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Recent Results from PEP and PETRA

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Abstract

We discuss recent results from PEP and PETRA with emphasis on compositeness limits, gluon fragmentation, the determination of the Standard Model parameters α_s, M_Z, and $\sin^2\theta_W$ from the total hadronic cross section, estimates for the top mass within the SM, the renormalization scale dependence of various α_s determinations, hadronic jet multiplicities as compared to the second order QCD matrix element, and new results on the 'tau-decay' problem.

1 Introduction

Several of the PEP and PETRA experiments are still very active and summarizing all results from last year (around 50 publications!) on fragmentation[1], intermittency and correlations[2], lepton asymmetries[3], quark asymmetries[4], lifetime measurements[5], new particle searches[6], QED tests[7], and two photon physics[8] is far outside the scope of the present talk. Instead I will concentrate on a few selected topics, namely limits on compositeness scales, gluon fragmentation, estimates for the top mass, choice of scale for the strong coupling constant, R measurements and tau decays.

2 Compositeness limits from Bhabha scattering

In compositeness models the fundamental fermions have substructure. Bhabha scattering is particularly simple, since initial and final state particles are the same and no assumptions on the constituents of different fermions have to be made. A general parametrization of the interaction at the subconstituent level can be formulated by adding to the SM Lagrangian a contact term of the form[9]:

$$L_{\text{eff}} = \frac{g^2}{2\Lambda^2}\left(\eta_{\text{LL}} j_{\text{L}} j_{\text{L}} + \eta_{\text{RR}} j_{\text{R}} j_{\text{R}} + 2\eta_{\text{RL}} j_{\text{R}} j_{\text{L}}\right), \qquad (1)$$

where the contact interaction is assumed to be flavour-diagonal and helicity-conserving. The parameter Λ characterizes the mass scale of compositeness subject to the definition $g^2/4\pi = 1$ (strong coupling). As usual, j_{R} and j_{L} denote right-handed and left-handed currents. The interference between this contact interaction and the γ and Z exchange in the standard theory is responsible for terms proportional to the η's. Since the data is in good agreement with the SM, which presupposes pointlike particles without contact interaction, one can only give upper limits on the compositeness scale. Various kinds of helicity structure for the currents can be considered: LL coupling: $\eta_{\text{LL}} = \pm1, \eta_{\text{RR}} = \eta_{\text{RL}} = 0$; RR coupling: $\eta_{\text{RR}} = \pm1, \eta_{\text{LL}} = \eta_{\text{RL}} = 0$; VV coupling: $\eta_{\text{LL}} = \eta_{\text{RR}} = \eta_{\text{RL}} = \pm1$ and AA coupling: $\eta_{\text{RR}} = \eta_{\text{LL}} = -\eta_{\text{RL}} = \pm1$. The RR and LL couplings cannot be distinguished at this energy because they differ only by the negligible interference terms with the weak effect.

CELLO[10] has repeated the usual analysis of determining the limits on Λ from the Bhabha differential cross section. One of the difficulties is that the normalization error causes a correlation between the data points, which have not always been taken into account correctly. Therefore, we have repeated the analysis for the combined data at $\sqrt{s}$=29 and 35 GeV[11] by taking all correlations into account via a complete error correlation matrix[12]. The 95% c.l. lower limits on the Λ parameter from such a fit are listed in Table 1; the

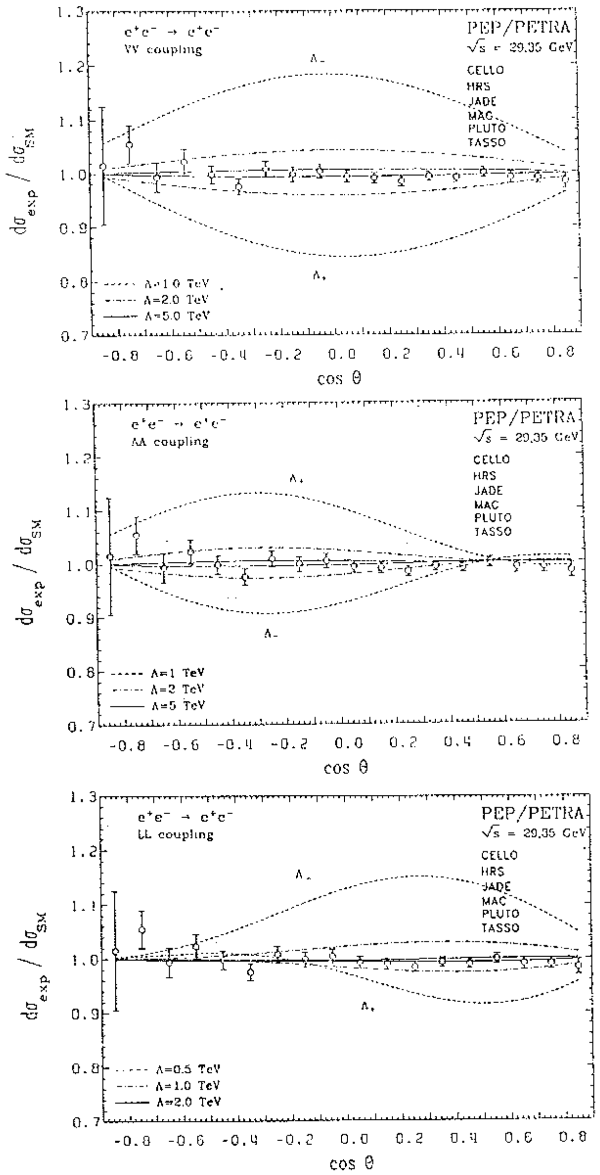


Figure 1: **The differential Bhabha cross section normalized to the standard model expectation with M_Z= 91.2 GeV and $\sin^2\theta_W = 0.23$. The curves show the deviations from compositeness interactions for different couplings: a) left handed or right handed coupling, b) vector coupling and c) axial vector coupling.**

current	η_{LL}	η_{RR}	η_{RL}	Λ^+ (TeV)	Λ^- (TeV)
VV	1	1	1	5.2	7.9
AA	1	1	−1	5.6	3.8
LL	1	0	0	2.0	3.0
RR	0	1	0	2.0	3.0

Table 1: Lower limits (95% c.l.) on the mass scale parameter Λ in composite models. The η factors are for Λ^+. They change sign for Λ^-.

deviations from the SM for various contact interactions are shown in Fig. 1. It should be noted that the limits from the combined data are not much better than some of the published limits. However, this is caused by not correctly treating the correlations in that case. If we take the correlations into account for data from such a single experiment, the limits are sometimes reduced by a factor close to two.

