant heat from the target.

Of the three alternatives, a rotating water-cooled target will operate at the lowest temperatures. There is little question of its ability to transfer heat. The most serious question is one of reliability, related to its mechanical complexity and the rotating water seals. Nevertheless, this basic approach is presently being developed for the neutron-producing target of the proposed German spallation neutron source.

The discussion thus far has centered on targets for the slow-extracted beam areas. Targets for the fast-extracted beam areas face an additional complication of thermal shock from the ultrahigh peak power loads of a pulsed beam whose pulse length ($\sim 1 \mu\text{s}$) is about five orders of magnitude shorter. In this regime the temperature change, hence thermal expansion, develops in a short time compared to the time it takes sound to traverse the beam spot. This leads to coherent stress waves that can constructively interfere to produce high stresses in the target. The phenomenon is now encountered in neutrino and antiproton production targets at CERN and BNL. The peak intensity from the LAMPF II booster is comparable, but the time-averaged beam is two orders of magnitude higher. Thus we will have comparable dynamic thermal stresses, but they will be superimposed on a much more severe steady-state load. Accurate modeling of dynamic thermal-stress phenomena for the LAMPF II ν target is a major effort that has not yet been attempted.

Preliminary analysis shows that graphite is an acceptable target material for the ν area. Estimates of thermal shock using simplified analytical models (Si 74) indicate that graphite targets exposed to the booster beam with a spot size larger than 5 mm diam will survive the dynamic thermal stresses. Furthermore, LAMPF I experience with cooling of graphite targets shows that steady-state heat transfer is achievable. Higher density targets are not mandatory for the ν area, since the ν beam quality for the dual horn system is not appreciably affected by the extra length of graphite.

It is clear that solutions can be achieved for the targeting requirements of LAMPF II. Opportunities exist for significant improvements in both performance and reliability. Development of higher density targets is desirable, since they result in improved performance of most secondary beams. Further study and development of targets should be undertaken on a timely basis.

5.2.4 Target-Cell Configuration

We refer to the immediate region around a thick production target or beam stop as a target cell. It contains the nearby beam-line components, including the first elements of secondary beams, which are certain to become highly activated in normal service. It is the region where most, if not all, of the remote-handling activity is expected to occur. Most of the special problems associated with high-intensity experimental-area facilities are concentrated in the target cells. Once constructed and in service, they are not easily reconfigured and tend to become as fixed and as permanent as any of the accelerator facilities. Not surprisingly, much of our early thinking has been directed to the configuration and design of target-cell components upon which so much else depends.

A one-interaction-length target will result in a sizeable fraction of the incident beam power being scattered into nearby components in the form of secondaries and other reaction products. These in turn will develop as a heat and radiation load in these components. To illustrate the heating or dose rate, we have used the Fermilab code CASIM (Va 75) to calculate the power deposition in an idealized but highly relevant model consisting of a 150-mm (one interaction length) copper target surrounded by a 2-m-thick spherical iron shield beginning one meter from the target. A conical hole 50 mm in diameter at the start of the shielding was provided to allow the uninteracted primary beam to pass through. A set of curves giving power density deposited in the shield as a function of polar angle and radial depth in the shield is shown in Fig. 5.2.4-1. These can be used to estimate the heat and radiation dose to components in a target cell. In the model calculation about 40% of the incident beam power, or 600 kW, is deposited in the shield. For comparison, about 150 kW are scattered into the target-cell components by the present A-2 target at LAMPF I. Active measures such as water cooling of components and close-in shielding are required to remove this heat from the target cell; free convection will not be sufficient. Furthermore, the high radiation dose requires highly radiation-hardened materials and designs for these components. Fortunately, the LAMPF I experience has resulted in the development of many of the needed technologies.

In each target cell the first component after a target is a water-cooled collimator or heat shield whose primary function is to intercept and attenuate the scattered beam power. It serves to protect more sensitive downstream components such as magnets, flanges, and instrumentation by an order-of-magnitude reduction in the heat load and radiation dose they would otherwise receive. Thermal calculations show that a well-designed copper collimator will be adequately cooled if placed about 1 m downstream from a one-interaction-length target. The apertures in the collimator for passage of the primary beam are large enough (>5 standard deviation) that they intercept very little of the residual (uninteracted) primary beam, only its halo. In this sense they are not true beam collimators, but rather special water-cooled shielding elements coaxial with the beam. However, they serve to clip beam halo that might otherwise spill downstream.

One of our primary goals has been to gain the maximum utility from each target station. This has led to our present emphasis on target-cell configurations with more than one secondary beam per target. The real challenge is to obtain two small-angle, charged-particle beams per target with minimal compromise of secondary beam channel performance. We believe this is possible for the lower energy beams in the new Area A by adroit use of special-purpose magnets, such as half-quads and combined-function magnets with holes in the yokes or pole tips to allow unimpeded passage of the primary beam. They function to collect as much solid angle as possible and bend the secondary beam away from the primary beam as soon as possible. More conventional transport elements are used for the rest of the beam line.

An example with two forward charged-particle secondary beams is shown schematically in plan view in Fig. 5.2.4-2 for the first target cell, A-1, in the new Area A. Because the low-energy (<1.5 GeV/c) kaon angular distribution is relatively flat out to around 10° , very little intensity is lost by pick-off angles of $\sim 5^\circ$ although the horizontal source spot size projected to the coordinate system of the secondary beam is somewhat larger depending upon target length.

Fig. 5.2.4-1. Power density deposited in spherical iron shield due to a 34- μ A beam bombarding a one-interaction-length copper target 1 m from the shield.

Fig. 5.2.4-2. Secondary beam front sections in target-cell A-1.

In the target cell of Fig. 5.2.4-2, the water-cooled copper collimator downstream of the target absorbs the primary heat load of scattered radiation. It has three separate apertures, one for each of the three beams that pass through it. The next elements encountered are two side-by-side half-quads, one for each secondary beam, as shown in cross section in Fig. 5.2.4-3. The primary beam passes through an aperture in the mirror plate between the half-quads. The half-quads do both focusing and bending of the beam and are a special case of a combined-function bending magnet. Following the half-quads are additional combined-function bending magnets with the opposite sign gradient for additional bending and focusing. The net effect of the half-quad and the combined-function bending magnet is focusing in both horizontal and vertical planes and bending of the secondary beam away from the primary beam.

In the K0.8 beam a conventional quadrupole singlet, which follows the bending magnet, is still part of the target cell. Here the three magnets have the focusing properties of a triplet and are used to provide a dispersed focus downstream at momentum jaws and a vertical focus a little further downstream at a mass slit outside the target cell. A dc separator, which is located between the quadrupole and the momentum jaws, provides the first stage of separation.

Beyond the target cell, each secondary beam line can use more conventional transport elements. In the target cell, all magnets and diagnostics are radiation hardened. Where necessary, components are water cooled. In the layout

Fig. 5.2.4-3. Cross-section of target-cell A-1 half-quadrupole magnets.

shown for target-cell A-1, which has the collimator/heat shield absorbing much of the heat load, the magnet cores should need only surface cooling at most.

The primary beam is transported and refocused to the downstream target by radiation-hardened quadrupole magnets. Radiation-hardened primary beam diagnostics currently in use at LAMPF can be easily adapted for use in the experimental area of LAMPF II. Some of these on vertical penetrations will also be located in the target cells.

