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The Electroweak Polarization Asymmetry: A Guided Tour

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ABSTRACT

A comprehensive review is provided of the electroweak polarization asymmetry at the Z°, a highly accurate measure of the Z° coupling to fermions. Its significance as a precision test of the Standard Model is explored in detail. Emphasized are the role of electroweak symmetry-breaking and radiative corrections; the non-decoupling of new physics beyond the Z°; and the testing of extensions of the Standard Model, such as supersymmetry, technicolor, new generations of fermions, grand unification, and new gauge forces. Also discussed are the relationship of the polarization asymmetry to other electroweak observables and its superiority to other Z° asymmetries. Experimental issues are briefly presented, stressing the importance of polarization at the SLC and LEP e^+e^- colliders.

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1. By Way of Introduction: Electroweak Polarization Physics

In their quest for discovering what the world is made of, nuclear and particle physicists have constructed an apparently successful theory of the fundamental forces and forms of matter, the Standard Model (SM). Based on the gauge symmetry SU(3)×SU(2)×U(1), the SM forces can be divided into two types. One is the strong force, experienced by quarks, with an underlying unbroken SU(3) gauge symmetry. The other is the electroweak (EW) force, affecting both quarks and leptons, whose gauge group SU(2)×U(1) is broken to the U(1) of electromagnetism. For the first time in a decade, pathbreaking tests of the SM will be carried out by a series of new accelerator facilities (SLC, LEP, KEK, Tevatron, HERA), probing the "321" theory in new energy ranges and searching for physics beyond the SM. The technique of electron-positron annihilation has acquired a special significance in illuminating the structure of the SM, EW physics in particular. e^+e^- collisions are "clean," with a well-understood initial state and subject only to EW interactions. These enjoy the property of being perturbative and thus calculable to arbitrary accuracy. The state-of-the-art e^+e^- annihilation will soon be provided by the new colliders at SLAC (SLC) and at CERN (LEP), both to study the Z weak neutral current resonance near 100 GeV center-of-mass energy. The key to detailed understanding of EW interactions is to take advantage of their violation of parity by the use of *polarized electron* (and positron) *beams.* Polarization will open up a bonanza of precision EW physics. In this lecture, I will try to explain the significance of one precious gem in this bonanza, the *polarization asymmetry* of the Z resonance, $A_{LR}(Z)$. While expounding this subject, I will draw on a now extensive body of recent literature, beginning with the seminal paper of Lynn and Stuart.[1–13]

What do precision tests of EW physics do for us? Since perturbative EW calculations can be carried out to any accuracy, careful comparisons of theory and experiment test both the EW SM and our cherished ideas about field theory in a way reminiscent of the Lamb shift and $g-2$ measurements. The *gauge* interactions of EW physics seem to be understood, but need to be checked beyond the current accuracies of a few percent. Tests of the weak gauge forces (the Z and the W) will also shed indirect light on the profound mystery of the SM, the *Higgs* sector, the source of EW symmetry-breaking (EW SB).[14] The SU(2)×U(1) symmetry is not manifest to us in everyday life — it is hidden by the Higgs mechanism. Not much is known about the Higgs sector of the SM. Its interaction with the gauge sector, giving mass to the Z and W, is controlled by the gauge symmetry. Its interaction with itself, producing the Higgs' own mass, is not understood. Nor is its interaction with the SM fermions (e^-, μ^-, quarks, etc.). The masses and number of generations of fermions are a complete mystery. To rationalize the Higgs sector, a great variety of untested physics has been postulated to accompany it: supersymmetry, technicolor, compositeness.[15] Different ideas about SB will all be tested by precision EW measurements in the 100 GeV region. EW SB has an important consequence for these experiments. Ordinarily, we expect the effects of heavy particles to "decouple" from measurements at energies well below their masses. In the presence of SB, this seemingly obvious property is evaded.[3,7] Through radiative corrections (higher-order loop corrections in perturbation theory), measurements in the 100 GeV region can tell us about physics far beyond this scale. This fact will prove of immense importance in testing the cornucopia of new physics proposed by theorists in the last decade and a half.

The polarization asymmetry $A_{LR}(Z)$ belongs to a select class of EW observables that are both theoretically important and experimentally measurable to high accuracy. Other such quantities include α (the fine structure constant), G_μ (Fermi's constant measured in μ decay), M_Z (the Z boson mass), M_W (the W boson mass), flavor mixing and CP violation in neutral mesons. These observables exhibit a sensitivity to radiative corrections and the EW SB not found elsewhere. The gauge coupling of the Z to fermions is an EW parameter of fundamental importance and directly controls the polarization asymmetry. The SLC and LEP polarization results will determine this coupling with unprecedented accuracy, a test of the SM far superior to current low-energy data. In the next section, I review the structure of the EW SM and the general properties of $A_{LR}(Z)$. Following a brief discussion of the experimental issues, I present some SM predictions for the polarization asymmetry and discuss extensively the study of new physics and EW SB through radiative corrections.

2. Polarization Asymmetry: Theoretical Structure

The first neutral-current polarization experiment was carried out by Taylor and Prescott at SLAC in 1978.[16]* Their polarization asymmetry measured the difference between e^- scattering from deuterons in the two electron polarization states. Since the scattering proceeds through the Z, it isolated the parity-violating coupling of the e^- to the neutral current. The other major test of this coupling has been neutral-current ν scattering from e^- beams and nuclear targets, the original way neutral currents were discovered in the early 1970's.[17] Together these experiments represent our knowledge of the weak neutral current as measured in low-energy scattering experiments, where the momentum transfer is small compared to the masses of the W and Z bosons (80–100 GeV). The weak charged current (the W) has of course been known for much longer: first discovered by Becquerel in β-decay in 1896, its theoretical significance was unravelled by Fermi in 1934 and reformulated in modern language by Feynman, Gell–Mann, Sudarshan and Marshak in 1958, after the addition of parity-violation by Lee, Yang and Wu.[18–21] The charged current was discovered first precisely because it is charge-changing; the neutral current for many years was hidden by background difficulties. Despite this historical gap in their discovery, the low-energy weak neutral and charged currents are essentially identical phenomena — interactions mediated by virtual heavy gauge bosons, "weak" at low energies because they are suppressed by the large boson masses. The production of real W's and Z's by the CERN ISR in 1983 confirmed, within errors, the predictions extrapolated from low-energy experiments.[22]

The polarization asymmetry discussed in this lecture is a direct descendant of the Taylor–Prescott asymmetry.[2] It measures the same parity-violating $e^- - Z$ coupling, but now on the Z resonance (90–96 GeV): s-channel e^+e^- annihilation to the Z (and the photon) carried at the SLC and LEP colliders in the next few years (Figures 1 and 2). With polarization, the e^- beams can be set in left- and right-handed helicity states, longitudinally polarized parallel or antiparallel to the direction of motion. (At high energies, the electron is essentially massless — chirality and helicity are identical.) The polarization asymmetry is then formed from the left- and right-handed annihilation cross sections:

$$A_{LR} = \frac{\sigma(e^+e^-_L \to f\bar{f}) - \sigma(e^+e^-_R \to f\bar{f})}{\sigma(e^+e^-_L \to f\bar{f}) + \sigma(e^+e^-_R \to f\bar{f})} \quad . \tag{2.1}$$

⋆ I ignore the atomic parity violation experiments performed in the late 1970's, as they were subject to substantial theoretical and experimental uncertainties. See Refs. 23 and 24 for discussion.

The asymmetry is defined for our purposes only to *charged* fermion-antifermion final states f *excluding* e^-. This restriction eliminates the presence of the t, or scattering, channel. We can measure $A_{LR}(Z)$ species by species, or simultaneously to all final states.

The weak neutral and charged currents are drawn together with the electromagnetic neutral current (the photon) in the gauge theory of Glashow, Salam and Weinberg, the EW SM.[25] The gauge group is SU(2)×U(1); the first being *left-handed weak isospin* ($\vec{I}$), and the second *weak hypercharge* (Y). Ordinary electric charge is $Q = I_3 + Y/2$. Each group has an independent coupling: g, g', respectively. These are usually expressed as the *electric charge*:

$$e^{-2} = g^{-2} + g'^{-2} \quad , \tag{2.2}$$

and the *sine of the weak mixing angle:*

$$\sin^2\theta_W \equiv s_\theta^2 = \frac{e^2}{g^2} = \frac{g^2}{(g^2 + g'^2)} \quad . \tag{2.3}$$

The three gauge bosons of SU(2) and the one of U(1) combine to form the four EW gauge bosons: photon, the Z and the $W^\pm$. The neutral gauge bosons are not separately from one group or another, but rather are mixtures of SU(2) and U(1). s_θ^2 measures this mixing. ($s_\theta^2 = 0$ implies no mixing.) The annihilation cross sections are formed from the neutral-current matrix element:

$$\mathcal{M}_{NC} = \frac{e^2 QQ'}{s} + \frac{e^2}{s_\theta^2 c_\theta^2} \frac{\left(I_3 - s_\theta^2 Q\right)\left(I_3' - s_\theta^2 Q'\right)}{s - M_Z^2 - i\sqrt{s}\,\Gamma_Z} \quad , \tag{2.4}$$

where (un)primed refers to final (initial) state couplings, $\sqrt{s}$ is the center-of-mass energy, and M_Z and Γ_Z are the mass and decay width of the Z. ($c_\theta^2 = 1 - s_\theta^2$.) The first term is the photon channel (Figure 1a), the second the Z (Figure 1b). Defining the left- and right-handed fermion-Z couplings:

$$g_L = I_3 - s_\theta^2 Q \ , \ g_R = -s_\theta^2 Q \quad ; \tag{2.5}$$

it is not difficult to show that:

$$A_{LR}(Z) \simeq \left(\frac{g_L^2 - g_R^2}{g_L^2 + g_R^2}\right)^e = \frac{2\left[1 - 4s_\theta^2\right]}{1 + \left[1 - 4s_\theta^2\right]^2} \quad . \tag{2.6}$$

Electrons, like all fermions, fall into the fundamental representation (**2**) of SU(2), so $I_{3e} = -1/2$, $Q_e = -1$. As promised, $A_{LR}(Z)$ is a direct measure of parity violation and the EW mixing s_θ^2. Note especially that it depends only on the initial-state couplings. Why is this so?