Traditionally any departure from QED has been parametrized by inserting time-like and space-like form factors at the respective vertices with cut-off parameters $\Lambda_\pm$(QED) [13] which are related to the Λ_{VV} by the following relation: $\Lambda(\mathrm{QED}) = \sqrt{\alpha} \cdot \Lambda_{VV}$. From the values in Table 1 we get as lower limits Λ_+(QED) > 680 GeV and Λ_-(QED) > 440 GeV. A cut-off parameter of 440 GeV verifies the validity of the SM down to distances of $4.5 \cdot 10^{-17}$ cm. This implies a limit of $1.1 \cdot 10^{-16}$ cm on the charge radius of the electron (using the relation $\Lambda^2 = 6/ < R_e^2 >$).

3 Gluon Fragmentation

Since the discovery of the 3-jet events, studies have been made to find a difference between quark and gluon fragmentation. QCD predicts larger multiplicities and consequently softer fragmentation. However, experimentally the results are inconclusive at present energies. The HRS Collaboration[14] quoted $1.25^{+0.17}_{-0.32} \pm 0.20$ for the ratio of gluon jet multiplicity to quark jet multiplicity. The JADE Collaboration[15] found that the least energetic jets in 3-jet events have the largest transverse momentum ($< p_{t3} > / < p_{t2} > = 1.16 \pm 0.02$), although preliminary results from the TPC[16] and CELLO[17] Collaborations are less striking. The AMY Collaboration[18] finds that gluon jets are less tightly collimated than quark jets. One difficulty with these studies is the definition of a gluon jet. E.g. in the string fragmentation model all partons are connected via strings and one cannot talk about a gluon jet fragmenting independently. This one can only do in independent fragmentation models, but these models are known not to describe all aspects from 3-jet events. String effects are only negligible at asymptotic energies. A nice way to circumvent the problem of having to define a gluon jet was introduced by the MARK-II Collaboration[19]. They compared the momentum spectra from symmetric 3-jet events at 29 GeV with 2-jet events at 19 GeV. In this case all jets have roughly equal energies and one knows the 3-jet events have one gluon jets, without having to specify which one is the gluon jet. One then compares the momentum spectra from these 2 and 3-jet events and the conclusion from the Mark-II Collaboration was that gluons do fragment softer than quark jets. This analysis has been repeated recently by the

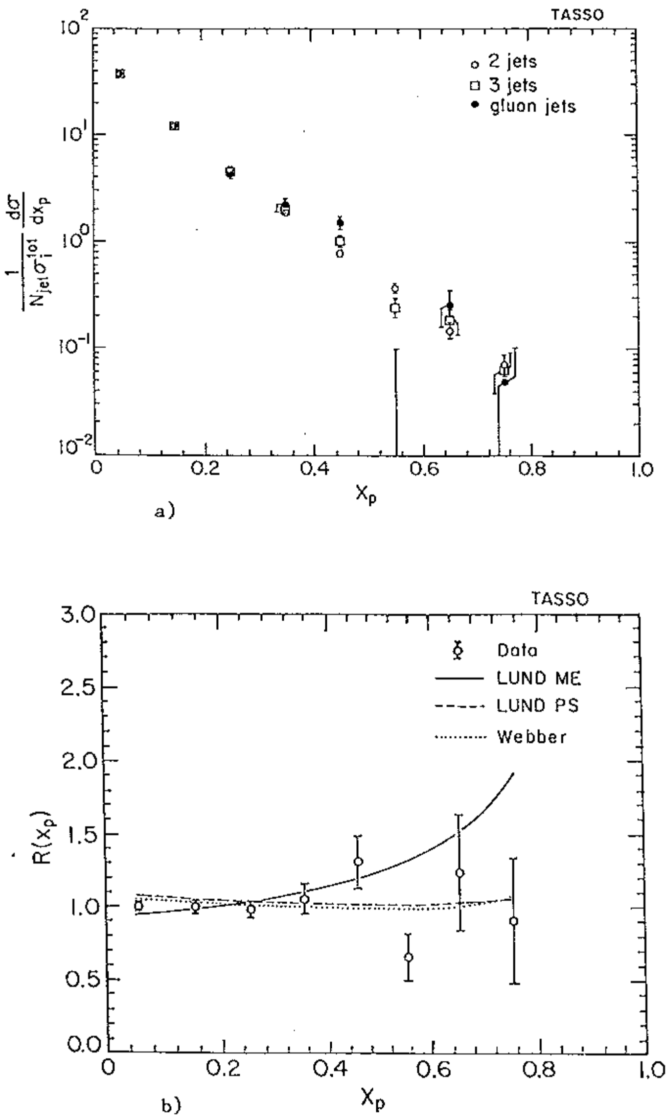


Figure 2: **a) The momentum spectra of 2- and mercedes-type 3-jet events and b) their ratio. The center of mass energies of the events are such that the jet energies in the 2- and 3-jet events are the same (about 11 GeV per jet). The momentum spectra for gluon jets have been obtained by subtracting the spectra of the 3- and 2-jet events.**

M_H (GeV/c^2)	M_{top} (GeV/c^2)
10	100 ± 35
100	115 ± 35
500	134 ± 33
1000	143 ± 32

Table 2: **M_{top} values for various M_H values from the Sirlin relation with M_Z=91.11±0.23 GeV/c^{2}and $\sin^2\theta_W$ =0.2293±0.003±0.0028. The errors on M_{top} stem from the uncertainty in $\sin^2\theta_W$. The uncertainty from M_Z introduces an additional error of 15 GeV/c^2.**

TASSO Collaboration[20] and their conclusion is opposite: they find that at PETRA energies there is no difference between quark and gluon fragmentation (see Fig. 2). They attribute the discrepancy with the MARK-II analysis to the fact that in that analysis the charged multiplicity cut for the 2-jet and 3-jet sample has not been identical.