5.2.5 Maintenance, Remote Handling, and Access Considerations

High intensity influences the design in several ways, but perhaps none have greater impact than the constraints imposed by maintenance and servicing of equipment in the vicinity of a primary beam target. At LAMPF II, the activation of equipment will be roughly 50 times higher than for comparable areas at the AGS (Brookhaven). A maintenance activity at the present AGS that is limited to a few hours by radiation exposure to personnel will be limited to a few minutes at LAMPF II. In practice, this means that many maintenance activities are no longer hands-on, but require remote-handling capability. Remote-handling equipment is costly and remote servicing so time consuming and manpower intensive

that maintenance requirements become a dominant design consideration. Clearly, a primary design goal is to minimize the need for and duration of remote-handling activities.

Effective guidelines for minimizing costs and downtime for remote handling have been developed. They are:

(1) vastly increase system reliability so that repairs are less frequently needed;

(2) give careful consideration to the ease of component removal and replacement to minimize remote handling of failed components;

(3) minimize the time and effort to access critical or more frequently repaired components; and

(4) develop reliable and thoroughly tested remote-handling procedures as part of the design process to minimize the time lost to unforeseen problems once remote servicing is under way.

Each of these has been followed as far as possible throughout the design studies to date. Shielding layout and access to target-cell components have been greatly influenced by the desire to access components quickly and easily and, where remote handling is necessary, to use the concepts, equipment, and procedures developed and proven in practice at LAMPF I. Remote-handling access through distances prescribed by high-density shielding allow us to continue to use the equipment and procedures previously developed. To do otherwise would incur new uncertainties and possibly large development costs. Moderately sized components that need more frequent replacement or repair are on vertical penetrations for quick access and removal without need for full-blown remote handling.

Highly reliable component designs for the harsh environment of a high-intensity target are a long-term ongoing development effort at LAMPF. A good deal of it carries over to LAMPF II, but much remains to be done before the desired level of reliability is reached. Thorough prototype development and testing are needed to ensure that the final designs will have the necessary reliability of performance and ease of servicing.

Future development of the LAMPF II proposal will require continued attention to means of reducing the costs that are driven by maintenance and remote-handling requirements. In our studies thus far, we have resisted the temptation to reduce initial construction costs at the expense of future operating costs and beam availability. In this proposal and in future optimization of the design, we will continue to strive for a good balance between the two.

5.3 FAST-EXTRACTED BEAM AREA (N)

Fast-extracted beam from the booster initially and from the main ring eventually is transported below grade to an underground target station at the beginning of Area N as shown in the layout of Fig. 5.3.1-1. A thick (~400 mm) graphite or possibly beryllium target located mostly inside a focusing horn is nearly optimal for producing an intense ν beam. It will be possible to locate the first few centimeters of target in front of the horn enclosure to provide secondaries for a low-energy pulsed π or μ beam line.

Access and remote servicing is through high-density removable shielding over the target cell as described in Sec. 5.2.2. The shielding over the target cell accommodates both servicing needs and the need for an experimental hall to house low-energy pulsed beams, such as a pulsed-muon beam facility, which can be located either at grade or below grade at the target level. At this stage the preference is for a horizontal beam located at the target level.

5.3.1 Neutrino Facility

5.3.1.1 Experimental requirements. A neutrino facility of the kind we visualize for LAMPF II consists of three major components: a target cell and focusing system; the decay region and shield; and the experimental enclosure. The layout for the neutrino area based on in-flight decay of secondaries from a production target is shown in Fig. 5.3.1-1. The iron shield at the end of the decay tunnel serves to stop all particles (muons, neutrons, protons, etc.)

Fig. 5.3.1-1. Experimental Area N layout. Drawing of the 6-GeV, in-flight, pion-decay neutrino facility. Decay tunnel length = 40 m, muon and neutron shielding = 15 m, and detector center is at 60 m.

except neutrinos. The neutrino flux is typically increased 15 fold by employing a charged-secondary focusing device (normally a magnetic horn) around the production target. The desired charged secondaries (pions) are focused parallel to the decay tunnel, thus enhancing the probability of their flight path being long enough to allow a decay. From kinematics, the resulting neutrinos tend to be forward peaked. This effect is more pronounced as the energy is increased.

The decay region and shield. Neutrinos are produced most copiously at an accelerator through the decay of pions, muons, and kaons. The principal decay mode of the pion is $\pi \rightarrow \mu + \nu_\mu$, and in any practical detector the neutrinos are detected near the forward direction. In this case, the neutrino energy is given by

$$E_\nu = \frac{(m_\pi^2 - m_\mu^2)}{m_\pi^2 + p_\pi^2 \theta_\nu^2} p_\pi \quad .$$

At $\theta_\nu = 0$ and $p_\pi = 2 \text{ GeV/c}$, the muon-neutrino energy is roughly $E_\nu = 800 \text{ MeV}$. The neutrino spectrum reflects the pion spectrum passed by the focusing system, but somewhat broadened by this angular effect and weighted by the decay probability of the pions as a function of momentum. The pion decay length is 60 m at 2.0 GeV/c and an optimum decay length is about 40 m, although this optimum is very broad. The shield length must be sufficient to range out the most energetic muons from pions produced at the target in the forward direction, about 15 m for a 6-GeV proton beam.

Absolute normalization and ν_e contamination. With the above dimensions the approximate ν_e contamination is given by

$$\frac{1}{0.8} \cdot \frac{m_\mu d/2}{p_\pi c \tau_\mu} = 3 \times 10^{-3} \quad .$$

This contamination can be slightly reduced by decreasing the decay length at the expense of ν_μ flux. With nominal geometry, the resulting ν_μ spectrum will peak at a neutrino energy of 800 MeV, allowing the use of quasi-elastic scattering for normalization and spectral measurement. The cross section for the reaction

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

is $\sim 0.5 \times 10^{-38} \text{ cm}^2$ at the peak energy, and an acceptance of 0.25 could be

attained by a selection based on the outgoing muon angle. If we assume a proton current of 142 μ A, then the ν flux at the detector center (60 m from the target) is $4.5 \times 10^{11}/\text{m}^2\text{-s}$. If we further assume a detector with a 100-ton fiducial volume with equal numbers of protons and neutrons, then the rate of quasi-elastic events detected is approximately 2.5/s. We assume an experiment is the equivalent of 100 days of full beam running, which amounts to a total of 2×10^7 quasi-elastic events.

ν_μ e elastic scattering. We have taken the reaction $(\bar{\nu}_\mu)e \rightarrow (\bar{\nu}_\mu)e$ as our prototype reaction for the experimental area with the assumption that when LAMPF II is operational, there will be a continuing need to measure this reaction (and others) with sufficient precision to test the standard model or its derivatives with rigor. This experiment is conceptually simple and of fundamental significance in present physics, providing sufficient precision can be achieved. The cross section for ν_μ e scattering is $1.6 \times 10^{-42} \text{ E}_\nu$. If we assume an experimental acceptance of 0.5, then our 100-day-equivalent experiment yields 10^4 events. The yield of $\bar{\nu}$ events would be 4000 in a similar period, leading to a 2% statistical error in the ratio of cross sections.

A serious limitation on the accuracy of the experiment arises in the normalization process. The cross section for quasi-elastic scattering is well known on free nucleons, but the nuclear effects may be large for typical detector materials (carbon, aluminum, etc.). The use of deuterium may be required to improve the systematic uncertainty, and the event rate indicated above would not easily be maintained. Of course, 10^7 events in 100 tons will allow for attrition from detection constraints.