The key to understanding the power of $A_{LR}(Z)$ is appreciating the interplay of parity violation and Z resonance dominance. Let us begin at the tree level, with only Figure 1, and no radiative corrections. First assume that Γ_Z is zero, so that the Z is infinitely dominant when $\sqrt{s} = M_Z$. Then both the numerator and denominator of $A_{LR}(Z)$ contain only Z quantities at the pole. Since polarization is applied only to the *initial* state, the final-state couplings common to the numerator and denominator cancel. Hence $A_{LR}(Z)$ depends only on the initial state (when $\Gamma_Z = 0$), as we see in Eq. 2.6. This includes independence from the effect of *detector acceptance cuts* on the final-state fermions, when these are symmetrical between the forward and backward directions. If we now restore $\Gamma_Z \neq 0$, we lose perfect Z dominance and photon channel effects appear in the denominator, which *conserves parity.* This effect is slight, suppressed by $(\Gamma_Z\ /M_Z)^2 \simeq (0.03)^2$. Symmetrical detector cuts continue to cancel, but a small dependence on the final-state flavor is introduced (Figure 3). This dependence in understood and computable.[6]

The tree level is only the beginning, as quantum corrections play a crucial and interesting role in $A_{LR}(Z)$. Radiative corrections are conveniently classified into a number of classes.[3,8] The first are the "oblique," or vector boson self-energies (Figure 4a). The second are the "direct" corrections: initial and final weak vertices (Figure 4b with W's and Z's); initial and final QED vertices (Figure 4b with photons); weak boxes (Figure 4c with W's and Z's); and QED boxes (Figure 4c with at least one photon). The last are the bremsstrahlung corrections (Figure 4d). The bremsstrahlung and direct QED corrections must be combined into QED *initial, initial/final* and *final* sets. With quarks in the final state come strong-interaction effects: gluon radiation and the hadronization of quarks (Figure 5). Let us begin with final-state effects, following Lynn and Verzegnassi. The final-state QED and QCD corrections are both parity-conserving and would cancel under perfect Z dominance. Since $\Gamma_Z \neq 0$, these effects are present. The QED contribution is negligible; the QCD one is not, but can be bounded although not reliably computed. The final-state strong interaction effect introduces the first small theoretical uncertainty into the prediction of $A_{LR}(Z)$. Note that this error cannot be circumvented or improved upon. However, the use of final-state hadrons is still valid and no definition of jet axes is needed. I will return to the final-state error later. Apart from this minor problem, final-state effects continue to drop out.

Consider next *parity-conserving* initial-state effects, which we should expect to cancel in $A_{LR}(Z)$. The only one of this type is initial-state QED, vertex and radiation (Figures 4b and 4d). Initial-state radiation, unlike final-state QED, has a dramatic effect on individual cross sections, because it redistributes the energy spectrum of the initial state (Figure 6). On average, the e^+e^- beams radiate away a small part of their energy before annihilating. With perfect Z dominance, this correction would factorize and cancel in the ratio of $A_{LR}(Z)$.[6,26] In fact, were it not for the photon channel, $A_{LR}(s)$ would be constant. With the photon channel, initial-state QED has a small effect, calculable either analytically or by Monte Carlo (Figure 7).[8]

Please note that all of these wonderful cancellations disappear off Z resonance.

The polarization asymmetry would seem impervious to any correction. In fact, the corrections I have shown you so far do not alter the basic parity-violating coupling that controls $A_{LR}(Z)$. *Anything* that changes this coupling has a direct effect on $A_{LR}(Z)$. In the SM, this is accomplished by two very important loop corrections I have not yet discussed. The first is the initial weak vertex set — as they violate parity, they shift g_L and g_R by different amounts. New parity-violating particles coupled to electrons will enter here also.

The other set, essential to the understanding of EW physics, are the oblique corrections (Figure 4a). These consist of the photon self-energy (vacuum polarization) and the W and Z self-energies. A great deal of discussion has centered around these corrections, but their effects can be summarized in a simple way using the general properties of gauge theories, as shown by Kennedy and Lynn.[7] At the tree level, absent any oblique corrections, three (or four) independent parameters are needed to fix the EW gauge interactions. Two are the gauge couplings, e^2 and s_θ^2. The third is the scale of EW symmetry breaking, for which we use G_μ. The fourth is the "rho-parameter" ρ, which controls the relationship between M_Z and M_W. If the Higgs sector contains only SU(2) *doublet* vacuum expectation values (v.e.v.'s), $\rho = 1$ automatically. With non-doublet v.e.v.'s, ρ becomes arbitrary and model-dependent, but we assume only doublets. Now add the oblique corrections, which contain divergent parts. Carefully re-expressing the matrix element in terms of physically measurable quantities, we find that the divergences cancel. The theory is said to be *renormalizable.* To predict anything now requires, in addition to the tree-level inputs, *the oblique corrections due to all particles with EW couplings,* even ones too heavy to produce at the Z pole. The oblique corrections can be summed up in one stroke by using the "starred" running functions developed by Kennedy and Lynn. These are effective EW couplings the change, or run, with s (Figure 8). The gauge couplings become $s_*^2(s)$ and $e_*^2(s)$. G_μ becomes $G_{\mu_*}(s)$ and ρ, $\rho_*(s)$. (Note that even if $\rho = 1$, $\rho_* \neq 1$ because of radiative corrections.) Then:

$$\begin{aligned} M_Z^2 \rightarrow M_{Z_*}^2 &= \frac{e_*^2}{s_*^2 c_*^2} \frac{1}{4\sqrt{2} G_{\mu_*} \rho_*} \quad , \\ M_W^2 \rightarrow M_{W_*}^2 &= \frac{e_*^2}{s_*^2} \frac{1}{4\sqrt{2} G_{\mu_*}} \quad ; \end{aligned} \tag{2.7}$$

in the matrix elements, and now:

$$A_{LR}(Z) \simeq \frac{2\left[1-4s_*^2(Z)\right]}{1+\left[1-4s_*^2(Z)\right]^2} \quad . \tag{2.8}$$

We expect $s_*^2(Z) \simeq 0.21$–0.23, from low-energy νN, νe, and eD scattering; and the W and Z mass measurements.[27] Thus $A_{LR}(Z) \simeq 0.2$–0.3. Note the large amplification of shifts in $s_*^2(Z)$: $\delta A_{LR}(Z) \simeq -8\delta s_*^2(Z)$. $A_{LR}(Z)$ is now sensitive to all particles with EW couplings. I will discuss the significance of this fact in Section 4.

The best way to predict $s_*^2(Z)$ is from the Z mass:

$$s_*^2(Z)c_*^2(Z) = \frac{e_*^2(Z)}{4\sqrt{2}M_Z^2}\frac{1}{G_{\mu *}(Z)\rho_*(Z)} \quad . \tag{2.9}$$

SLC and LEP will measure M_Z to ± 50 MeV, a negligible error in $s_*^2(Z)$.[28] We need $G_{\mu *}(Z)$ and $\rho_*(Z)$, but these involve heavy particle contributions, which I postpone. $e_*^2(Z)$ is not affected by particles heavier than the Z. It can be computed by running the vacuum polarization (Figure 4a, with a photon) from $s = 0$ (where $e_*^2(0) = 4\pi\alpha$) to $s = M_Z^2$, if we know the masses and electromagnetic couplings of all particles between these two energies. For the hadronic resonances coupled to the photon between 0 and 10 GeV (ρ, ω, ϕ, etc.) the vacuum polarization contributions cannot be accurately computed from the present state of QCD. We are forced to use experimental data. The vacuum polarization can be expressed using a dispersion relation in terms of the measured $e^+e^- \to hadrons$ cross section. Unfortunately, these data contain uncertainties which creep from $e_*^2(Z)$ into $s_*^2(Z)$ and thence into $A_{LR}(Z)$.[5] A similar error occurs in the calculation of M_W from M_Z. The vacuum polarization error is the second theoretical uncertainty in the prediction of $A_{LR}(Z)$ and the more significant of the two. Unlike the final-state error, this error is remediable with better data and reanalysis. The recent work of Burkhardt et al., has improved our knowledge of the hadronic vacuum polarization and reduced the error.[29] It is:

$$\delta A_{LR}^{vacpol} = 0.002 \quad , \tag{2.10a}$$

while the final-state hadronization error is smaller:

$$\delta A_{LR}^{final} = 0.0005 \quad , \tag{2.10b}$$

leading to a total theoretical error of:

$$\delta A_{LR}^{theo} = 0.0025 \quad . \tag{2.11}$$

A reduction of δA_{LR}^{vacpol} to ± 0.001 might make the difference between barely detecting deviations in $A_{LR}(Z)$ and seeing them unequivocally. Reanalysis of the available e^+e^- hadroproduction data (a good thesis project!) or, better, a remeasurement of the hadronic resonances in the 0–10 GeV region (perhaps the new accelerator in China) would be a great service to precision EW physics.

Let us reiterate the main points. The polarization asymmetry at the Z pole, $A_{LR}(Z)$, directly measures the initial-state parity-violating coupling of the e^- to the Z. It is essentially free of final-state corrections and detector cut dependence; what effects are present, are calculable, except for final quark hadronization. This causes a slight error. Initial-state QED corrections are computable, but mostly cancel. *Any* modification of the parity-violating initial-state coupling changes $A_{LR}(Z)$. Initial vertex corrections do this, if they violate parity. Oblique corrections, the most important and interesting loop corrections, do also. The general effect of radiative corrections on $A_{LR}(Z)$ is summarized and compared with other Z quantities in Table I. $A_{LR}(Z)$ is expected to be in the range of 0.2–0.3. One oblique correction, the vacuum polarization, receives hadronic contributions that introduce a second and more significant theoretical error. The total theoretical uncertainty is ± 0.0025. The hadronic vacuum polarization depends on low-energy e^+e^- hadroproduction data, and the associated error could be improved by a remeasurement of the $e^+e^- \rightarrow$ *hadrons* cross section. Nevertheless, we are already able to make a serious test of the EW SM with $A_{LR}(Z)$, if we can measure it.