4 Estimating the top mass in the minimal SM

The radiative corrections at the highest PETRA and TRISTAN energies depend on the unknown top and Higgs mass. From the recent precise M_Z determinations, combined with present knowledge about $\sin^2\theta_W$, one gets constraints on the top mass through the Sirlin relation[21]:

$$\sin^2\theta_W = \frac{1}{2}\left[1 - \sqrt{1 - \frac{4\pi\alpha}{\sqrt{2}G_F M_Z^2(1-\Delta r)}}\right] \qquad (2)$$

The radiative correction factor Δr to the Fermi constant G_F, as determined from muon decay, depends on the top and Higgs mass, so if M_Z and $\sin^2\theta_W$ are known, one can determine M_{top} for an assumed value of M_H, of course under the assumptions of the minimal SM. The most precise value of $\sin^2\theta_W$ stems from the ratio of neutral and charged current cross sections in deep inelastic neutrino quark scattering. This ratio has been measured to a 1% accuracy and the extracted value of $\sin^2\theta_W$ from such a ratio is insensitive to M_{top}[22,23,24]. For large top masses and assuming $\rho = 1$ one finds: $\sin^2\theta_W = 0.2293 \pm 0.0030 \pm 0.0028$[23]. The corresponding values of M_{top} are shown in Table 2 for various Higgs masses. The values range between 100 and 140 GeV, so taking this range into account in the errors, one estimates:

$$\mathrm{M_{top}} = \mathbf{120 \pm 35 \pm 20}\ \mathrm{GeV}.$$

The first error corresponds to the error in $\sin^2\theta_W$, while the latter error covers the uncertainty from M_Z and M_H. The first error can be reduced slightly to 30 GeV, if one includes all measurements on $\sin^2\theta_W$, taking into account the top dependence of $\sin^2\theta_W$ [24]. For these fits I have used the programs from Burgers and Hollik[25] to calculate Δr. The Higgs dependence found with these programs is appreciably larger than the one given by Ellis and Fogli[24].

5 Scale dependence of α_s determinations

5.1 Need for renormalization and definition of α_s

The coupling constant is not constant, but varies with Q^2, both in QED and in QCD. However, in QED the coupling constant increases as function of Q^2, while in QCD the coupling constant decreases. The running of the coupling constants can be calculated from the loop diagrams of virtual fermions in the gauge boson propagator. These diagrams are divergent for large Q^2, which implies that one has to introduce two scales in the coupling: in addition to the physical Q^2 of the reaction one needs a 'cut-off' like scale to regulate the divergencies. For the QCD coupling constant one finds from the one loop calculations:

$$\alpha_s(Q^2) = \alpha_s(\mu^2)\left[1 + \frac{\alpha_s(\mu^2)}{4\pi}\left(11 - \frac{2n_f}{3}\right) ln\frac{Q^2}{\mu^2}\right]^{-1} \quad (3)$$

Note that α_s decreases with increasing Q^2 if $11 - 2n_f/3 > 0$ or $n_f < 16$, thus leading to asymptotic freedom at high energy, if the number of active quark flavours n_f is less than 16.

One can trade the renormalization scale μ for the 'confinement' scale Λ, i.e. $\alpha_s \to \infty$, if $Q^2 \to \Lambda^2$. In this case $\alpha_s(\mu^2)$ drops out of Eq. 3 and one gets the well known first order result:

$$\alpha_s(Q^2) = \frac{4\pi}{(11 - 2n_f/3) ln(Q^2/\Lambda^2)} \equiv \frac{1}{b\, ln(Q^2/\Lambda^2)} \quad (4)$$

If one includes the next higher order terms, one finds[26]:

$$\alpha_s(Q^2) = \frac{4\pi}{\beta_0 L}\left[1 - \frac{\beta_1}{\beta_0^2}\frac{lnL}{L} + \left(\frac{\beta_1}{\beta_0^2}\right)^2 \frac{1}{L^2}\left\{(lnL - \frac{1}{2})^2 + \frac{\beta_2\beta_0}{\beta_1^2} - \frac{5}{4}\right\}\right] \quad (5)$$

with[27]

$$\begin{aligned} L &= ln(Q^2/\Lambda^2_{\overline{MS}}) \\ \beta_0 &= 11 - \tfrac{2}{3}n_f \\ \beta_1 &= 2(51 - \tfrac{19}{3}n_f) \\ \beta_2 &= \tfrac{2857}{2} - \tfrac{5033}{18}n_f + \tfrac{325}{54}n_f^2 \,. \end{aligned}$$

Instead of using a series expansion in α_s, one can solve the renormalization group equation exactly[28]. However, the solution of α_s for a given Λ can only be solved numerically in that case. If one includes the third order $(1/L^2)$ term in Eq. 5, this expression is numerically very close to the exact solution. On the other hand, if one neglects this term, as is done e.g. in the particle data book, the Λ value for a given α_s value deviates typically by 7% from the exact solution. Therefore, we will use the relation between α_s and $\Lambda^{(5)}_{\overline{MS}}$ as given above.

In QED the higher order terms are small, because of the smallness of the coupling constant and different renormalization schemes (RS) give very similar results. For QCD the higher order corrections are not necessarily small and the choice of RS becomes more critical. Many schemes and scales have been discussed[29,28]. Before proceeding, one should define the distinction between renormalization conventions and scales: a convention implies everything needed for the definition of the coupling constant except the renormalization scale. Both, different conventions and different scales lead to different coupling constants and usually different higher order corrections. As soon as we get a different coupling, we call it a different RS. Well known conventions are MOM-schemes or minimal subtraction schemes[29,28]. The latter define a gauge independent coupling constant and are the easiest framework for QCD; these are among the reasons why they are most widely used.

A few well known particular choices of scale are:

- Fastest Apparent Convergence. In this case the scale is choosen such that all higher order terms are zero (Fastest Apparent Convergence)[30].
- Principle of Minimal Sensitivity. Stevenson[28] proposed to choose for each process a renormalization point such that the observable shows minimal sensitivity to the RS, i.e. $\partial\Re/\partial(RS)$=0 or $\partial\Re/\partial ln\mu = 0$. The idea behind it is that at all orders each observable is independent of the RS and by choosing the PMS scale, one forces the approximate result to be in a region where the dependence on the RS is minimal. Stevenson[28] gives a detailed prescription how this can be implemented in practical applications to any desired order of perturbation theory.
- BLM Choice. Brodsky, Lepage and Mackenzie (BLM)[31] choose the scale, which absorbs all vacuum polarization contributions into the coupling in analogy with the usual practice in QED.