$\nu_\mu \rightarrow \nu_e$ oscillations. The ν beam produced by pion and muon decays has a contamination of ν_e of about 3×10^{-3} , as we estimated previously. The ν_e arise principally from the $\pi\text{-}\mu\text{-}e$ decay chain and from the Ke_3 decay mode. The spectra are calculable to good precision, providing the geometry of the beam is well defined and the ν_μ spectrum is well measured.

The probability that a ν_e will be observed in a ν_μ beam via oscillation is

$$P = \sin^2(2\alpha) \sin^2\left(1.27 \frac{\Delta m^2}{E_\nu} \ell\right) .$$

$\sin^2(2\alpha)$ is the mixing angle factor, Δm^2 is the difference of the ν_μ and ν_e masses in eV^2 , ℓ is the distance from the source in km, and E_ν is the neutrino energy in GeV. Then for values of the argument

$$1.27 \frac{\Delta m^2}{E_\nu} \ell \ll \frac{\pi}{2} ,$$

the amplitude for $\nu_\mu \rightarrow \nu_e$ goes as ℓ^2 . In order to optimize the sensitivity for observing oscillations for small Δm^2 , the detector should be far enough away so that the background from naturally occurring ν_e in the beam is at a low level. To verify that the background is as expected, we chose to detect 100 events in our canonical experiment. In a 100-ton detector the number of ν_e events in 100 days at distance d is $300/d^2$ (d in km) from both kaon decay and the π - μ - e chain. We can solve for d , giving $d = 1.7$ km.

There has been an indication that there may be an oscillation effect with $\sin^2(2\alpha) = 0.2$ and $\Delta m^2 = 0.2$ (Ko 84). Then the oscillation rate at 1.7 km is 0.05, which gives 800 events for the transition in the 100-day run. The position for which the oscillation transition is maximum is when the argument

$$1.27 \frac{\Delta m^2}{E_\nu} \ell = \frac{\pi}{2} ;$$

then $\ell = 4$ km, the background rate would be 20 events, and the oscillation signal 600 events. The experimental area we have outlined here produces a maximum flux at the close detector position for 6-GeV protons with a relatively small ν_e contamination.

5.3.1.2 The beam characteristics. The code for simulating the neutrino beam characteristics is a Los Alamos version of the CERN high-energy neutrino beam-line simulation code NUBEAM. The Los Alamos version has been extensively modified and now runs on the VAX. It has been validated by comparison with 30-GeV neutrino spectra from BNL experiment 734. NUBEAM is also capable of calculating the fluxes, energy distributions, and radial divergences for all species of neutrinos at any desired incident proton energy. It produces accurate distributions to be used by the experimenter in designing and optimizing an experiment. Estimates are that with a positive-focus horn system, an average ν_μ energy of ~ 0.8 GeV would be obtained and muon-neutrino flux increases of as much as $15\times$ over a bare target are possible. A Monte Carlo study of the neutrino flux using the pion horn focusing system has been completed and the results are given in Table 5.3.1-I. Figure 5.3.1-1 shows the calculated wide- and narrow-band ν_μ spectra from the two-horn system. Narrow-band (roughly monochromatic) neutrino spectra are obtained by placing plugs and collimators between the two horns such that a single pion momentum can be selected. Figure 5.3.1-2 shows π^+ trajectories through the horns at a current of 250 kA.

5.3.1.3 The pion focusing horns. A two-horn system was chosen to obtain the best magnetic focusing of pions. The first horn removes the depth-of-field problem associated with a long target (0.4 m). For pion trajectories that do not cross over between the horns, the first horn creates a virtual image of the target upstream shorter than the real target. For pion trajectories that cross over, a real but shorter image of the target is created between the horns at the focal point of the second horn. The second horn is then a point-to-parallel system.

TABLE 5.3.1-I

 ν RATES FROM A MONTE CARLO CALCULATION^a

	Flux ν/m^2 (π^+ Focus)	Ratio to Bare Target
Muon neutrinos	$4.5 (\pm 0.3) \times 10^{11} \nu/\text{m}^2\text{-s}$	~ 15
Electron neutrinos	$2.7 (\pm 0.1) \times 10^9 \nu/\text{m}^2\text{-s}$	~ 15
Antimuon neutrinos	$3.4 (\pm 0.6) \times 10^9 \nu/\text{m}^2\text{-s}$	
Antielectron neutrinos	$2.1 (\pm 0.7) \times 10^7 \nu/\text{m}^2\text{-s}$	
	Flux ν/m^2 (π^- Focus)	Ratio to Bare Target
Muon neutrinos	$6.3 (\pm 1.0) \times 10^9 \nu/\text{m}^2\text{-s}$	
Electron neutrinos	$1.8 (\pm 0.4) \times 10^8 \nu/\text{m}^2\text{-s}$	
Antimuon neutrinos	$3.1 (\pm 0.7) \times 10^{11} \nu/\text{m}^2\text{-s}$	~ 15
Antielectron neutrinos	$1.5 (\pm 0.4) \times 10^9 \nu/\text{m}^2\text{-s}$	~ 15

^aNormalized to 140 μA of 6-GeV protons on target, this table gives the neutrino flux averaged over a detector with a radius of 2.0 m located 60 m from the production target. The decay tunnel has a radius of 2.0 m and a length of 40 m. The horn current was ± 250 kA. Angular distribution of π^+ , π^- , K^+ , and K^- are from Sandford-Wang fits for 10-GeV/c incident protons on a beryllium target. Scaling to 7 GeV/c was accomplished by holding P_t constant.

Fig. 5.3.1-1. ν_μ spectra for positive horn focusing. A possible narrow-band spectra is also shown. The ν_μ flux was averaged over a detector radius of 2 m located 60 m from the production target. The incident proton beam was 140 μA of 6-GeV protons on a 0.4-m carbon target. The horn focus was positive (π^+), at a current of 250 kA.

DUAL HORN SYSTEM

Fig. 5.3.1-2. Trajectories of π^+ through a wide-band horn system. The horn current was 250 kA. Figure is not to scale.

5.3.2 Pulsed-Muon Beam Facility

For certain experiments, pulsed muon beams present tantalizing opportunities for improved results. These have been reviewed extensively in a recent design study aimed at exploiting PSR (He 83). For example, in μ -e conversion studies the motivation is an enormous reduction in prompt backgrounds induced by pion interactions. An example of a far different use is the line narrowing of muonium hyperfine levels from the use of "old muonium."

A pulsed-muon channel would be implemented in Area N. The muon channel could view the first several centimeters of graphite target used for the ν beam. The low-energy pions of interest for muon beams are produced more or less isotropically, hence any convenient production angle will suffice. To avoid interferences with the ν horn, a production angle of $\sim 80^\circ$ or larger is manageable.