3. Some Experimental Issues

A brief look at the experimental aspects of $A_{LR}(Z)$ is worthwhile bringing the discussion down to earth.[13] The detection of the final M_Z decay products is no different with polarization; symmetric detector cuts cancel in $A_{LR}(Z)$, as do detector efficiencies. Three issues peculiar to $A_{LR}(Z)$ are: (1) *getting* a beam of polarized electrons; (2) *keeping* it polarized; and (3) *measuring* its polarization. Notice that the e^+ beam does not have to be polarized: the M_Z is a vector particle, and an e^- of a given helicity will interact only with an e^+ of opposite helicity.*

Linear accelerators, such as the SLC at SLAC, do not disturb the polarization of an incoming beam, so that the beam can be polarized before acceleration. An efficient source of polarized electrons to be used at the SLC is the gallium arsenide (GaAs) crystal. The electrons can be ejected from the crystal using a laser. They fall into degenerate triplet and singlet states in the crystal, so the resulting beam is not fully polarized. In fact, the maximum polarization possible is $P = 3-1/3+1 = 50\%$. (Full polarization may be obtained using a stressed uniaxial crystal like cadmium gallium arsenide, where spin-orbit couplings split the e^- states according to spin.) The SLC does have curved arcs in the final section of the collider. The beams are bent by transverse magnetic fields, which cause the electron spins to process. A potential disaster is deftly avoided by adjusting the length of the arcs so that the e^- spins have returned to their initial direction when they arrive at the interaction point.[10,13] Such a ruse cannot work in an electron synchrotron such as LEP. Initial polarization would be wrecked as the electrons circulate around in the transverse magnetic field. Instead, initially unpolarized e^- and e^+ beams will build up transverse polarization in the magnetic field. A last-minute B field can then rotate the polarization in the longitudinal direction at the interaction point. The major obstacle to this approach is the possibility of depolarizing resonances in the synchrotron.[12,30]

Since the polarization of the e^- beam is never perfect, we must actually measure it. For $P < 100\%$, the measured $A_{LR}(Z)$ is just $P \times A_{LR}(P = 100\%)$ $(0 \leq P \leq 1)$. The basic approach to measuring P is to scatter part of the e^- beam from a polarized target. Møller scattering of e^- from a magnetized iron target is an example. The beam interacts with the two outer Fe electrons, which are in definite spin states. The background e^-–Fe nucleus scattering limits this method to an accuracy of $\Delta P/P = 3\%$. Another approach is the Compton scattering of polarized laser light from the polarized e^- beam. This allows $\Delta P/P = 1\%$.[10,13,31]

⋆ This assumes only spin-1 channels in the annihilation. Colliding electrons and positrons of the *same* polarization might be an effective way of searching for new spin-0 states.

At the CERN LEP, the synchrotron arrangement means that the beam luminosity and energy can be measured quite well, by online beam sampling and the bending magnetic field. The SLC, being a one-pass collider, requires indirect approaches. Both the luminosity and energy will be checked by periodically extracting the beam for testing.[32] Neither can be measured as well at SLC as they can at LEP. $A_{LR}(Z)$ once again rises to the challenge. The absolute luminosity cancels in the ratio, although the relative luminosities from run to run must be known. Fluctuations in the beam energy cancel in $A_{LR}(Z)$ for the same reason that the initial-state radiation cancels. The numerous advantages of $A_{LR}(Z)$ make it much superior to the forward-backward asymmetry $A_{FB}(Z)$ as a measure of the fermion–Z coupling. At the Z:

$$A_{FB}(Z) \simeq \frac{3}{4}\left(\frac{g_L^2 - g_R^2}{g_L^2 + g_R^2}\right)^e \left(\frac{g_L^2 - g_R^2}{g_L^2 + g_R^2}\right)^f \quad . \tag{3.1}$$

$A_{FB}(Z)$ varies depending on the final fermion f. It can be measured cleanly only to $\mu^+\mu^-$ pairs; but these have a branching ratio at the Z of only 3%, increasing the statistical error. The wonderful cancellation of radiation and hadronization in $A_{LR}(Z)$ does not work for $A_{FB}(Z)$, making the hadronic $A_{FB}(Z)$ difficult to predict and measure. Initial-state radiation does not cancel, so neither will beam energy fluctuations — another source of error.[6]†

The final upshot for the experimental uncertainties in $A_{LR}(Z)$ is that only the polarization and statistical errors matter.

$$\delta A_{LR}^{exp} = \sqrt{A_{LR}^2\left(\frac{\Delta P}{P}\right)^2 + \frac{1}{P^2}\frac{1}{N}} \quad , \tag{3.2}$$

where N is the number of events in the sample.[13] An optimistic but feasible goal is: $P = 45\%, \Delta P/P = 1\%$, with $N = 10^6$ Z events. Note that all the hadronic data can be used. Then $\delta A_{LR}^{exp} = 0.004$. Hence, the total error, combining theoretical and experimental uncertainties:

$$\delta A_{LR} = 0.0065 \quad . \tag{3.3}$$

Since $\delta A_{LR}(Z) = -8\delta s_*^2(Z)$, this will measure $s_*^2(Z)$ to ± 0.0008, about an order of magnitude better than current low-energy data.[27] The combination of M_Z, M_W,

† Measuring $A_{FB}(Z)$ to b or c quarks would still be a good check on their quantum numbers. One can also define a polarized forward-backward asymmetry that measures only the final-state couplings (at the Z, just Eq. 2.6, with "e" replaced by "f"). This quantity is subject to many of the same limitations as the unpolarized A_{FB}.[33]

and $A_{LR}(Z)$ in fact constitutes a considerably more refined test of the SM than comparison of M_Z and M_W to νN scattering, for example. The superiority of $A_{LR}(Z)$ to $A_{FB}(Z)$ and low-energy data is summed up in Figure 9.[8] Reducing δA_{LR}^{theo} to ± 0.001 gives δA_{LR} of ± 0.005, a worthwhile improvement, as we shall see.

4. The Standard Model and Beyond: Broken Symmetry and Radiative Corrections

$A_{LR}(Z)$ is controlled by $s_*^2(Z)$, which can be calculated using Eq. 2.9. Rewriting it:

$$s_*^2(Z)c_*^2(Z) = \frac{e_*^2(Z)}{4\sqrt{s}G_\mu M_Z^2}\left[1 - 4\sqrt{2}G_\mu(\Delta_\rho(0) + \Delta_3(Z))\right] \quad . \tag{4.1}$$

Recall we need *three* inputs, α, G_μ, M_Z, plus the radiative corrections necessary to run these parameters to the Z pole. Using G_μ and eliminating $\rho_*(Z)$, I have to introduce the loop functions $\Delta_\rho(0)$ and $\Delta_3(Z)$. The significance of these quantities will become clear as we proceed. We can divide the radiative corrections into two groups. One class are the corrections due to particle states *lighter* than M_Z, which go into $e_*^2(Z)$, $\Delta_\rho(0)$ and $\Delta_3(Z)$. They are computable, either via field theory or, for the hadronic contributions to $e_*^2(Z)$, by a dispersion relation. They need to be checked, of course, but are not of vital interest. The real stakes are in the class of corrections due to particles *heavier* than M_Z. These contribute *only* to $\Delta_\rho(0)$ and $\Delta_3(Z)$, so these functions are central in exploring heavy particle effects in the SM. *All particles with EW couplings will contribute to* $\Delta_\rho(0)$ *and* $\Delta_3(Z)$, *regardless of their masses.*

Let us begin with the minimal SM: SU(2)×U(1) with three generations of fermions and a single neutral Higgs boson. The top quark and Higgs masses are unknown. Our first view of heavy particle effects can be seen in Figure 10, showing the effect on $A_{LR}(Z)$ of heavy top and Higgs masses, an amazing, counterintuitive result. Calculating the Lamb shift, for example, we do not need to know the top mass. This common-sense wisdom is enshrined in the Appelquist–Carazzone theorem: in an unbroken gauge theory, heavy particle effects at q^2 are suppressed by q^2/m^2, $m^2 \gg q^2$.[34] When the symmetry is broken, however, this wisdom fails. Heavy particles will affect physics at lower energies *if* this heaviness is due to a large *dimensionless* parameter that breaks a *global* symmetry.[35,36] The SM satisfies this criterion. The SB scale is set by $G_\mu \sim 1/2\sqrt{2}\phi^2$, where $\phi \sim 250$ GeV. The masses of all particles arise from multiplying ϕ by the appropriate dimensionless couplings. M_W, $M_Z \sim g\phi$, where g represents known gauge couplings. Note that this relationship is enforced by the the gauge symmetry. Fermion masses are $\sim G_Y\phi$, where G_Y represents arbitrary Yukawa couplings. The Higgs mass is $\sim \sqrt{\lambda}\phi$, where λ is the arbitrary Higgs self-coupling. The Higgs and fermion sectors possess no known symmetry controlling λ and G_Y, or the number of fermion generations. The reason for EW SB in the Higgs sector is itself an enigma.