The RS dependence is most easily studied by varying the argument of α_s. Varying this scale is equivalent to varying the QCD scale Λ, since α_s depends only on the ratio Q/Λ. The effect of changing the α_s-scale can be studied easily as follows. Suppose a variable is given in a certain renormalization scheme and for a given Q^2 scale to be:

$$R = r_1\alpha_s + r_2\alpha_s{}^2 + r_3\alpha_s{}^3 + O(\alpha_s{}^4) \quad (6)$$

If we choose the coupling at a different scale, we get:

$$R' = r'_1\alpha_s{}' + r'_2\alpha_s{}'^{\,2} + r'_3\alpha_s{}'^{\,3} + O(\alpha_s{}'^{\,4}) \quad (7)$$

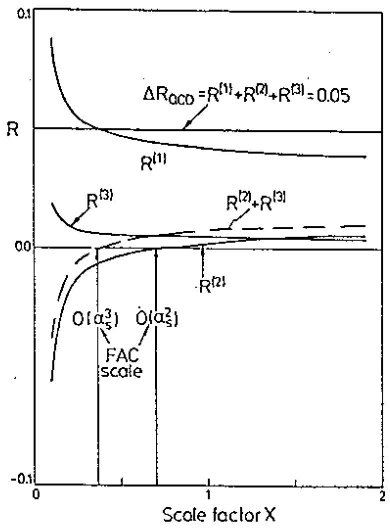
If one neglects the terms of order $O(\alpha_s{}^4)$, then

$$R' - R = dR = r_1 d\alpha_s + \alpha_s dr_1 + \alpha_s{}^2 dr_2 + 2\alpha_s r_2 d\alpha_s + \alpha_s{}^3 dr_3 = 0 \quad (8)$$

This can only be zero, if the coefficient for each power of α_s equals zero, which yields 3 equations for the 3 unknowns r'_1, r'_2, r'_3. After calculating $d\alpha_s$ from the renormalization group equation: $\mu\partial\alpha_s(\mu)/\partial\mu = -\beta'_0\alpha_s{}^2(\mu) - \beta'_1\alpha_s{}^3(\mu) + O(\alpha_s{}^4)$, we find at a different scale $Q' = xQ$:

$$\begin{aligned} \alpha_s{}' &= \alpha_s - \beta'_0\, ln\, x\, \alpha_s{}^2 - \beta'_1\, ln\, x\, \alpha_s{}^3 \\ r'_1 &= r_1 \\ r'_2 &= r_2 + r_1\beta'_0\, ln\, x \\ r'_3 &= r_3 + r_1\beta'_1\, ln\, x + 2r_2\beta'_0\, ln\, x \end{aligned} \quad (9)$$

The β-factors are renormalization scheme independent and given by $\beta'_0 = \beta_0/(2\pi)$ and $\beta'_1 = \beta_1/(4\pi^2)$. Here we have neglected terms of higher order in $ln\, x$ as well as the higher order terms in the renormalization group equation. These terms are only important for very small scales, in which case the scale dependence becomes so strong, that the approach of differential calculus is questionable anyway.



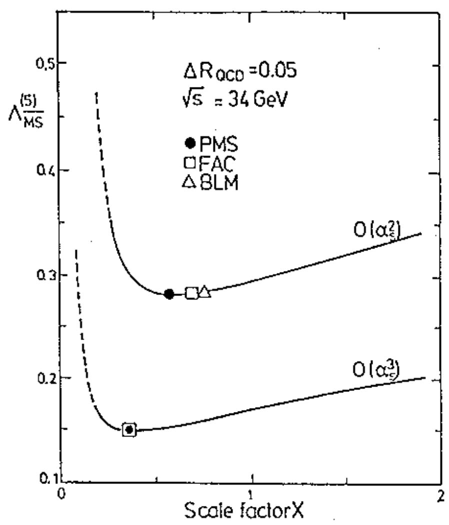


Figure 3: **Contributions to R (a) and Λ determinations in second - and third order (b) as a function of the renormalization scale factor x.**

5.2 Scale dependence of α_s from R

Studying the RS dependence of R is interesting for two reasons: a) it is the first process for which both the second and third order QCD contributions have been calculated, so one can check if specific choices of renormalization scales in second order would yield small contributions in third order, and b) the third order contributions are larger than the second order contribution for the usual scale $Q=\sqrt{s}$[32]. Since other choices of scale are equally well justified, one can see if there are scales for which the higher order contributions are better converging. We have studied these contributions as function of scale in contrast to previous results for a few specific scales[33]. The various QCD contributions to R as function of the renormalization factor x are shown in Fig. 3a, assuming the total contribution to be constant ($R_{QCD}=0.05$). One observes that at $x=1$ the third order contribution is indeed larger than the second order contribution, but at small and large x the absolute value of the second order contribution is larger. However, at all scales the first order contribution is dominant. We have indicated the scales at which the second order or the second plus third order contributions become zero. These are the FAC scales (Fastest Apparent Convergence)[30]. After recalculating the coefficients at a new scale (using Eq. 10), one can redetermine the corresponding α_s from the measured R-value and recalculate the corresponding Λ value. The result is shown in Fig. 3b. The minimum in this curve is the PMS scale corresponding to the point of minimal sensitivity[28]. One observes that at all scales the Λ values in third order are roughly a factor two below the Λ values in second order, indicating that for this reaction there is nothing like an optimum scale, where the higher orders are not important.

This clearly indicates that all the heated discussions about choosing a certain scale mean

nothing more than betting on the future: you only can say something seriously about higher order contributions by calculating them, not by fiddling with renormalization scales or schemes.

5.3 Scale dependence of α_s from jetmultiplicities

Studies on jet multiplicities[34,35,36] in hadronic events in e^+e^- annihilation revealed that $\mathcal{O}({\alpha_s}^2)$ QCD models underestimate the production rates of 4-jet events. Subsequently, it was shown that leading log QCD models can describe the jet multiplicities much better, especially if a combination of the first order matrix element and the leading log calculations is used, as implemented in the newer version of the LUND Monte Carlo[37]. At first this has been interpreted as a need for higher QCD calculations. However, one did not realize that the difference between the second order and leading log QCD calculations is not only the higher order diagrams with multiple jets, but the scale used as argument for α_s in this newer Monte Carlos is typically in the order of the gluon virtuality, i.e. in the order of a few GeV or even less. If one uses such a small scale for α_s, its value becomes correspondingly larger and especially the 4-jet multiplicity, being proportional to ${\alpha_s}^2$, will be enhanced[38]. Recently it has been shown, that second order QCD models can describe the jetmultiplicities extremely well, if the scale of α_s is choosen to be in the order of 1 GeV[35,36], provided one adjusts $\Lambda^{(5)}_{\overline{MS}}$ correspondingly as shown in Fig. 4 for the Mark-II data[36,39]. It should be noted that a scale $\mu^2=0.0017\ E^2_{cm}$ and $\Lambda^{(5)}_{\overline{MS}}$=95 MeV corresponds to α_s(1.2 GeV)=0.26. For such low scales and large values of α_s, the exact definition of α_s becomes important, as well as the specified number of active flavours. Furthermore, it has been pointed out[36] that a) both the 3-jet and 4-jet rate decrease with increasing $\sqrt{s}$, as expected in QCD from the running of α_s and b) the jetmultiplicities are badly described by abelian vector theories. This is encouraging, but before these effects are considered to be a proof for the running of α_s and the gluon self coupling in QCD, one has to make absolutely sure that fragmentation models cannot mimic such rather small effects.