The beam rates possible for a decay beam using the solenoid channel

TABLE 5.3.2-I

PULSED-MUON CHANNEL PARAMETERS AND RATES

Target = 4-mm graphite	$\Delta\Omega_\pi = 42 \text{ msr}$
5-m-long solenoid	$\frac{\Delta P}{P} \pi = 16\%$
π^- rate into solenoid = $7.6 + 10^9/\text{s}$	$I_p = 170 \mu\text{A}$
μ^- Rate at output = $2.4 \times 10^8/\text{s}$	

considered in earlier design studies (SA 82, HE 83) are shown in Table 5.3.2-I. These were estimated using data measured at BNL for 200 MeV/c π^- produced at 30° by 6 GeV/c protons incident on several different targets (Be 73). A reduction by about a factor of two was taken in going to 90° production, based on the observed angular distribution over the range $66-133^\circ$ measured for 200 MeV/c π^+ produced by 29 GeV/c protons incident on iridium and tungsten targets (Fo 79).

5.4 SLOW-EXTRACTED BEAM AREAS (A AND H)

The configuration and general features of these areas have been described earlier. In this section we will focus on the secondary beam lines from a user point of view. We will describe important design criteria, major design features, and essential performance parameters. Promising ideas for enhanced performance will be presented and assessed. A summary of the main parameters for these beam lines is presented in Table 5.4-I. Plan-view layouts for the two slow-extracted beam areas A and H are shown in Figs. 5.4.2-1 and 5.4.3-1, respectively.

In choosing the finite collection of beam lines for the slow-extracted beam areas, we have been guided by the major requirements discussed in the introductory Sec. 5.1. We have chosen a set that would provide beams of pions, kaons, and antiprotons of every momentum from 0.4 up to 35 GeV/c. Electrostatically separated beams are provided for kaons up to 6 GeV/c and antiprotons up to 10 GeV/c. With the successful development of rf separators one could have

TABLE 5.4-I

SECONDARY BEAM LINE PARAMETERS

Beam Line	θ_0 (degrees)	P (GeV/c)	$\Delta\Omega$ (MSR)	$\frac{\Delta P}{P}$ (%) (Res.)	$\frac{\Delta P}{P}$ (%) (Bite)	Length (m)	Separation
K0.8 hyper-nuclear & stopped kaon	~7	0.4-0.8	5	0.1	5	18	2-stage dc
K1.2 + EPICS II	~7	0.7-1.2	2	0.05	3	25	2-stage dc
K2 + spectrometer	~4	1.0-2.0	1.0	0.5	5	35	2-stage dc
K6 general purpose	0	2-6 K 2-10 $\bar{p}$	0.07	1	5	75	1-stage dc
K ⁰	3	wide-band	0.05	TOF	wide-band	40	neutron absorbers
K35 unseparated branch	0	up to 35	0.5	1	5	85	none
K35 separated branch	0	up to 35	0.05	1	1	~190	1-stage rf
Stopped muon	~60	0.2	42	5	15	~20	1-stage dc
Test beam	~2	up to 15	0.1	~1	~3	~40	none

separated beams up to 35 GeV/c. A K^0 beam, a stopped-muon channel, a high-energy double-armed spectrometer area, and a thin-target facility are also included. With a limited number of beam lines it is not possible to provide a beam optimized for and dedicated to every special purpose; some compromises were made so that all the important requirements are included. The result, we think, is a powerful set of facilities that are capable of addressing the broadest range of physics questions in this energy regime. The flexibility is there to adapt to new developments and opportunities that the future is certain to bring.

5.4.1 General Principles

While details of optimized designs are incomplete, it is possible to deduce many of the important performance parameters on the basis of general principles and past achievements in beam line design. Source-brightness characteristics, secondary-beam pick-off angles, beam purity, channel acceptance, energy range, and beam rates will be discussed.

5.4.1.1 Source-brightness optimization. For π , K, and $\bar{p}$ beams the high-brightness requirement is met by targeting full-intensity, full-energy beam on the highest density target that will survive while choosing beam-line pick-off angles in very forward directions. From the previous discussions of targeting, we can conclude that a useful water-cooled, fixed copper target can be developed. Higher brightness by a factor of ~ 1.6 can be obtained with a tungsten target, provided certain technical uncertainties with target cooling and operational reliability can be mastered. Pick-off angles for the K and $\bar{p}$ beams are chosen as close to 0° as is possible ($< 7^\circ$), consistent with reconstitution of the primary beam and the geometrical constraints of other secondary beams from the same target. Lower energy muon beams are less sensitive to pick-off angle considerations and therefore can effectively be located at large angles to the primary beam.

5.4.1.2 Beam-purity considerations. High beam purity for K and $\bar{p}$ beams is a crucial requirement that has major impact on our ability to use high intensity effectively. In general, some intensity will be traded off for higher purity. Our goal here is K/ π or $\bar{p}$ / π ratios better than 1:1. For some energies this requires π/μ rejection factors of 10^4 . Below 5 or 10 GeV/c, dc separators are effective; above this region we are forced to consider rf separators.

A single-stage electrostatically separated beam is, in practice, limited to π/μ rejection factors of 100-200 principally because of the halo of π 's and μ 's generated by decays of short-lived particles in the vicinity of the production target and scattering from pole tips and collimators. In principle, these can be eliminated by careful collimation and trimming of the source halo at an intermediate focus prior to the separators. A second stage of separation can accomplish the same purpose more effectively in a comparable length of extra beam transport. In fact, two-staged, separated beams for bubble chamber work have achieved π rejection factors as high as 10^5 (Eb 60, Bu 66). We choose two-stage separation, where rejection factors greater than 200 are required.

Electrostatic separators rapidly lose effectiveness as the secondary beam momentum, P , is raised. The maximum separable phase space in the separation phase plane falls as P^{-3} for $P \gg M$. As the momentum is raised, the maximum separable phase space becomes much smaller than could otherwise be transported. This effect sets a practical limit of around 6 GeV/c for K and 10 GeV/c for $\bar{p}$

beams, although in principle very small phase space (therefore low intensity) beams of higher energy can be separated.

Conventional rf separation techniques do not suffer from the same rapidly falling momentum dependence, basically because the crucial transverse deflections of the two species are in opposite directions rather than in the same direction, as is the case with electrostatic separators. Thus they are capable of providing π -K separation at all the higher energies available at LAMPF II. They do, however, suffer from two limitations:

(1) limited acceptance, because of the small beam apertures allowed in high-frequency cavities; and

(2) short duty factor, because of the very high power (10-20 MW) rf source required to achieve high fields in room-temperature cavities.

The geometrical limit on acceptance is tolerable, even though it is considerably less than what can otherwise be transported by a beam channel optimized for transmission. A reduction in acceptance by an order of magnitude is typical. Acceptance is more likely to be limited by the maximum deflection possible in the cavity.

More serious for LAMPF II is the duty-factor limitation, which falls short of the needed cw operation by two or three orders of magnitude (for room-temperature cavities). The development of reliable, high-field superconducting microwave cavities changes the picture dramatically. Since beam loading would be negligible, the rf power-source requirements become quite manageable. Further study and development of superconducting rf separators is highly desirable for LAMPF II. We assume the necessary R&D will be undertaken and therefore include this option in one leg of our proposed highest energy secondary beam line, K35.

5.4.1.3 Beam-line acceptance and length. In order to estimate secondary beam rates, we must know the channel acceptance and channel length. These may be estimated from general principles as follows. In separated beams the acceptance in the plane of separation (vertical) is usually limited by the separator deflection. In the dispersion plane (horizontal) the acceptance is more complicated, but is often limited by the momentum resolution required by the physics and the volume of high-quality magnetic field one is willing to provide. For rf separation both planes are likely to be limited by the separator geometrical acceptance with the deflection plane having about 1/3 the acceptance of the nondeflection plane (Be 63). These principles were used to set the solid-angle acceptance of the beams shown in Table 5.4-I. They are consistent with the numbers for existing beams at BNL and KEK (Ya 81, Bu 78).