The EW SM has two broken global symmetries. The first is well-known, the *global SU(2) custodial isospin symmetry*, which overlaps with the local SU(2)$_L$

gauge symmetry. If the Higgs sector contains only $SU(2)_L$ doublet (**2**) v.e.v.'s, the symmetry is preserved. Then $\rho = 1$, as mentioned in Section 2. Mass splittings in weak isospin multiplets break this symmetry. Their effect leaks into radiative corrections and controls $\Delta_\rho(0)$. This is the same quantity measured in the "rho-parameter" of low-energy neutral-current scattering:

$$\rho_*(0) = \frac{1}{1 - 4\sqrt{2}G_\mu \Delta_\rho(0)} \quad . \tag{4.2}$$

The measured $\rho_*(0)$ appears to differ slightly from one.[27] Invoking "naturalness," I will make the reasonable assumption that only doublet v.e.v.'s exist and $\rho_*(0) \neq 1$ because of calculable radiative corrections in $\Delta_\rho(0)$. The top mass effect is due to its splitting from the bottom quark:

$$\Delta_\rho(0) = \frac{3}{64\pi^2}\left[m_t^2 + m_b^2 - \frac{2m_t^2 m_b^2}{m_t^2 - m_b^2}\ln\left(\frac{m_t^2}{m_b^2}\right)\right] \quad . \tag{4.3}$$

The top effect is not only not suppressed, but grows quadratically! The Higgs effect is not so dramatic:

$$\Delta_\rho(0) = -\frac{3}{64\pi^2}\left[M_Z^2 \ln\left(\frac{m_H^2}{M_Z^2}\right) - M_W^2 \ln\left(\frac{m_H^2}{M_W^2}\right)\right] \quad . \tag{4.4}$$

The isospin splitting in this case is $M_W \neq M_Z$, a result of the mixing between $SU(2)_L$ and $U(1)_Y$. (Note the opposite sign characteristic of gauge bosons.) The other broken symmetry of the EW SM is *global chiral symmetry,* the independence of the left- and right-handed components of the fermion fields. Once the fermions acquire a Dirac mass, the two components are mixed and chiral symmetry is broken. This breaking is summed up by the quantity $\Delta_3(Z)$. Unlike the isospin effect, the chiral breaking contributions of heavy fermions are constants, independent of mass, and non-zero *even if the fermions are degenerate.*[3,7] $\Delta_3(Z)$ has never been measured before and cannot be measured at low energies, because $\Delta_3(0) = 0$. It represents new information about EW physics.

If we extend the SM with new EW multiplets of scalars and fermions ("matter"), maintaining the $SU(2)\times U(1)$ gauge structure, the relationship Eq. 4.1 is unchanged. The effect of new matter heavier than the Z is contained completely in $\Delta_\rho(0)$ and $\Delta_3(Z)$ and can be understood in terms of the broken global symmetries. A potpourri of new physics is shown in Table II.[3,8] New isospin multiplets announce their presence if they have large mass splittings, through $\Delta_\rho(0)$. Such is the case for a new fermion generation or for the squarks and sleptons of supersymmetry. As $\rho_*(0) \simeq 1$, large isospin splittings in the SM are already ruled out.

This limits m_{top} to no more than about 180 GeV.[27] Also, most technicolor theories are eliminated, because they require a large number of pseudo-Goldstone bosons with substantial splittings. The question of splittings is more subtle: scalars and fermions contribute to $\Delta_\rho(0)$ with the same sign, but the gauge boson-Higgs contributions (Eq. 4.4) have the *opposite* sign (but only a weak logarithmic dependence on the Higgs mass). The comparison of $A_{LR}(Z)$ to experiment through Eq. 4.1 will place stronger bounds on splittings, equivalent to measuring the rho parameter to $\pm 0.3\%$ (splittings $\sim M_W$).[7] Even when splittings are turned off, the chiral effect is clearly seen in Table II. The contributions of heavy degenerate scalars disappear, but the fermions, Cheshire-like, persist. Such is a new generation of degenerate fermions or a set of supersymmetric gauginos and Higgsinos (listed as "Winos"). From this type of radiative correction data, we cannot deduce the theory of the world, only place limits on hypothesized models. The most general analysis will allow us to separate $\Delta_\rho(0)$ and $\Delta_3(Z)$ without any further discrimination among models. To do this, we need a second independent measurement. The obvious choice is M_W, which will be measured by LEP2 to approximately ± 100 MeV.[37]

$$M_W^2 = \frac{e_*^2(W)}{s_*^2(W)} \frac{1}{4\sqrt{2}G_\mu} \left[1 - 4\sqrt{2}G_\mu \Delta_1(W)\right] \quad . \tag{4.5}$$

$e_*^2(W)$ and $s_*^2(W)$ are computable using low-energy physics, once $s_*^2(Z)$ is known. $\Delta_1(W)$ plays the same role in charged-current interactions that $\Delta_3(Z)$ does in the neutral current. For heavy fermions, $m_f^2 \gg M_W^2, M_Z^2$:

$$\begin{aligned} \Delta_1(s) &= C_1 s \quad , \\ \Delta_3(s) &= C_3 s \quad . \end{aligned} \tag{4.6}$$

Assuming no substantial isospin splitting exists in the SM, $C_3 = C_1 = 0$ for degenerate scalars; while for degenerate fermions, $C_3 = C_1 \neq 0$. Measuring M_W thus gives us C_3, hence $\Delta_3(Z)$; combining with $A_{LR}(Z)$, we can separate out $\Delta_\rho(0)$. With C_3, we can *bound the total number of fermion multiplets with weak isospin.* Each multiplet of isopsin i and N_C colors contributes:[7]

$$C_3 = -\frac{i(i+1)(2i+1)N_c}{144\pi^2} \quad . \tag{4.7}$$

This bound is more restrictive than counting ν's with the Z width, as C_3 responds to fermions of arbitrary mass. For example, with the quoted errors for $A_{LR}(Z)$ and M_W, we can count the total number of heavy doublet fermion generations to a resolution of about ± 1.5. The two functions $\Delta_\rho(0)$ and $\Delta_3(Z)$ represent the complete knowledge of EW SB obtainable in low-energy four-fermion processes. Qualitatively new radiative correction effects, arising from the same broken symmetries, will appear in other processes, such as $e^+e^- \to W^+W^-$.[38]

To generalize beyond the SU(2)×U(1) structure, we can go in two different directions. One is to *unify* the low-energy gauge group into a single Lie group, *grand unification.* This effectively occurs only at a very high energy, typically 10^{15}–10^{18} GeV, called the *unification scale,* M_X. By relating the $SU(2)_L$ and $U(1)_Y$ groups, grand unification predicts a value of $s_*^2(Z)$, given M_X and the radiative corrections due to particles between M_Z and M_X.[9,39] $s_*^2(Z)$ is now sensitive to M_X, a new scale of SB. The simplest grand unified group is the minimal SU(5) of Georgi and Glashow, now ruled out by proton decay bounds.[40] Nevertheless, the general idea remains valid with other models and is a subset of more ambitious theories that include gravity, such as superstrings. In Figure 11, I show some results for supersymmetric SU(5) and superstring-inspired E_6 models, together with minimal SU(5) for comparison. The SUSY models include new matter multiplets taken at a common mass μ. The second direction away from SU(2)×U(1) is to *diversify,* by adding new low-energy gauge groups to EW interactions. The most popular of these are again superstring-inspired E_6 models, broken down to $SU(2)\times U(1)\times U(1)'$ or to $SU(2)_L\times SU(2)_R\times U(1)_{B-L}$.[4] Both models predict a new neutral-current resonance, the Z'. The Z' will mix with the Z and destroy the simple relationship between $s_*^2(Z)$ and $A_{LR}(Z)$ (Eq. 2.8) by $O(M_Z^2/M_{Z'}^2)$ effects. $M_{Z'}$ is restricted to be heavier than about 110 GeV, from current low-energy scattering limits, depending on the mixing.[27] But how do we separate a Z' from SU(2)×U(1) radiative corrections? The trick is to compare two different quantities that in SU(2)×U(1) are controlled only by $s_*^2(Z)$. We cannot use M_W, but $A_{FB}(Z)$ to some final-state species is a candidate. Shifts in $A_{LR}(Z)$ and $A_{FB}(Z)$ from the predicted GSW values due to the same oblique corrections are related in SU(2)×U(1) because of the common coupling $s_*^2(Z)$ between them. Thus:

$$a \cdot \delta A_{LR}^{GSW}(Z) + b \cdot \delta A_{FB}^{GSW}(Z) = 0 \quad , \tag{4.8}$$

where a, b are some constants depending on the final-state fermion quantum numbers. Such a "sum rule" as Eq. 4.8 tests the simple SU(2)×U(1) "star-everything-oblique corrections." If Eq. 4.8 is *not* zero experimentally, we have two possibilities:

(a) SU(2)×U(1) is the correct gauge group, but some new *direct* loop corrections are present (e.g., new Yukawa couplings); or

(b) SU(2)×U(1) is *not* right — a Z' exists. The non-zero value of Eq. 4.8 can tell us the underlying gauge group.

In Figure 12 are the values of the sum rule for the two different Z' models, using $A_{FB}(Z)$ to c and b quarks. Implementing this idea is not as easy as it sounds. The sum rule is zero automatically for $A_{FB}(Z)$ to mu pairs if the Z' coupling to leptons is universal. We are forced to use $A_{FB}(Z)$ to quark pairs, such as $b\bar{b}$ or $c\bar{c}$, with possibly large statistical and systematic errors. The difficulties are great but not insuperable.*

As promised, $A_{LR}(Z)$ is the key to a wide range of EW physics. Within in the SU(2)×U(1) gauge structure, the physics of heavy particles is linked through radiative corrections and SB to the measurement of M_W and $A_{LR}(Z)$. The general SB structure of the SM, including broken isospin and chiral symmetries, can be summed up with the two functions $\Delta_\rho(0)$ and $\Delta_3(Z)$. In the minimal SM, these can fix ranges for m_{top} and m_{Higgs} (Figure 13). $A_{LR}(Z)$ can be used to search beyond the SM, testing the prediction of grand unified models and to probe for a Z'. The non-decoupling of heavy particles in radiative corrections inverts experimental limits on new particles in a curious way — particle masses can be bounded from *above*. But having constrained EW SB, new gauge groups, and grand unification, we are still left with the question of why the Standard Model explains the low energy world as well as it does.

⋆ A_{FB} is a measure of C violation (asymmetry between e^+ and e^- beam directions), while A_{LR} measures P violation (asymmetry between e^- polarization states). In the gauge structure of the SM, C and P violation are linked so that CP is conserved. One can construct a CP-violating asymmetry using particles and *anti*particles.[41] In the SM, CP violation is confined to the Higgs-fermion sector.