5.4 Scale dependence of α_s from AEEC

The asymmetry of the energy weighted angular correlations (AEEC) is usually considered a good observable for the determination of the strong couling constant, since the contributions from 2-jet events are negligible and second order corrections are reasonably small. The results depend on the way the 3-jet events are fragmented (string fragmentation or independent fragmentation[40]). On the other hand, since the fragmentation contribution is small and negative, it can be used to set a lower limit on the value of α_s by ignoring completely the fragmentation contribution and comparing the experimental results directly with the expressions at the parton level. By studying the energy dependence from other variables with positive fragmentation contributions, one can obtain upper limits on α_s. Combining these limits yields rather stringent bounds on α_s values[40,41]. However, all these studies used $\sqrt{s}$ as argument for α_s. CELLO[42] has studied the renormalization scale dependence for the AEEC and finds that apart from the fragmentation dependence, a major uncertainty comes from the scale dependence, as shown in Fig. 5): if the scale of α_s is varied between $0.2\sqrt{s}$ and $\sqrt{s}$, $\Lambda^{(5)}_{\overline{MS}}$ varies a factor two in contrast to the example of R in the previous section, where $\Lambda^{(5)}_{\overline{MS}}$ varied only 10%, if the scale was changed in this range.

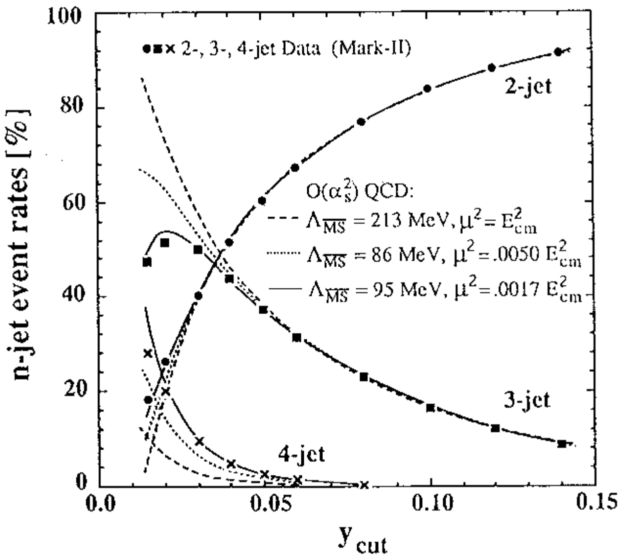


Figure 4: **Experimental data on jet multiplicities as function of the jet resolution parameter compared with different renormalization scales.**

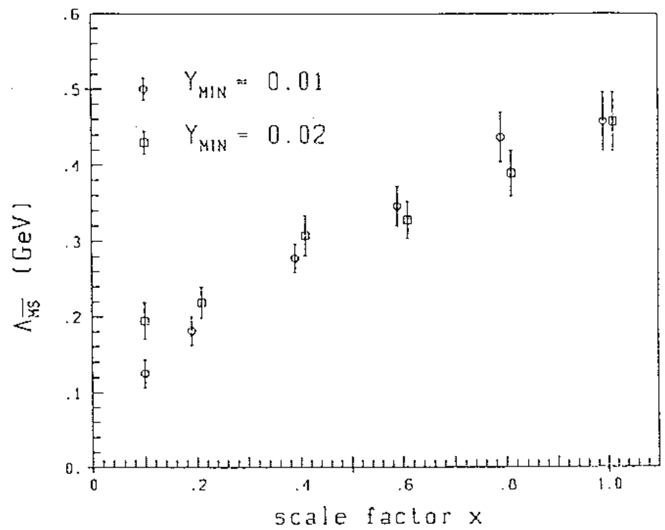


Figure 5: **Fitted $\Lambda^{(5)}_{\overline{MS}}$ values from the AEEC as function of the renormalization scale for two different values of the ymin parameter in the Monte Carlo.**

6 Results on R

6.1 Determination of M_Z

At the highest TRISTAN energies Z^0 exchange does increase R already by 50 %, thus allowing a direct measurement of M_Z. The radiative corrections depend on M_{top}. The M_Z values found for M_{top} between 60 and 180 GeV/c^2 range between 88.1 and 89.3 GeV/c^2[43,44,45]. Since the error on each value of M_Z is at least 1 GeV/c^2, there is no strong discrepancy (less than 3 s.d.) with the M_Z value of 91.11 GeV/c^2 from SLC[46]. Therefore, we will use this latter value to calculate the radiative corrections more precisely and use for M_{top} the value given in the previous section ($M_{top} = 120 \pm 35 \pm 20$ GeV/c^2).

The result from a 3 parameter fit to the data on R between 7 and 60.8 GeV is:

$$M_Z = 89.3 \pm 1.4 \text{ GeV}/c^2$$
$$\alpha_s(34\ GeV) = 0.142 \pm 0.018$$
$$\sin^2\theta_W = 0.217^{+0.024}_{-0.019}$$

The value of $\sin^2\theta_W$ agrees with the more precise value from deep inelastic neutrino scattering[23]: $\sin^2\theta_W = 0.2293 \pm 0.003 \pm 0.0028$. If we use this value for $\sin^2\theta_W$, we find:

$$M_Z = 88.9 \pm 1.2 \text{ GeV}/c^2$$
$$\alpha_s(34\ GeV) = 0.137 \pm 0.017$$

The uncertainty from the top mass entering into the radiative corrections is less than 0.1 GeV/c^2, if we use $M_{top} = 120 \pm 35 \pm 20$ GeV/c^2[45].

Here we included the data as presented at the Topical Conference at KEK[47]. These data are of a preliminary nature and the results should be considered accordingly.

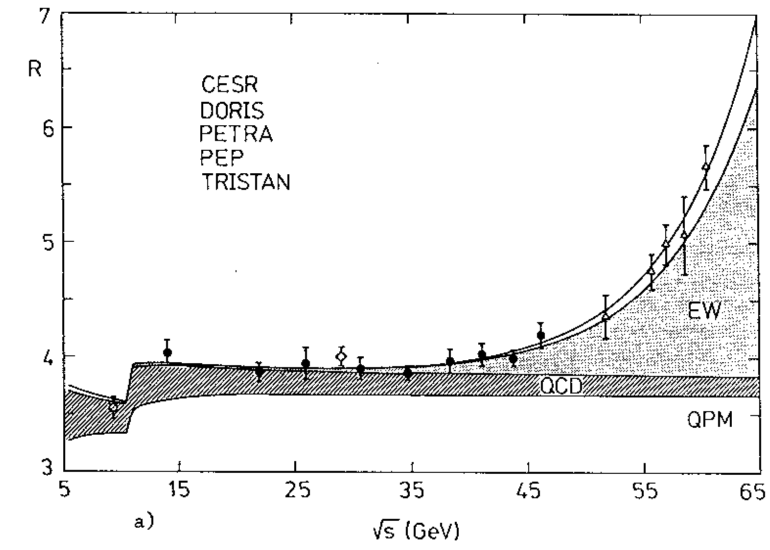
Fig. 6a shows the fit results together with the expectations from the SM with M_Z=88.9 and 91.11 GeV/c^2, Γ_{tot}=2.5 GeV/c^2, and $\sin^2\theta_W$ =0.2293. For clarity we have averaged the data points within certain energy bins. The error bars represent the total errors including systematic errors and the correlations.