Channel length is influenced by many considerations, and a rough layout is needed for making reasonable estimates. It is convenient for our analysis to separate a channel into three sections:

(1) a front-end or input section, which collects the secondaries from the target, sets the solid angle and momentum bite, bends this beam away from the primary beam, and "matches" it to the next section;

(2) a separator section, which purifies the beam; and

(3) an output section, which transports the beam and matches it to the users' requirements.

The length of all but the separator section typically scales with $p^{1/2}$ for fixed momentum resolution and fixed aperture magnets. Note that this implies a solid angle that goes as p^{-1} .

The length of the separation section scales differently. For electrostatic separation the maximum separable phase space, ϵ_m (deflection-limited case), is given by

$$\epsilon_m = \frac{1}{4} \frac{|f|}{R} eVL \cos \delta \approx \frac{1}{4} \frac{|f|}{R} eVL ,$$

where

$$\sin \delta = \frac{1}{2} \frac{|f|}{R} \frac{eVL^2}{a^2} < \frac{1}{\sqrt{2}}$$

and

$$f = \frac{1}{p^2} (W_2 - W_1) \approx \frac{M_2^2 - M_1^2}{2p^3} \text{ when } p^2 \gg M^2 .$$

V is the separator voltage, L its length, a the separator half aperture, R the separation factor (>1), W_i the total energy of the i^{th} particle, P the channel momentum, and e the charge of the electron. Given the strong dependence of ϵ_m on momentum, it is desirable to have the separator length increase with P to keep the acceptance from decreasing as rapidly as it would otherwise.

The formula discussed above also demonstrates one reason why we have chosen a factor of two or less variation in the useful momentum range of dc-separated channels. The magnetic-field optics most readily handle a fixed acceptance, which in this case would be set by the separator operating at the highest momentum. At the lower momenta the separator could handle much more phase space, but the rest of the channel cannot. To avoid a large mismatch between separator sections and the rest of the channel, the momentum range is restricted.

There is another reason to restrict the momentum range of the channel. Since kaons have a short lifetime, the channel length should be kept to a minimum to avoid excessive decay losses. Channel length increases with the maximum momentum. For separated beams this is a stronger function than for

unseparated beams. Thus, decay losses are optimized at the maximum momentum and are increasingly less than optimal at lower momenta. Once again, to avoid a large departure from optimal conditions, the momentum range is restricted.

For rf separation, the distance L between cavities sets the length and is given by

$$L = \frac{\lambda P_0^2}{M_2^2 - M_1^2},$$

where P_0 is the design momentum for a phase slip of π between the two particle species and λ the wave length of the rf field. To keep the losses from decay constant, lengths should not increase more rapidly than P ; therefore, λ should decrease as P^{-1} . For our high-energy beam line K35, we choose $P_0 = 20$ GeV/c and choose the rf frequency in the C band ($\lambda = 6$ cm). For $K-\pi$ separation these choices imply $L \approx 100$ m.

5.4.1.4 Production cross sections. To estimate particle-production cross-sections at 45 GeV, we have fitted a collection of data from 10 to 200 GeV with a slightly modified version of a model put forward by Hojvat and Van Ginneken (Ho 83) and used by them for estimating $\bar{p}$ production. In this model the invariant cross section for inclusive production from a nucleus is given by a product of three factors:

$$\frac{E}{\sigma_{\text{abs}}} \frac{d^3\sigma}{dp^3} = F_s \cdot F_{\text{aps}} \cdot F_{\text{nuc}},$$

where

(1) $F_s = K(1 - X_R)^m (1 + P_t^2/g^2)^{-4}$ is the cross section for inclusive production by p-p collision valid in the scaling region;

(2) $F_{\text{aps}} = 1 + B S^C \exp(A \cdot X_R)/(1.001 - X_R)$ is an empirical "approach to scaling factor" that accounts for deviation from scaling at lower energies; and

(3) $F_{\text{nuc}} = a(1 + bP_t^2) \exp(cX_R + dX_R^2)$ is an empirical factor that accounts for modifications due to nuclear matter.

The kinematic variables used are P_t , the transverse momentum (an invariant) of the secondary; the radial scaling variable $X_R = E^*/E_{\text{max}}^*$ introduced by Taylor (Ta 76); and S , the invariant total energy of the p-p system consisting of the incoming proton and one proton (or neutron) of the target nucleus. E^* is the total energy of the secondary in the p-p center of mass (c.m.) and E_{max}^* is the maximum c.m. energy of the secondary consistent with the conservation laws in p-p or p-N collisions.

The virtue of using radial scaling rather than Feynman scaling is that the scaling region is reached for lower values of S and is always approached from above. Taylor has shown that radial scaling is reached by $\sqrt{S} = 10$ GeV for all P_t and X_R , and sooner for small values of P_t and X_R .

Hojvat and Van Ginneken used similar expressions for F_S and F_{aps} , but with fewer adjustable parameters. They did not do a global fit, but instead used projections onto planes of $P_t = 0$ and $X_F = 0$.

The parameters in F_S are taken from fits to p-p data in the scaling region made by Taylor et al. (Ta 76, Jo 78). We have obtained estimates for the parameters A, B, and C in F_{aps} by least-square fits to p-p production data from 10 to 200 GeV. With the parameter of F_S and F_{aps} fixed by the program described above, we then fit the product formula to available production data from nuclei, varying only the parameters of F_{nuc} . For $\bar{p}$ production we have used the results of Hojvat and Van Ginneken directly.

Reasonable fits were obtained to a large collection of data that, unfortunately, includes no data between 26 and 67 GeV for production from nuclei. A sample result is shown in Fig. 5.4.1-1, where we compare the formula to some data for K^- production from beryllium for $P_t \approx 0$. We use the formula as a convenient means of summarizing available data and as a means of interpolating and extrapolating where there are no data. We judge that the formula is good to a factor of 2 for most of the region of interest at 45 GeV. It is less reliable for large values of $X_R > 0.8$ and for very low energy secondaries (in the lab system).

Low energies in the lab correspond to high X_R and backward angles in the p-target nucleon center-of-mass system. Symmetry about 90° in the p-p center-of-mass system is built into the formulas, and although it is required in the elementary p-p interactions, it is not required in p-nucleus interactions. Some enhancement of low energy (lab system) production is indicated in the few data points available. Some of our very low energy yields may be underestimated by a factor of 5.

5.4.2 Low-Energy Secondary Beam Area (A)

A fraction of the slow-extracted beam from the main-ring accelerator is transported to a reconfigured Area A via a beam tunnel that joins the present switchyard on the north side. The facilities included in Area A were chosen in part to make optimal use of existing features of this area. Area A with its above-grade shielding configuration and size is best suited for lower energy secondary beams where high-density shielding is used to advantage with relatively short beam lines. An estimated 80% of the existing shielding can be reused in our proposed configuration. Some simple extensions of the building space and crane coverage are also planned. A schematic plan view layout of Area A and its beam lines is shown in Fig. 5.4.2-1.

5.4.2.1 Thin-target facility. The present thin-target facility at LAMPF is served by a thin target ($<10^{-4}$ radiation lengths) located in the switchyard upstream of Area A. This facility will be preserved at LAMPF II for use with the 45-GeV slow-extracted beam. It will be possible to use thicker targets since multiple scattering of the 45-GeV primary beam will not be a problem for downstream users.