5. Conclusions and Prospects

The polarization asymmetry is thus an extraordinary piece of information to have about EW physics, even more powerful when combined with the W mass and the forward-backward asymmetry of the Z. Its importance underscores the crucial need for polarized e^- beams at both SLC and LEP. The analyzing power extended to these machines by polarization presents a unique opportunity for precise tests of the Standard Model. The constraints placed on new physics from the SLC/LEP results will influence the physics programs of the next generation of accelerators. Polarization naturally reappears in plans for new e^+e^- colliders, as the preferred way to study new gauge couplings and constrain still higher energies via radiative corrections and symmetry-breaking.[42]

Peeling the layers of the Higgs mechanism will be the great challenge of particle physics for at least the next decade. Our understanding of gauge symmetry is good enough that we can imagine grand unification and new gauge forces of all sorts. But we have no principle to guide us in breaking these symmetries. Similarly, we lack good reasons for the different fermion families and their masses. We understand forces, but not matter. Unwrapping the mystery of electroweak symmetry breaking is the next decisive step towards the final goal of particle physics, understanding all matter and energy in terms deeper and simpler than what meets the eye from the world around us.

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TABLE CAPTIONS

All tables from Ref. 8.

I. Effect of radiative corrections at the Z^0.

II. Sample of new physics effects on electroweak measurables.

TABLE I

Effect of Radiative Corrections on Z Peak

Example: $M_Z = 94$ GeV, $m_{top} = 60$ GeV, $m_{Higgs} = 100$ GeV

	$\sigma^{\mu^+\mu^-}$(pb)*	Peak (GeV)*	$A_{\rm FB}^{\mu^+\mu^-*}$	$A_{\rm LR}^{\mu^+\mu^-}$
Tree level	1930	94.00	0.129	0.414
Changes due to:				
Initial QED	−500	+130 MeV	−0.014	−0.004
Final QED	+20	none	negligible	negligible
I/F QED	negligible	negligible	negligible	negligible
Oblique	−55	negligible	−0.073	−0.136
Initial weak	negligible	negligible	−0.001	−0.005
Final weak	negligible	negligible	−0.001	negligible
Weak boxes	negligible	negligible	negligible	negligible
Uncertainties due to cuts:				
Endcap	±0.4%	negligible	±0.001	negligible
Acollinearity	negligible	negligible	negligible	negligible
Realistic experimental goals:				
SLC/LEP	±3%	±50 MeV	±0.05	±0.004

* Unpolarized.

"Negligible" = "$\leq$ one part in 10^3."

TABLE II

Shifts to Asymmetries and M_W from New Physics
Results are generic

One-Loop Physics	$\delta A_{LR}^{\mu^+\mu^-}$	$\delta A_{FB}^{\mu^+\mu^-}$	δM_W (MeV)
Heavy quark pair			
a) Large splitting	0.02	0.01	300
b) Degenerate	−0.004	−0.002	−42
Heavy lepton pair			
a) Large splitting $m_\nu = 0$	0.012	0.006	300
b) Degenerate	−0.0013	−0.0006	−14
Heavy squark pair			
a) Large splitting	0.02	0.01	300
b) Degenerate	0	0	0
Heavy slepton pair			
a) Large splitting	0.012	0.006	300
b) Degenerate	0	0	0
Winos			
a) $m_{3/2} \ll 100$ GeV	0.005	0.0025	100
b) $m_{3/2} \gg 100$ GeV	<0.001	<0.001	<10
Technicolor			
$SU_8 \times SU_8$	−0.04	−0.018	−500
O_{16}	−0.07	−0.032	−500
Strong Interaction Uncertainty	±.002	±.002	±25 MeV

FIGURE CAPTIONS

All figures from Ref. 8 unless otherwise noted.

1. Electroweak neutral currents. (a) Photon; (b) Z^0 (Ref. 7).
2. Z^0 resonance peak. $M_Z = 94$ GeV.
3. Final-state flavor dependence of A_{LR}. $M_Z = 94$ GeV.
4. Radiative corrections. (a) Oblique; (b) Direct vertex; (c) Direct box; (d) Bremsstrahlung.
5. Hadronization of final-state quarks at the Z^0 (Ref. 6).
6. Effect of bremsstahlung on Z^0 resonance shape (unpolarized). $M_Z = 93$ GeV. Dotted: No radiation. Solid: With initial-state radiation.
7. Cancellation of initial-state bremsstrahlung from A_{LR}. Lynn and Stuart: $M_Z = 94$ GeV, no radiation. BREM5: same, one-photon initial-state radiation. EXPOSTAR: same, initial-state radiation to all orders.
8. Starred running functions, for $M_Z = 94$ GeV, $m_{top} = 60$ GeV, $m_{Higgs} = 100$ GeV. (a) $s_*^2(s)$; (b) $e_*^2(s)$; (c) $G_{\mu_*}(s)$; (d) $\rho_*(s)$.
9. Comparison of errors in $s_*^2(Z)$ as measured by $A_{LR}(Z)$ and $A_{FB}(Z)$ as a function of number of Z^0 events, with errors from low-energy measurements.
10. $A_{LR}(Z)$ as a function of m_{top} and m_{Higgs}. $M_Z = 94$ GeV.
11. Predictions of grand unification models for $A_{LR}(Z)$ as a function of the unification scale M_X (Ref. 9).
12. Sum rule value (Eq. 4.8) for charm (solid) and bottom (dotted) final states as a function of the Z' mass. (a) $SU(2)_L \times U(1) \times U(1)'$; (b) $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ (Ref. 4).
13. $A_{LR}(Z)$ versus M_W for different values of m_{top} and m_{Higgs}. $M_Z = 94$ GeV.

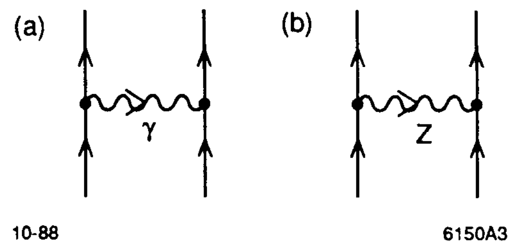
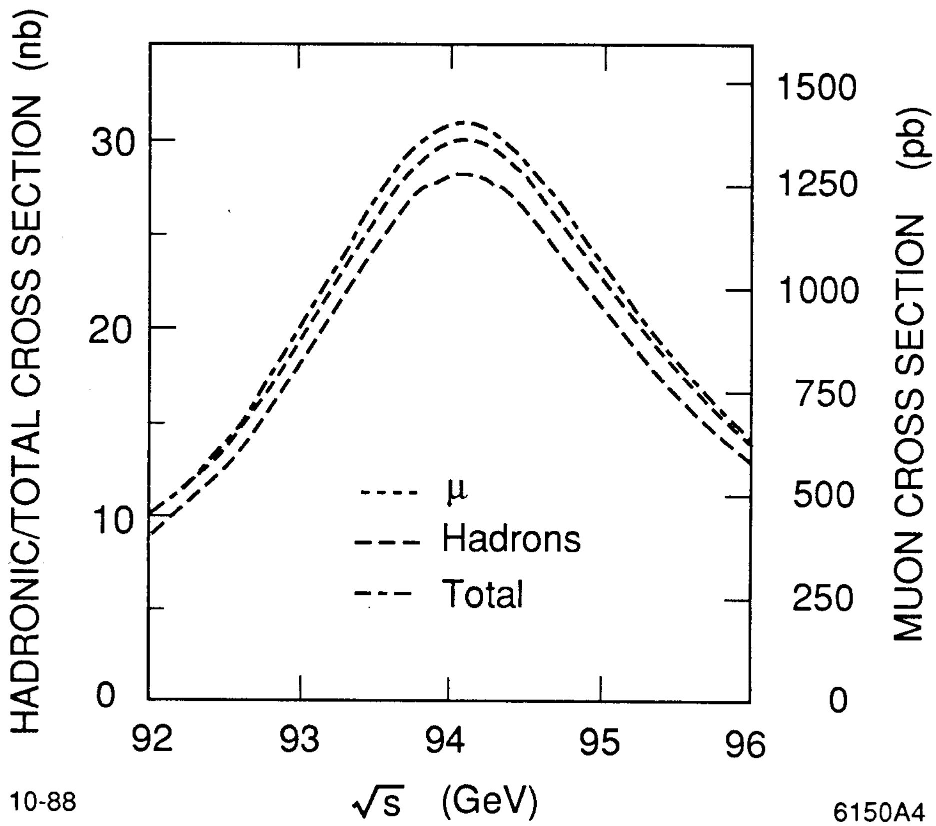


Fig. 1



HADRONIC/TOTAL CROSS SECTION (nb)
30
20
10
0
1500
1250
1000
750
500
250
0
MUON CROSS SECTION (pb)
μ
Hadrons
Total
92
93
94
95
96
10-88
√s (GeV)
6150A4