One clearly sees that the TRISTAN data is above the fit expected for M_Z=91.11 GeV/c^2. The discrepancy is about 2 standard deviations, which either originates from a statistical fluctuation or a more interesting origin would be the contribution of new physics. One possible suggestion is the existence of a new heavy neutral gauge boson (Z')[48], which would be able to explain simultaneously a too high hadronic cross section and a too low leptonic cross section. The different behaviour of leptons and hadrons is caused by their different vector couplings (the interference between the photon and gauge bosons is only sensitive to the vector couplings). The μ-pair cross section from the combined TRISTAN data is indeed on the low side[49].

6.2 Determination of α_s

For the determination of α_s it is convenient to define the cross section with the G_F parametrization, which makes the results insensitive to M_Z and the unkown top mass[43]. From the fitted scale parameter for five active flavours $\Lambda^{(5)}_{\overline{MS}} = 0.260^{+0.160}_{-0.130}$ GeV one finds in third order:

$$\alpha_s(34\ GeV) = 0.144 \pm 0.015$$
$$\alpha_s(91\ GeV) = 0.122 \pm 0.011.$$



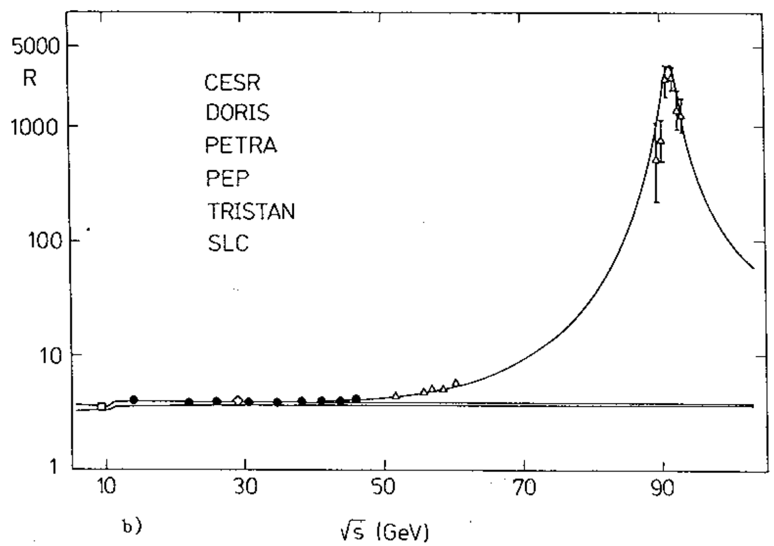


Figure 6: **Experimental data (fully corrected for electroweak contributions) and the Standard Model expectations for $M_Z = 91.11$ GeV/c^2, $\sin^2\theta_W = 0.2293$ and $\Gamma_{tot}=2.5$ GeV/c^2. The highest curve in a) uses the best fitted value with $M_Z = 88.9$. The two lowest curves indicate the contributions from the parton model and the QCD contribution, respectively.**

Here we have taken the third order QCD corrections into account, which lower the α_s values about 10% with respect to the second order one. It is interesting to determine the QCD contribution independently from the definition of α_s. If one assumes a linear energy dependence within the energy range, one finds this contribution to be[43]: $f_{QCD}(\sqrt{s} = 34\ GeV) = 1.057 \pm 0.008$. An extrapolation of this value to the LEP/SLC energy range yields[43] $f_{QCD}(\sqrt{s} = 91\ GeV) = 1.046 \pm 0.006$, which is an experimental number for the infinite series $(1 + \alpha_s/\pi + ...)$.

6.3 Comparison between PEP/PETRA/TRISTAN and SLC data on R

Fig. 6b shows all data on R for centre of mass energies between 7 and 93 GeV. In order to compare the 'low' energy data with the published SLC data, we subtracted the small leptonic 'background' contribution in the latter and applied the second order exponentiated radiative corrections using the program ZHADRO from G. Burgers and W. Hollik[25]. We used M_{top}=115 GeV/$c^2$and M_H=100 GeV/c^2 and assumed that the detection efficiency, calculated for first order radiative corrections, does not change in second order, which is certainly a good approximation with the present statistical errors.

The curve in Fig. 6b corresponds to $M_Z = 91.11$ GeV/c^2 and $\Gamma_{tot} = 2.5$ GeV/c^2. Here we have used the SM width (Γ_{tot}=2.5 GeV/c^2), which is close to the width calculated from the fitted number of neutrino species: $\Gamma_{tot} = 2.63 \pm 0.23$ GeV/c^2 corresponding to $N_\nu = 3.8 \pm 1.4$[46]. The directly fitted width ($1.61^{+0.60}_{-0.43}$ GeV/c^2[46]) is slightly lower, because of the different assumptions involved: in the first case one assumed all couplings known, while in the fit of Γ_{tot} one only assumed the lepton couplings to be known. In the latter case the couplings to hadrons are determined by: $\Gamma_f = \Gamma_h + f(\Gamma_\mu + \Gamma_\tau) = \Gamma_{tot} - (1-f)(\Gamma_\mu + \Gamma_\tau) - \Gamma_e - N_\nu\Gamma_\nu$ with f=0.556. Γ_{tot}=1.61 leads to $\Gamma_h \approx 0.8$ GeV/c^2, which has to be compared with 1.6 GeV expected for 5 quarks and 3 colours. However, neither value of Γ_{tot} (1.6 or 2.6) deviates by more than 1.5 s.d. from the expected SM value of 2.5 GeV/c^2, so there is no disagreement. One just has to wait for more statistics, especially to get better limits on the number of neutrino species.

7 New results on tau decays

7.1 Definition of the 'tau-decay' problem

The production properties of the τ are in good agreement with the expectations of the SM[50]. However, the decays are not completely understood: the results on two completely different kind of measurements, namely the fraction of inclusive decays into one charged particle plus anything is about $5 \pm 1.5\%$ larger than the sum of the exclusive, completely identified decay fractions with one charged particle. This discrepancy is called the 'tau-decay' problem[51,52].