Fig. 5.4.1-1. K^- invariant cross sections for production from beryllium for $P_t \approx 0$. S- and P_t -dependent factors of the model were removed by multiplying both the model and the data by the factors shown in the ordinate label. Data are shown for $\theta < 3.6$ mr; $P_t < 0.63$ GeV/c. Most of the data are from $\theta = 0^\circ$ measurements. Solid curve is the model.

At some future time it may be desirable to relocate the thin-target area to a point upstream of all primary beam splitting to take advantage of continuous use of the full-intensity primary beam. Sufficient space will be left in the new switchyard to accommodate such a facility should it prove advisable.

5.4.2.2 K0.8 hypernuclear and stopped-kaon facility. This is a low-energy channel designed to meet the needs of high-resolution hypernuclear physics spectroscopy as well as provide high rates of stopping kaons for a wide variety of studies, including K-mesic atoms and rare decays. The useful momentum range of 0.4 to 0.8 GeV/c is optimal for these purposes. Two stages of separation are needed to bring the π/K ratio below 1. A momentum resolution of about 0.1% can be achieved by momentum analysis accomplished with the aid of data from multiwire chambers before and after the last bending magnet. The chambers in these locations are downstream of the separators and will therefore be exposed to beam rates that are manageable.

Fig. 5.4.2-1. Plan view layout of low-energy Area A.

At this incomplete stage of design, K0.8 is envisioned as a three-bend channel with an achromatic output focus. As discussed earlier, a kaon production angle of about 7° is feasible. An electrostatic separator and a mass slit are located after each of the first two bends. Overall length is about 18 m to the output focus. A spectrometer for analysis of reaction products completes the hypernuclear facility. The present EPICS spectrometer has about the correct momentum range and resolution to serve this purpose; however it has the disadvantage of a rather long flight path (about 8 m) from target to detector plane. Stopped-kaon experiments would run with the spectrometer moved out of the way.

The beam rates for this channel at its output focus, shown in Figs. 5.4.4-1, -2, and -3 were calculated assuming the parameters of Table 5.4-I, the full-momentum bite of 5%, and a 0.5-interaction-length target. A pion rejection factor of 10^4 was also assumed corresponding to two-staged separation.

5.4.2.3 K1.2 and EPICS II. Nuclear physics studies with kaons would not be complete without high resolution beams of momentum higher than provided by the K0.8 facility. The next momentum range is 0.7-1.2 GeV/c. A momentum resolution of 0.05% is the goal for K1.2 and its spectrometer (EPICS II), corresponding to a kaon energy resolution of 285 keV at 0.7 GeV/c and 555 keV at 1.2 GeV/c.

The current thinking is that this channel should operate as an energy-loss spectrometer with a vertically dispersed beam on the scattering target. This design has two advantages:

(1) Since resolution is not set by the channel momentum bite, a sizable bite can be used, resulting in higher beam rates.

(2) The dispersion plane and scattering plane are orthogonal, thereby decoupling momentum analysis from measurement of scattering angle.

A production angle of $5-7^\circ$ can be achieved in the A-1 target cell by the use of half-quads and combined-function magnets, as was discussed earlier in Sec. 5.2.4. Two stages of separation are needed for adequate beam purity. Channel length from production target to scattering target might be about 25 m.

Beam rates at the scattering target for this channel are also shown in Figs. 5.4.4-1, -2, and -3. They were calculated assuming the parameters in Table 5.4-I, the full-momentum bite of 3%, a 0.5-interaction-length target, and a length of 25 m. A pion rejection factor of 10^4 is assumed for the two stages of separation.

5.4.2.4 K2 and spectrometer. The next range of momentum in our progression of separated beams covers 1.2-2.0 GeV/c. In this range we anticipate the need for a more general-purpose channel that can meet the needs of programs studying hadron-nucleon interactions and those hadron-nucleus interactions where resolution of individual nuclear bound-state levels is not needed. We judge that a momentum resolution of about 0.5% will be adequate for these purposes. High purity implies that there be two stages of separation, which in this case would be accomplished with electrostatic separators.

This channel will view a one-interaction-length target in the second target cell in Area A. A beam pick-off angle in the range of $3-5^\circ$ will suffice, depending upon the details of working around a K^0 beam taken off at 3° on the same side of the primary beam. For this beam, as well as for $K0.8$, we imagine the optics to be based on a three-bend channel with an achromatic output focus and separator sections after each of the first two bends. The medium-resolution spectrometer being proposed for use in EPB at LAMPF I is well matched to the spectrometer requirements of K2, and we propose to reuse it for this purpose at LAMPF II. Some extension of the present Area A to the south and east is needed to house this spectrometer.

Beam rates shown in Figs. 5.4.4-1, -2, and -3 for this channel at its output focus were calculated using the K2 parameters of Table 5.4-I, the full momentum bite of 5%, and a one-interaction-length target. A pion-rejection factor of 10^4 is assumed for two stages of separation. When this channel is used at its best resolution of 0.5%, beam rates will be down a factor of 10, corresponding to a factor of 10 reduction in momentum bite.

5.4.2.5 K6 general-purpose beam line. This will be the highest energy charged-particle beam in Area A covering the range 2 to 10 GeV/c. As a general-purpose beam line, a momentum resolution of about 1% appears adequate. One stage of electrostatic separation is incorporated. It will provide adequate separation for kaons up to 6 GeV/c, but up to 10 GeV/c for antiprotons. Since A-2 is the last target in Area A, one of the charged-particle secondary beams

can be at 0° . We have chosen it to be K6 in order to optimize the rates of the higher energy antiprotons.

This channel will operate at the limit of what is practical with electrostatic separators. We assume that we can use dc separators at 1 MV, as has been achieved at KEK (Is 80). Even with these voltages it will require 20 m of separator length. With careful collimation and trimming of halo from the vertical source image at an intermediate focus ahead of the separator, we judge that a pion rejection factor of 500 can be obtained.

One improvement that might increase the rates at lower momenta is to provide a front end with larger acceptance than that limited by the separator at its highest operating momentum. This would be followed by a tunable matching section that would serve to properly fill the separator with more phase area, hence more beam, at lower momenta. This would have to be balanced against the decay losses from the extra length of the matching section. More study of this option is needed before a definite conclusion can be reached.

It should be noted that a cw rf separator might also work here. As was pointed out earlier, high-duty-factor rf separators require considerable development. Since electrostatic separators with the needed characteristics have been operated at existing machines, they are the more conservative choice. However, future development in rf separator technology might make them more attractive. In that case, the channel could be modified to use the new technology. It would not be necessary to modify the front-end elements in the target cell.

The optics of this beam could also be based on a three-bend design with an achromatic output focus. An intermediate dispersed focus after the first bend would be used to set the momentum bite and trim the halo on the vertical image of the source. The separator section would be located after the second bend and the mass slit located before the final bend, so that background particles from the slit are less likely to reach the final focus. Another desirable feature would be two output legs, each with an achromatic focus. This might be accomplished with optics similar to that employed in the P^3 at LAMPF I, but modified to accommodate a long separator. In this scheme the two output sections could not, of course, have beam in them simultaneously; rather, the beam is switched between them by the last bending magnet. In the long-duration experiments likely at such a channel, it is not efficient to use beam in one continuous run. One needs down time for analysis and/or modification of the equipment. The channel will be more efficiently utilized if two such experiments can be interlaced, as could be the case with the proposed two output ports.