Fig. 2

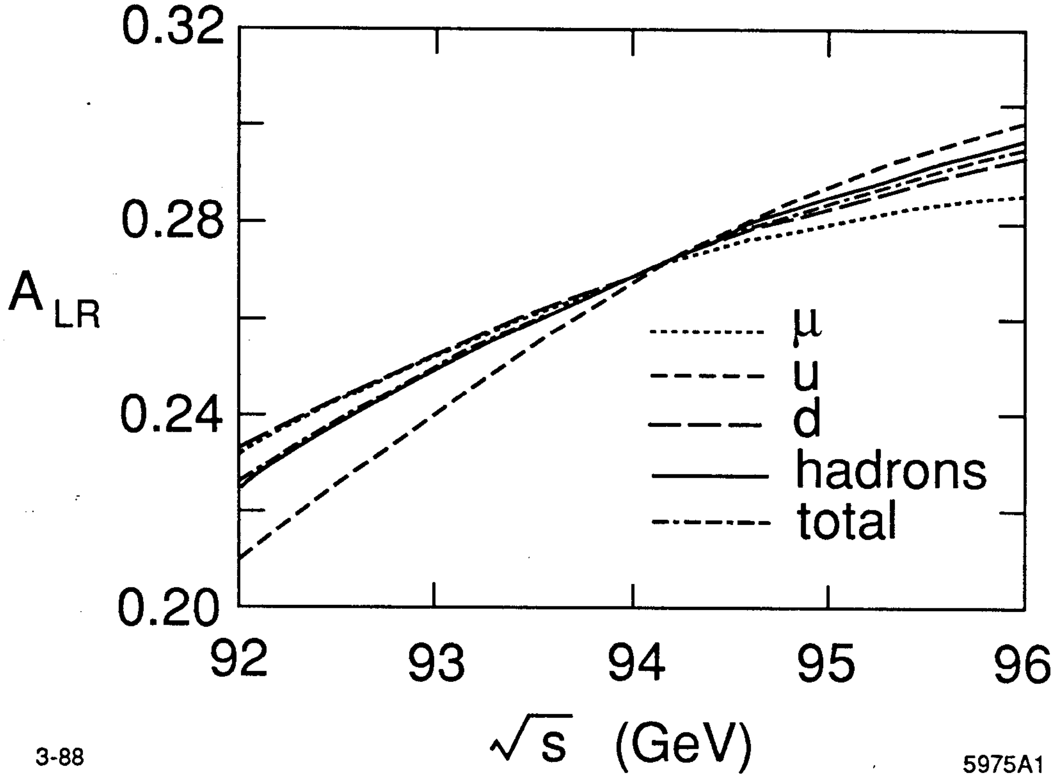

0.32
0.28
0.24
0.20
A LR
μ
u
d
hadrons
total
92
93
94
95
96
√s (GeV)
3-88
5975A1

Fig. 3

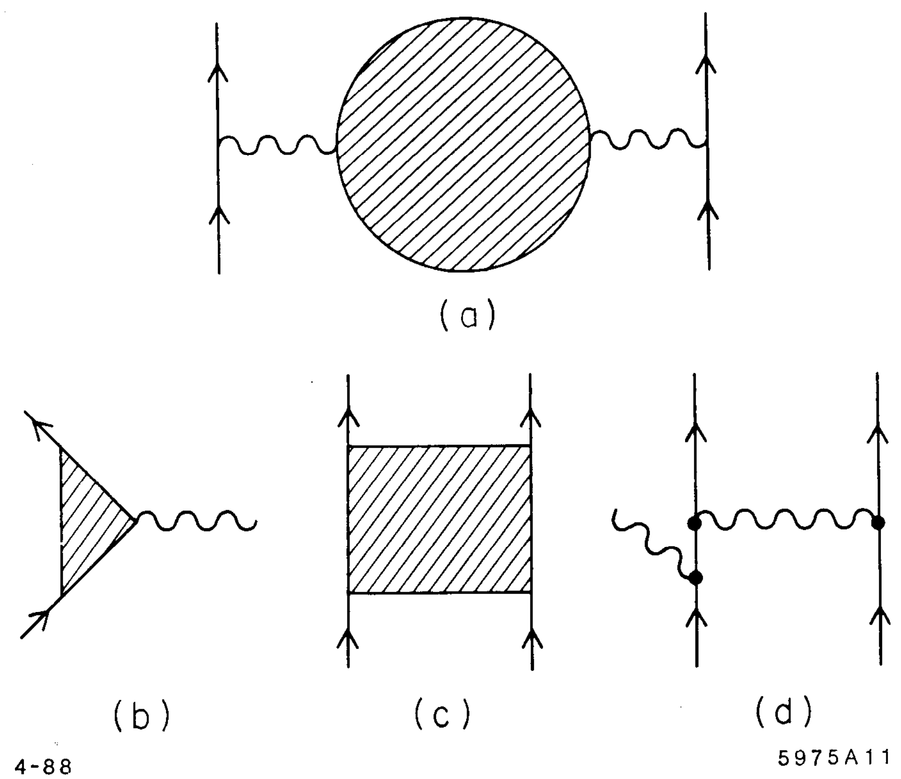


Fig. 4

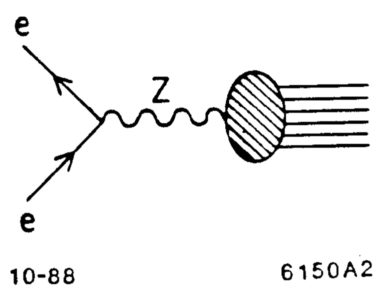


Fig. 5

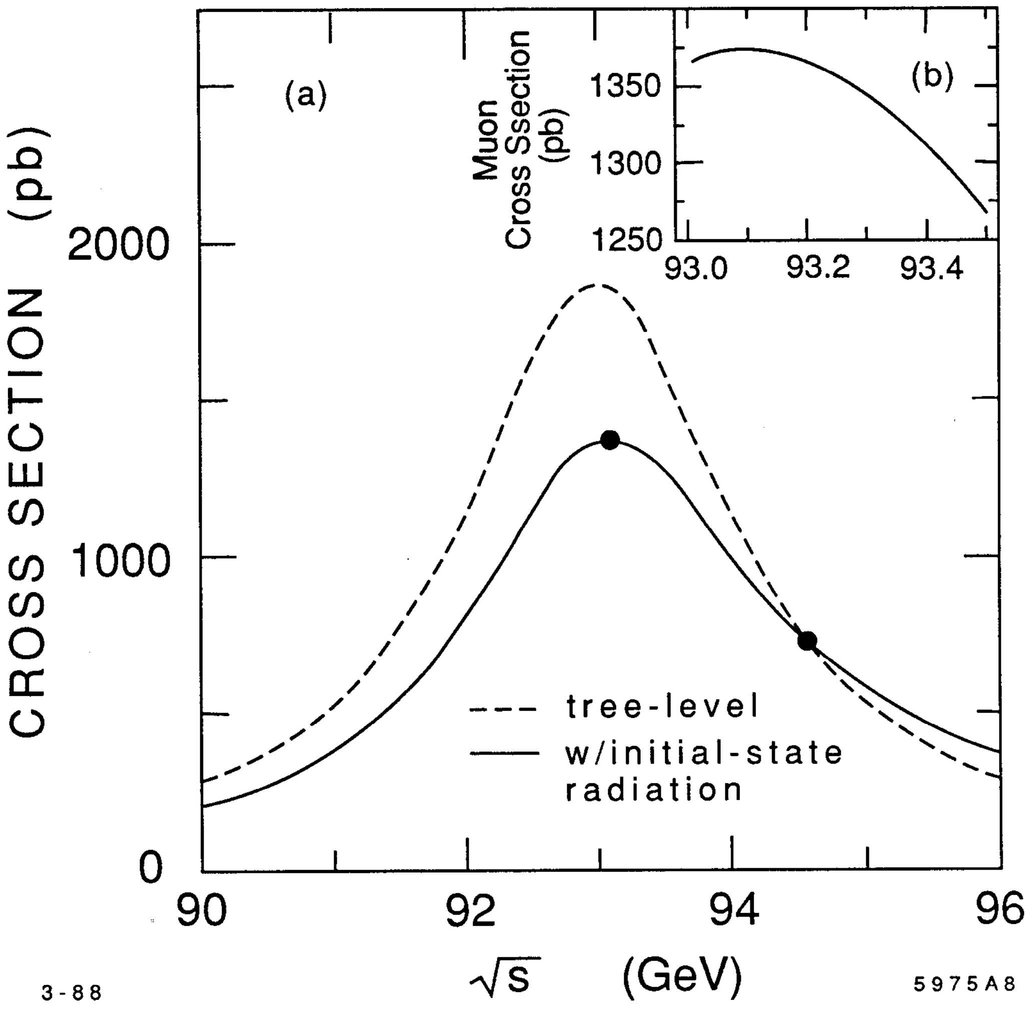
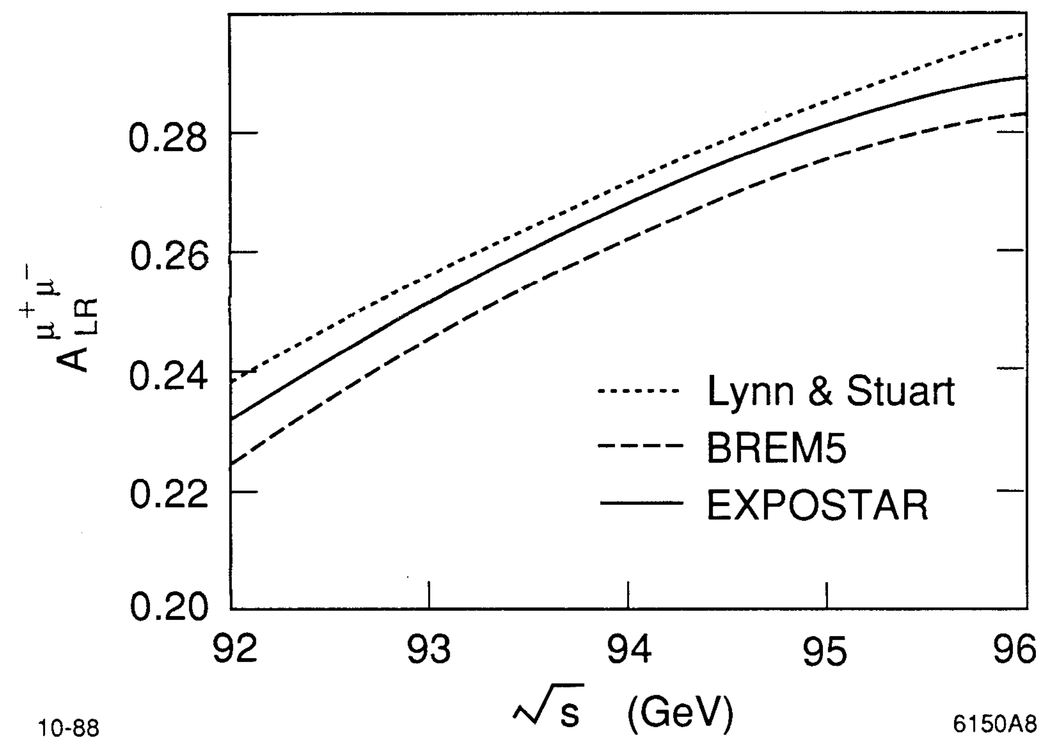
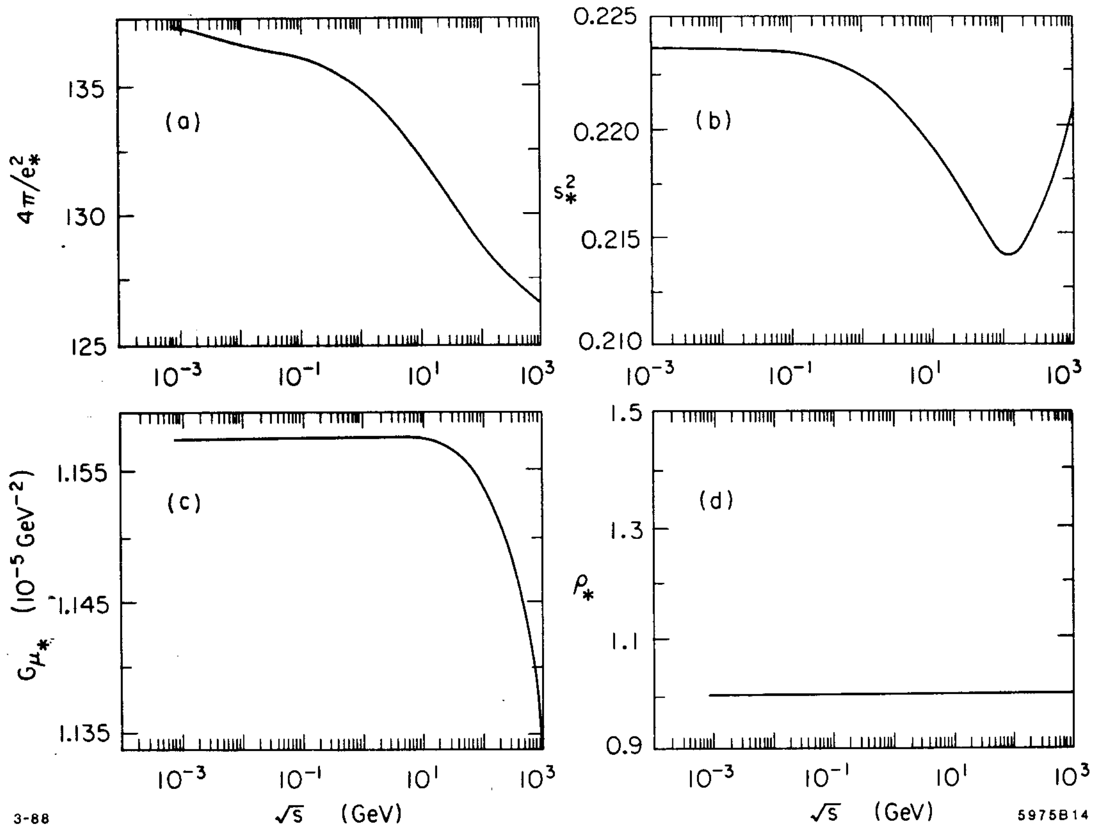