We will proceed with shortly summarizing the results on the branching ratios, which can be found in several recent tau reviews[50], and discuss two new analysis by the CELLO-[53] and HRS Collaboration[54] in more detail.

The interest in the CELLO analysis stems from the fact that it is a high statistics analysis, in which both the topological and exclusive branching ratios have been measured

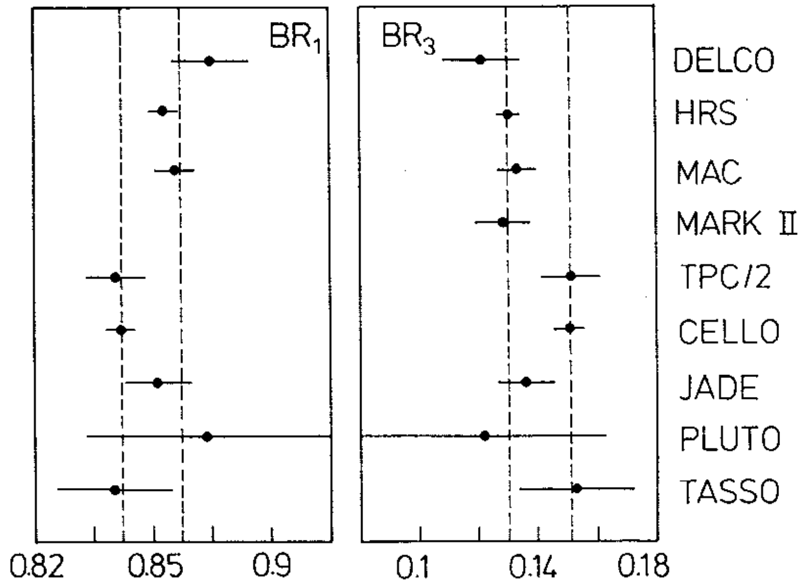


Figure 7: **Experimental results on the topological one prong branching ratios.**

simultaneously for the first time. Up to now, every experiment delivered only partial results on the decay modes. The disadvantage is clear: For a complete picture one has to combine several experiments, most of them being dominated by systematic errors, either particular for a given experiment or systematic errors, common to all experiments. Errors of the latter type are caused e.g. by the Monte Carlo input of the non-measured or badly measured tau decay modes. If all decay signatures are being measured simultaneously, one can iterate the dependence of the experimental results on the Monte Carlo input until a stable situation has been reached.

7.2 Results on topological branching ratios

The topological branching ratios are defined as

$$B_{1(3)} = \frac{\Gamma(\tau \rightarrow 1(3)\ charged\ particles + X)}{\Gamma_{tot}} \tag{10}$$

Measuring them is usually considered easy, because it is a partially inclusive measurement of charged particles only, in which unknown neutrals X do not have to be identified.

However, as shown already by the discrepancies between the various experiments (see Fig. 7), such measurements are not really that easy. For example, the two most precise (and new) measurements from CELLO ($B_1 = 0.849 \pm 0.004 \pm 0.003$) and HRS ($B_1 = 0.864 \pm 0.003 \pm 0.003$) differ several standard deviations[53,54].

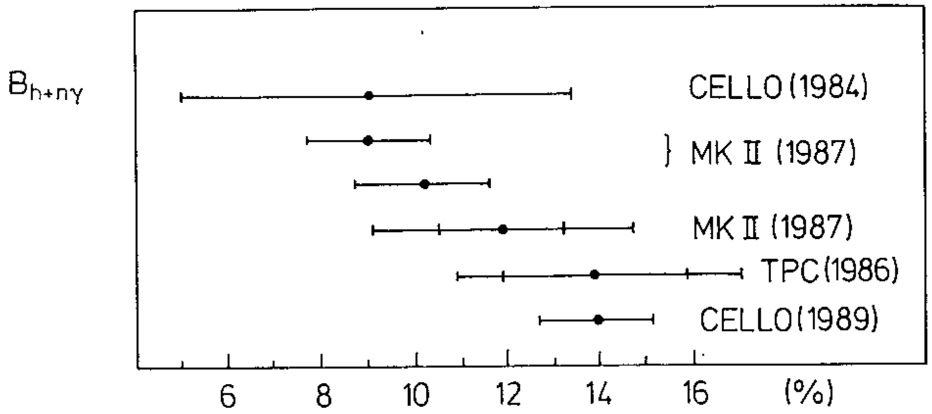


Figure 8: **Experimental results on the topological one prong branching ratios with multiple π^0's.**

One of the difficulties lies in the fact that one has to make rather hard cuts in the visible energy, acoplanarity and invariant mass to eliminate possible background from QED-type reactions and multihadronic final states. For example, the two prong (1-1) topology constitutes 72% of all tau decays. The efficiency for this topology is only 30% in the CELLO - and 11% in the HRS analysis. Obviously, one cuts rather strongly in various distributions and therefore depends on Monte Carlo simulations. Note too, that although the topological branching ratios are defined with charged particles only, one has to understand the electromagnetic calorimeters very well too, since they are used e.g. in the visible energy cuts. Understanding calorimeters is not always easy, as exemplified e.g. by the 'observation' of a large tau branching ratio into $\eta\pi\nu$ by the HRS Collaboration[55], which has been proposed originally as a solution to the tau-decay problem, but such large $\eta\pi\nu$ branching rations could not be reproduced by many other experiments[56].

7.3 Results on exclusive branching ratios

In this case one needs particle identification in the final state. For easily identifiable final states, like $\pi\nu_\tau$, $\rho\nu_\tau$, or $\mu\overline{\nu_\mu}\nu_\tau$ the experimental results are in excellent agreement with theoretical expectations, as shown in the last two columns of Table 3. The world average for the branching ratio into $e\overline{\nu_e}\nu_\tau = 17.7 \pm 0.4\%$, which is about 3 standard deviations below the theoretical expectation of 18.9%.

Indirect measurements of B_e are in better agreement with the theoretical expectation: a) from $e-\mu$ universality one expects $B_e = B_\mu/0.973 = 18.1 \pm 0.4\%$ and b) from the measured muon and tau lifetimes ($\tau_\tau = 3.027 \pm 0.078 \cdot 10^{-13}$) sec, where we have included the recent lifetime measurements by JADE[57] and TASSO[58]) one finds: $B_e = \frac{\tau_\tau}{\tau_\mu}\left(\frac{M_\tau}{M_\mu}\right)^5 = 18.9 \pm 0.5\%$.