The beam rates shown in the curves of Figs. 5.4.4-1, -2, and -3 for this channel were based on the parameters shown in Table 5.4-I, a full-momentum bite of 5%, a one-interaction-length target, and use of electrostatic separation. A pion-rejection factor of 500 in $K-\pi$ separation can be assumed for estimating beam purity.

5.4.2.6 The K^0 beam. The particle physics capabilities would not be complete without a well-designed K^0 beam. We think that there is enough room to pick off a neutral beam at around 3° (50 mr). This angle is chosen to reduce the neutron contamination relative to K^0 . Further reduction of neutron contamination relative to K^+ s would be done with absorber material that absorbs

neutrons more readily than K^0 's. Of course, the kaon rate will be reduced by absorption, but the purity will be improved.

The minimum channel length will be about 40 m. It takes a space of about 10 m between target and beam stop to accommodate the other beam lines, in particular the 0° beam, K6. About 25 m of beam stop and shielding are required to stop the forward high energy muons. An additional few meters are likely to be needed from the exit hole in the shielding to the center of the experimental setup. The maximum solid angle of the beam is determined by the channel length and the largest tolerable beam spot, which we judge to be about 0.3 m diam. In this way we estimate the maximum solid angle to be 50 μ sr. Adjustable collimating jaws upstream in the channel will allow the beam size to be reduced.

The K_L^0 beam rates and momentum spectrum shown in Fig. 5.4.4-4 were calculated using the beam parameters discussed above and shown in Table 5.4-I for the full solid angle. Production cross sections are taken from a scaling model fit to higher energy data (Ho 84).

5.4.2.7 Stopped-muon channel. High rates of stopping muons are possible at LAMPF II. Both decay-in-flight and "surface muon" beams are possible from a single channel viewing the one-interaction-length A-2 target at a wide angle. Design studies have been made incorporating the concepts developed for a pulsed-muon beam at PSR (He 83, Sa 82). Central features of this design include a superconducting solenoid for the decay section and two output ports. Channel parameters and μ^- beam rates are given in Table 5.4.2-I below.

5.4.2.8 Test beams. Every well-developed experimental area has test beams of modest cost and performance that can be used by experimenters to test and develop equipment. In Area A, a small-acceptance beam can be picked off at small angles from the A-1 target. It would be unseparated.

5.4.3 High-Energy Secondary Beam Area (H)

A new experimental area is planned for work with high-energy secondaries. As mentioned earlier, the two target cells will be below grade to minimize shielding costs. High-energy secondaries are produced mainly in very forward

TABLE 5.4.2-I

STOPPED-MUON CHANNEL PARAMETERS AND RATES

Target = one-interaction-length copper

5-m-long solenoid decay section

$$\Delta\Omega_\pi = 42 \text{ msr}$$

$$\frac{\Delta P}{P}|_\pi = 16\%$$

$$I_p = 34 \text{ } \mu\text{A}$$

$$\pi^- \text{ rate into solenoid} = 4 \times 10^{10} \text{ s}^{-1}$$

$$\mu^- \text{ rate at output} = 1.2 \times 10^9 \text{ s}^{-1}$$

directions. This fact, plus the need for long spectrometers and beam line, leads to a layout much longer than it is wide. A schematic layout for this area is shown in Fig. 5.4.3-1.

Beam splitting for the division of primary beam between Areas A and H would occur in the new beam switchyard after the slow-extraction channel of the main 45-GeV ring. Splitting would be done in the conventional manner with an electrostatic septum used to produce the initial separation. This would be followed by further separation in a Lambertson septum magnet some distance downstream of the electrostatic septum.

The beam scattered from the septum wires has the potential of causing spills that produce sufficient activation to compromise hands-on maintenance of the switchyard. We plan to use appropriately located collimation to absorb the scattered beam in a controlled manner and thereby localize any activation.

5.4.3.1 Double-armed spectrometer area. The first target cell in Area H would be configured to serve the needs of high-energy spectrometer studies. Both high- P_t studies of p-p scattering and studies of high-energy μ pairs in the Drell-Yan process require double-arm spectrometers. We are now in the process of examining the overlap in spectrometer requirements. There is a reasonable chance that the forward spectrometer arm needed for high- P_t work will also serve as one arm of the μ -pair spectrometer.

The spatial requirements are more certain; both programs require long (~ 40 m) spectrometers that operate in the forward direction. We should look at designs where the magnets need not move to change production angle of the observed particles. Angles out to several degrees are desired.

5.4.3.2 K35 high-energy general-purpose beam line. A high-energy general-purpose beam line that will transport the highest energy (~ 35 GeV/c) secondaries produced in usable quantities completes the proposed initial set of beam lines. As presently planned, it enjoys 0° production from the one-interaction-length target in the second target cell of area H. A transport section composed of quadrupoles and vertical bends would bring the beam from below grade to an experimental hall above grade where experiments are performed.

Fig. 5.4.3-1. Layout of Area H.

The front-end elements would be designed to accept as much phase space as possible. Downstream requirements such as those for rf separation may introduce lower limits, but the front end, especially those components in the target cell, should not be the limiting apertures. Future developments may change the downstream components to improve performance, but we plan that front-end acceptance should accommodate all foreseeable improvements.

Separation by electrostatic separators is not practical at the momenta covered by K35. The best hope for improving beam purity is the development of high-duty-factor rf separators along the lines discussed earlier in Sec. 5.4.2.2. Our plan is to assume that the necessary development will occur and be successful. Two output legs (ports) can be provided. In one we could have a large-acceptance unseparated beam of the highest possible intensity. The other output leg would be longer and configured to contain the rf separator section of about 100 m length. Intensity will be reduced by the extra decay losses and the acceptance of the rf separator. Pion rejection by a factor 10^3 is assumed, since it has been achieved elsewhere in rf separated beams (Fi 75).

Beam rates for both legs are shown in Figs. 5.4.4-1, -2, and -3. They were calculated assuming the parameters of Table 5.4-I, a one-interaction-length target, and full-momentum bite appropriate for each leg.

5.4.4 Secondary Beam Rates

We have gathered together in one place a collection of curves showing expected beam rates as a function of secondary beam momentum, P , for the various secondary beam lines. They were calculated using production cross sections estimated from the formula described in Sec. 5.4.1.4 and beam-line parameters given in Table 5.4-I. A 0.5-interaction-length copper target is used for K0.8 and K1.2 rates, whereas a one-interaction-length target is used for the remaining beam lines. For these curves the full beam intensity of 34 μ A at 45 GeV is assumed to strike each target, although in shared use approximately one-half that amount would be typical.

We show curves for K^- , π^- , $\bar{p}$, and K^0 rates in Figs. 5.4.4-1, -2, -3, and -4. Rates of π^+ are nearly the same as π^- , whereas K^+ rates are typically a factor of 3 higher than K^- . A plot of the K^+/K^- ratio at 0° as a function of P is shown in Fig. 5.4.4-5.

We show contamination for K^- beams in Fig. 5.4.4-6. We have not shown contamination for other beams as separate curves. They can be estimated by assuming that the pion rejection factor is 10^4 for two stages of separation in K0.8, K1.2, and K2; a factor of 500 for dc separation in K6; and a factor of 1000 for rf separation in K35.