(a)
(b)
Muon
Cross Ssection
(pb)
1350
1300
1250
93.0
93.2
93.4
CROSS SECTION (pb)
2000
1000
0
90
92
94
96
tree-level
w/initial-state
radiation
√s (GeV)
3-88
5975A8

Fig. 6



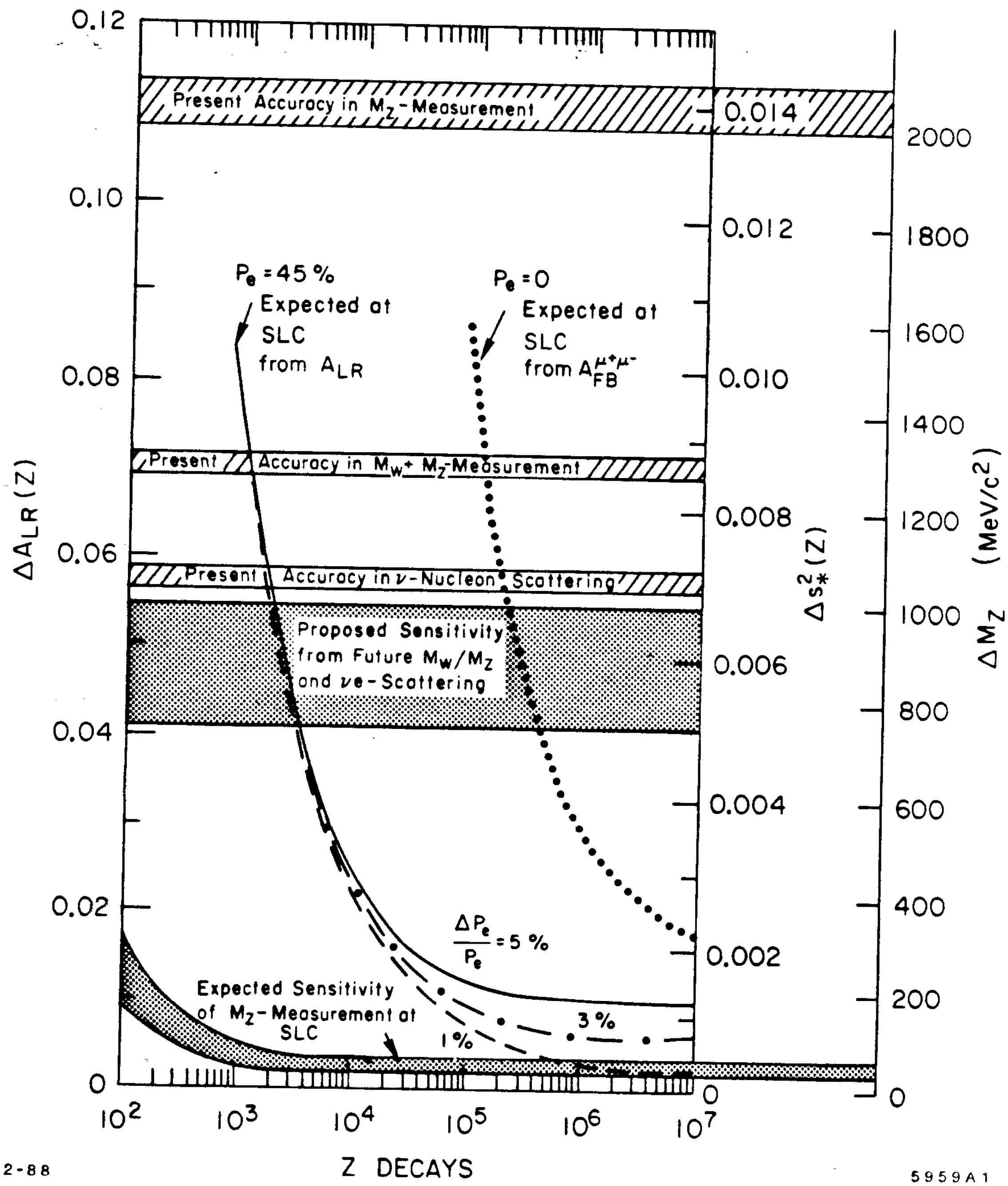
0.28
0.26
0.24
0.22
0.20
$A_{LR}^{\mu^+\mu^-}$
Lynn & Stuart
BREM5
EXPOSTAR
92
93
94
95
96
$\sqrt{s}$ (GeV)
10-88
6150A8

Fig. 7



(a)
(b)
(c)
(d)
135
130
125
0.225
0.220
0.215
0.210
1.155
1.145
1.135
1.5
1.3
1.1
0.9
$4\pi/e_*^2$
s_*^2
G_{μ_*} (10^{-5} GeV^{-2})
ρ_*
10^{-3}
10^{-1}
10^{1}
10^{3}
$\sqrt{s}$ (GeV)
3-88
5975B14

Fig. 8



0.12
0.10
0.08
0.06
0.04
0.02
0
$\Delta A_{LR}(Z)$
Present Accuracy in M_Z-Measurement
0.014
2000
0.012
1800
$P_e = 45\%$
Expected at
SLC
from A_{LR}
$P_e = 0$
Expected at
SLC
from $A_{FB}^{\mu^+\mu^-}$
1600
0.010
1400
Present Accuracy in M_W+ M_Z-Measurement
0.008
1200
ΔM_Z (MeV/c^2)
Present Accuracy in ν-Nucleon Scattering
$\Delta s_*^2(Z)$
1000
Proposed Sensitivity from Future M_W/M_Z and νe-Scattering
0.006
800
0.004
600
400
$\frac{\Delta P_e}{P_e} = 5\%$
0.002
Expected Sensitivity of M_Z-Measurement at SLC
200
3%
1%
0
0
10^2
10^3
10^4
10^5
10^6
10^7
Z DECAYS
2-88
5959A1