Fig. 5.4.4-1. K^- beam rates; low-energy beams shown in upper figure and high-energy beams in lower figure.

Fig. 5.4.4-2. π^- beam rates; low-energy beams shown in upper figure and high-energy beams in lower figure.

Fig. 5.4.4-3. Antiproton beam rates; low-energy beams shown in upper figure and high-energy beams in lower figure.

Fig. 5.4.4-4. K_L^0 beam rates for K^0 beam line. Fig. 5.4.4-5. K^+/K^- at 0° .

Fig. 5.4.4-6. K^-/π^- ratio; low-energy beams shown in upper figure and high-energy beams in lower figure.

5.5 FUTURE DIRECTION OF EXPERIMENTAL AREA DEVELOPMENT FOR LAMPF II

In this proposal we have outlined the scope and general features of a set of experimental facilities adequate for a broadly based exploitation of the LAMPF II accelerators. It should not be considered a well-developed conceptual design, but rather a glimpse of what might be done, a good starting point for a sound conceptual design study. Future efforts will be devoted to further development of the conceptual design with emphasis on optimization, cost reduction, and the R&D needed to resolve outstanding technical issues. The goal of these efforts will be an optimized conceptual design including reliable cost estimates.

In the preceding sections we have identified several opportunities for possible cost reductions in shielding, remote servicing, and experimental area configuration. We consider cost reduction to include some consideration of future operating costs as well as initial construction costs. The optimization process will require the development of an information base that includes the functional dependence of cost and performance on important design variables. In some cases this will be done by development of the viable alternative designs to the point where expected performance is understood and reliable cost estimates can be made. Optimization is then a matter of choosing the alternative that best meets the cost and performance criteria.

Technical uncertainties in target design, high-duty rf separators, and a number of issues related to innovative secondary-beam design require additional in-depth research and development. These should begin soon to have beneficial influence on the conceptual design. In particular, the evaluation and development of high-field, high-duty-factor, superconducting rf separators will take considerable time and effort to come to fruition. Our first step will be a sound assessment of the present state of the art. Innovative beam design is especially needed for high-resolution, low-energy kaon beams required for nuclear physics with kaons. Although this problem has attracted the attention of a number of workers, a stronger effort is clearly warranted.

Cost estimates for the secondary beam lines are perhaps the most uncertain costs in this proposal. Parameters presented for the secondary beam lines were deduced from general principles rather than from specific designs. A major effort in the next year will be to produce more specific and optimal conceptual designs. We expect active consultation and collaboration with the user community in this area.

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CHAPTER 6.

COST ESTIMATE AND SCHEDULE

6. COST ESTIMATE AND SCHEDULE

6.1 COST ESTIMATE

The detailed cost estimate for LAMPF II was prepared by first making a work-breakdown structure and then costing each item in that structure. For each item we used recent experience at Los Alamos or other laboratories, vendor-supplied prices, or standard reference costs from construction cost-estimating handbooks. Labor rates applicable in the Los Alamos area were used for manpower cost estimates. All costs are given in January 1, 1985 dollars.

This cost estimate was prepared jointly with a consulting team led by Paul Reardon of the Science Applications International Corporation's Princeton office. In the course of making the joint estimate, numerous differences between the Los Alamos and SAIC estimates appeared. All of these differences were examined and resolved, usually in favor of taking the larger estimate. The work-breakdown structure, cost estimate, and backup justification are presented in SAIC reports that are available on request. Only the summary information is presented here.

The cost estimate includes installation, EDI, project management, and a large contingency. This contingency is appropriate for the present preconceptual design stage. Further work on the conceptual design can be expected to result in a reduction of the contingency.

The cost estimate includes the booster, main ring, and three experimental areas. These experimental areas include all the beam lines defined in Chap. 5. No detectors or spectrometers are included. The summary of the cost estimate is presented in Table 6.1-I.

6.2 SCHEDULE

We have prepared a detailed schedule for LAMPF II construction. This schedule, which is documented in the SAIC Management Plan, shows a seven-year construction period. The long construction time was chosen to minimize peak funding and manpower requirements.

The project will be completed in two stages. The low-energy stage, consisting of booster and Experimental Area N, will be completed in 3-1/2 years. The booster can be operated without interference to the construction of the remainder of the project, which will be done during the second 3-1/2 years. If we assume a 1988 construction start, then booster operation can commence in mid-1991 and full operation will occur in 1994.

In the proposed seven-year construction schedule, no shutdown of the LAMPF linac is necessary. Experimental Area A would have to be shut down to reconfigure it for high-energy beam. A three-year shutdown of Area A starting in 1991 is suggested.

Many other construction schedules are possible. Given sufficient funding and manpower, we could construct the complete facility in four years. At this early stage without funding information it is appropriate to discuss a range of scenarios.

December 18, 1984
KY/PCB/HAT

Table 6.1 - 1
LAMPF II
COST SUMMARY
(JAN. 1985 DOLLARS)

LOW ENERGY- STAGE I	MATERIAL	LABOR	TOTAL	ED&I	CM/PM	TOTAL	CONTINGENCY	TOTAL
CONV FACILITIES	\$ 21,853,219	\$ 7,261,009 ^a	\$ 29,114,228	\$ 4,487,003	\$ 2,523,939	\$ 36,125,170	\$ 8,874,830	\$ 45,000,000
HARDWARE/BOOSTER	35,630,340	6,769,765 ^b	42,400,105	12,010,000	6,005,000	60,415,105	14,584,895	75,000,000
EXPERIMENTAL AREA N	14,930,000	2,836,700 ^b	17,766,700	3,600,000	1,800,000	23,166,700	5,833,300	29,000,000
<hr/> SUBTOTAL	<hr/> \$ 72,413,559	<hr/> \$16,867,474	<hr/> \$ 89,281,033	<hr/> \$20,097,003	<hr/> \$10,328,939	<hr/> \$119,706,975	<hr/> \$29,293,025	<hr/> \$149,000,000
HIGH ENERGY- STAGE II								
CONV FACILITIES	\$ 16,490,688	\$ 9,223,197	\$ 25,713,885	\$ 4,114,221	\$ 2,314,250	\$ 32,142,356	\$ 7,857,644	\$ 40,000,000
HARDWARE/MAIN RING	67,314,650	12,789,784 ^b	80,104,434	10,950,000	5,475,000	96,529,434	23,470,566	120,000,000
EXPERIMENTAL AREAS	74,990,000	14,248,100 ^b	89,238,100	17,000,000	8,500,000	114,738,100	28,261,900	143,000,000
<hr/> A AND H	<hr/>	<hr/>	<hr/>	<hr/>	<hr/>	<hr/>	<hr/>	<hr/>
SUBTOTAL	\$158,795,338	\$36,261,081	\$195,056,419	\$32,064,221	\$16,289,250	\$243,409,890	\$59,590,110	\$303,000,000
TOTAL	\$231,208,897	\$53,128,555	\$284,337,452	\$52,161,224	\$26,618,189	\$363,116,865	\$88,883,135	\$452,000,000

^aINCLUDES AC SUBSTATION INSTALLATION (\$5,634,000 x 0.19) \$1,070,460

^bINSTALLATION @ 19% OF MATERIAL

LOS ALAMOS NAT'L LAB.
IS-4 REPORT SECTION
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