Fig. 9

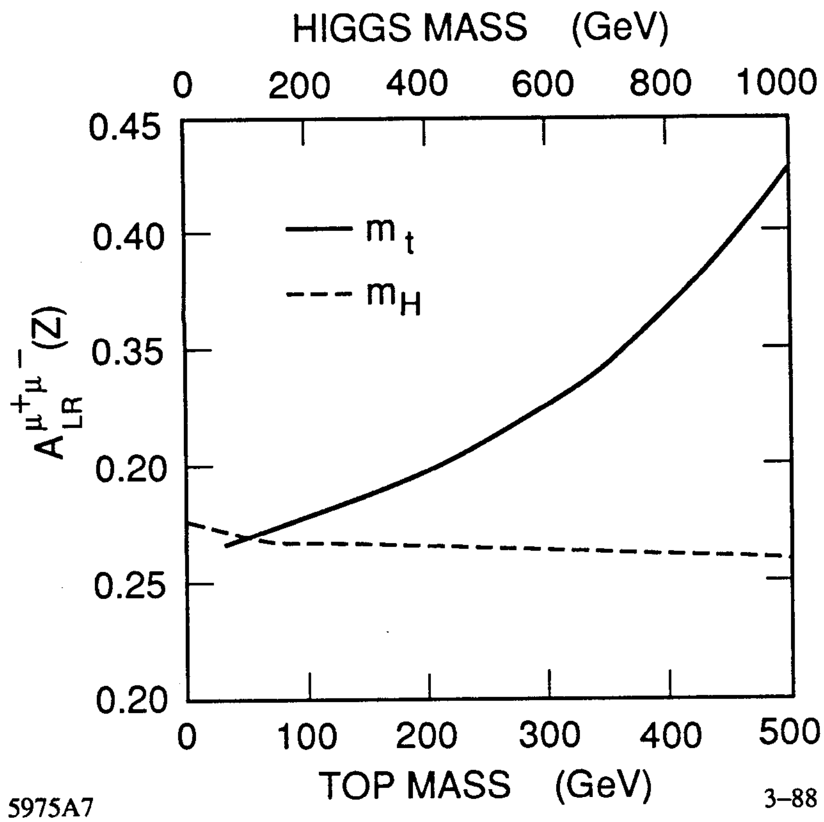
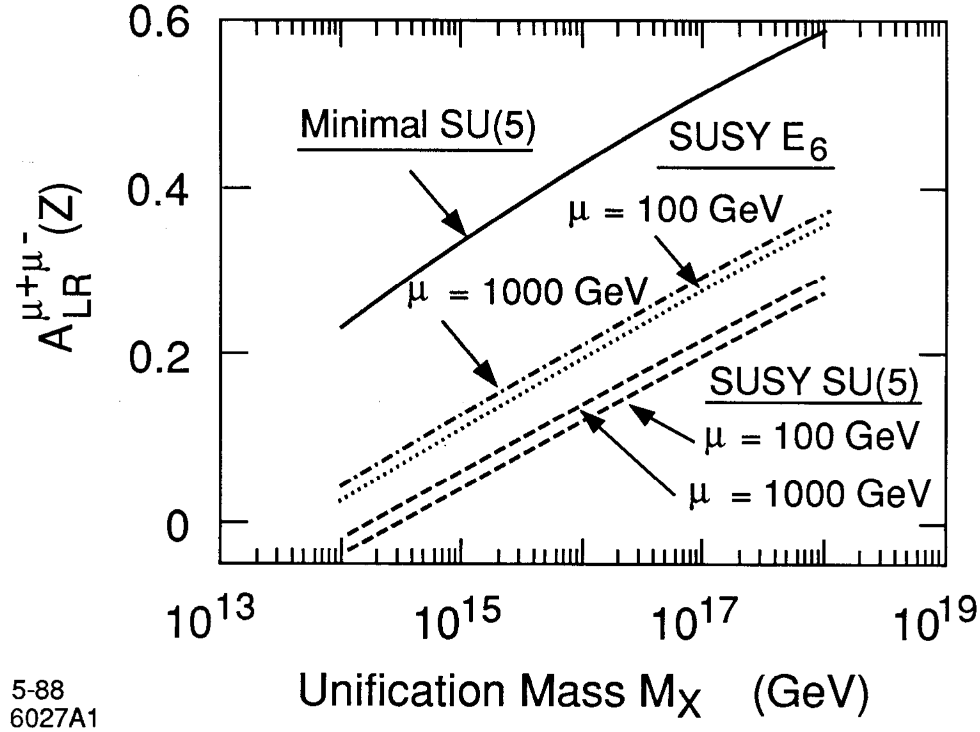
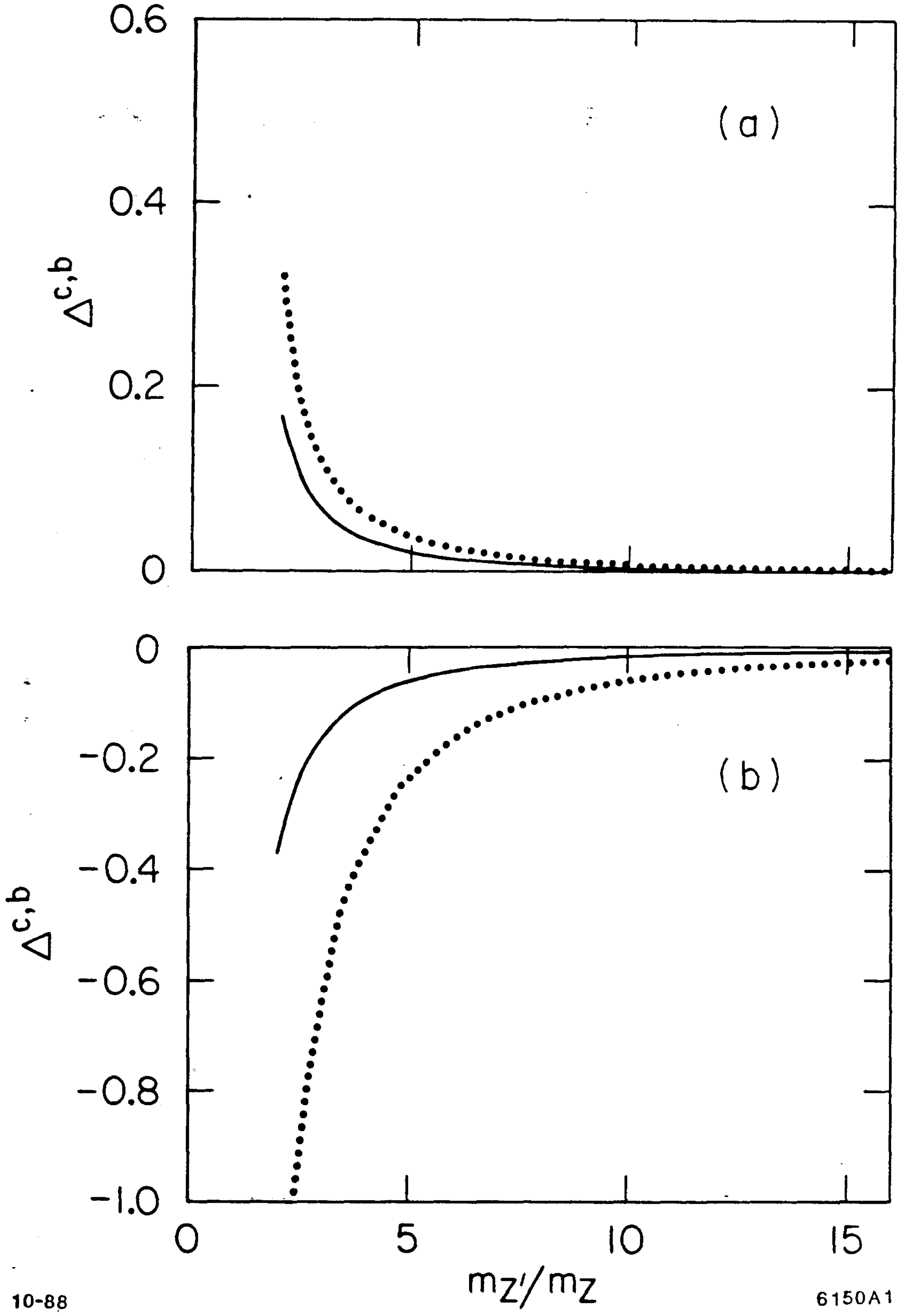

HIGGS MASS (GeV)
0
200
400
600
800
1000
0.45
0.40
0.35
0.20
0.25
0.20
m_t
m_H
$A_{LR}^{\mu^+\mu^-}(Z)$
0
100
200
300
400
500
TOP MASS (GeV)
5975A7
3–88

Fig. 10



0.6
0.4
0.2
0
Minimal SU(5)
SUSY E6
μ = 100 GeV
μ = 1000 GeV
SUSY SU(5)
μ = 100 GeV
μ = 1000 GeV
$A_{LR}^{\mu^+\mu^-}(Z)$
10^13
10^15
10^17
10^19
Unification Mass M_X (GeV)
5-88
6027A1

Fig. 11



0.6
0.4
0.2
0
(a)
$\Delta^{c,b}$
0
−0.2
−0.4
−0.6
−0.8
−1.0
(b)
$\Delta^{c,b}$
0
5
10
15
$m_{Z'}/m_Z$
10-88
6150A1

Fig. 12

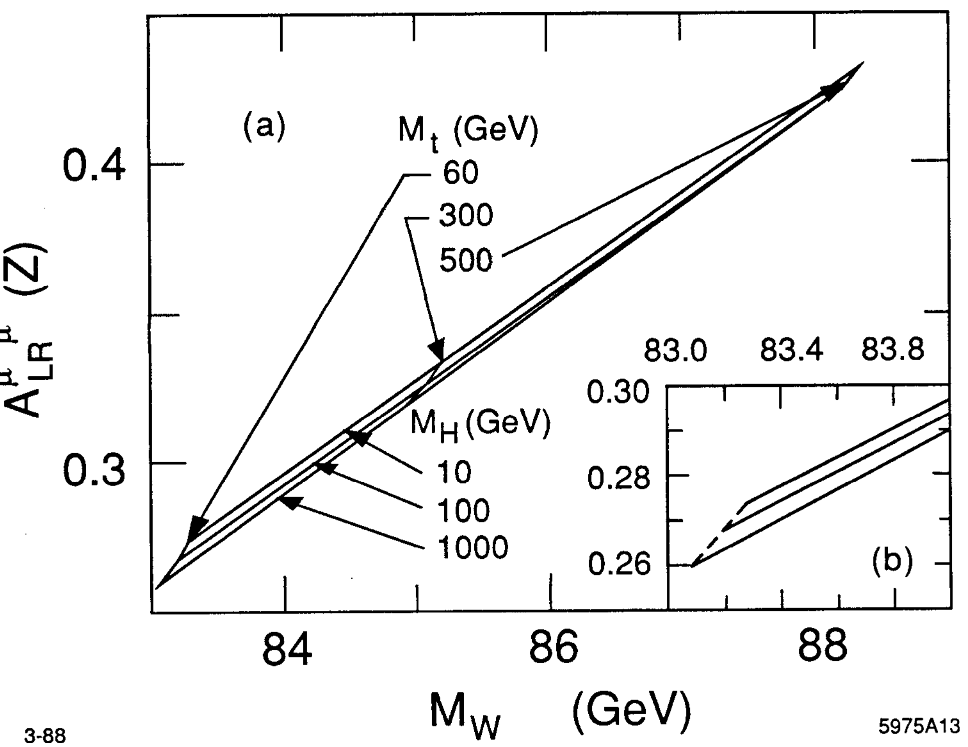

(a)
M_t (GeV)
60
300
500
M_H(GeV)
10
100
1000
0.4
0.3
$A_{LR}^{\mu^+\mu^-}(Z)$
84
86
88
M_W (GeV)
83.0
83.4
83.8
0.30
0.28
0.26
(b)
3-88
5975A13

Fig. 13