We will now discuss the experimentally difficult decay modes with multiple π^0s. The data have been summarized in Fig. 8. One observes a range of values between 8 and 16%. The experimental problems stem from the overlap of hadrons (charged and neutrals) in a τ-jet, which is a really high density jet due to the high boost. This makes the efficiency matrix

Decay channel	Class	CELLO experiment	world av.	theor. expect.
$\tau \to e\nu\nu$	1	$18.4 \pm 0.8 \pm 0.4$	17.5 ± 0.4	18.9 ± 0.5
$\tau \to \mu\nu\nu$	2	$17.7 \pm 0.8 \pm 0.4$	17.8 ± 0.4	18.4
$\tau \to$ hadron ν	3	$12.3 \pm 0.9 \pm 0.5$		
$\tau \to \pi\nu$		$11.1 \pm 0.9 \pm 0.5$	10.8 ± 0.6	11.4
$\tau \to K\nu$			0.7 ± 0.2	0.74
$\tau \to$ hadron $+ 2\gamma\ \nu$	4	$22.6 \pm 1.5 \pm 0.7$		
$\tau \to \rho\nu$		$22.2 \pm 1.5 \pm 0.7$	22.3 ± 1.1	23.2
$\tau \to K\pi\nu$			1.4 ± 0.3	1.2
$\tau \to$ hadron $+ > 2\gamma\ \nu$	5	$14.0 \pm 1.2 \pm 0.6$		
$\tau \to \pi 2\pi^0\nu$		$10.0 \pm 1.5 \pm 1.1$	7.5 ± 0.9	9.0
$\tau \to \pi \geq 3\pi^0\nu$		$3.2 \pm 1.0 \pm 1.0$	3.0 ± 2.7	1.2
$\tau \to$ 3 hadrons ν	6	$9.0 \pm 0.7 \pm 0.3$		
$\tau \to 3\pi\nu$		$8.7 \pm 0.7 \pm 0.3$	6.8 ± 0.6	9.0
$\tau \to$ 3(5) hadrons $+ \geq 1(0)\gamma\ \nu$	7	$5.8 \pm 0.7 \pm 0.2$		
$\tau \to 3\pi \geq 1\pi^0\nu$		$5.6 \pm 0.7 \pm 0.3$	4.4 ± 1.6	5.0
$\tau \to 5\pi \geq 0\pi^0\nu$		0.16 ± 0.14	0.12 ± 0.03	
$\tau \to K\pi\pi\nu$			0.22 ± 0.14	0.2
$\tau \to K\bar{K}\nu$				0.6
$\tau \to K\bar{K}\pi\nu$			0.22 ± 0.15	0.2
$\tau \to 5\pi\nu$			0.06 ± 0.02	0.2
$\tau \to 6\pi\nu$			0.05 ± 0.02	0.2
Sum		$99.8 \pm 2.6 \pm 1.2$	92.8 ± 3.6	100.0

Table 3: **Measurements of the exclusive branching ratios in %. The branching ratios within the classes 3,4,5 and 6 have been determined by subtracting the expected strange particle contributions. For comparison the world averages are also given. The theoretical expectations are based on the standard model and are normalised to the measurement of the τ lifetime.**

non-diagonal and therefore dependent on the Monte Carlo input. The most precise value has been obtained recently by the CELLO Collaboration from PETRA's last year of high statistics data taking[53]. The fine granularity and correspondingly good spatial resolution of their liquid argon calorimeter allows them to separate $2\pi^0$, $3\pi^0$, and $n\pi^0$ final states. The CELLO results for all exclusive branching ratios are shown in Table 3 together with the previous world averages and theoretical expectations. Differences with the world average are in $Br(\tau \to e\nu\nu)$, $Br(\tau \to h+ \geq 4\gamma\nu)$, $Br(\tau \to \pi 2\pi^0\nu)$,and $Br(\tau \to 3\pi\nu)$. Especially the difference in the last one is important, since this one has been used by Gilman and Rhie[51] to calculate an upper limit on the poorly known topological branching ratio $Br(\tau \to \pi 2\pi^0\nu)$ by assuming $Br(\tau \to \pi 2\pi^0\nu) = Br(\tau \to 3\pi\nu)$ (isospin invariance). The new CELLO data is in agreement with this last relation ($8.8\% \pm 0.8\%$ versus $10.2 \pm 1.5 \pm 1\%$, respectively), but the values are larger than the world averages. From Table 3 one concludes that there is no 'tau-problem' within the CELLO data: the exclusive branching ratios add up to $99.8 \pm 2.6 \pm 1.2\%$, of which $84.5 \pm 2.3 \pm 1.2\%$ are one prong decays. This latter number agrees excellently with their topological one prong branching ratio of $84.9 \pm 0.4 \pm 0.3\%$. It is difficult to conclude from a single experiment, that the 'tau-problem' has disappeared. One better concludes with the hope that another experiment will try to repeat the effort of measuring both topological and exclusive branching ratios simultaneously, which is certainly preferable over the combination of bits and pieces from many different experiments, as discussed e.g. in detail in Ref.[59]. Especially , if one measures the decay modes only partially, one cannot iterate the observed branching ratios as Monte Carlo input and make consistency checks.

8 Summary

Seven from the nine PEP and PETRA experiments are still very active and many interesting results are still coming out (around 50 publications last year!). The SM is in excellent shape and hard to beat: Limits on compositeness scales are in the TeV range, solutions for the τ decay problem appear in sight, and the gluon fragmentation appears to be in good agreement with the very simple picture of string fragmentation.

The strong coupling constant can be determined from the R data almost independently from M_Z. One finds for five active flavours: $\Lambda^{(5)}_{\overline{MS}} = 0.260^{+0.160}_{-0.130}$ GeV, which corresponds to (using Eq. 5)

$$\alpha_s(34\ GeV) = 0.144 \pm 0.015$$
$$\alpha_s(91\ GeV) = 0.122 \pm 0.011.$$

Here we have taken the third order QCD corrections into account, which lower the α_s values about 10% with respect to the second order one. Studying the scale dependence (i.e. the argument of α_s) has become popular and useful, since a strong scale dependence indicates that one has to worry about the higher order corrections. On the other hand, if the dependence is weak, this does not guarantee that the higher order corrections are small, as exemplified for R, which is the only obervable sofar for which both the second and third order corrections have been calculated.

One of the surprises has been that by choosing an α_s scale in the order of 1-2 GeV, the second order QCD matrix element can describe the hadronic jet multiplicities extremely well. Such a small scale is typically used in parton shower Monte Carlos too, so the claim one needs the higher order leading log contributions in the matrix element is premature at PEP and PETRA energies, since the second order matrix element Monte Carlo seems to do equally well.

A question mark is coming from the difference between M_Z from 'low' and 'high' energy data: from the increase in R at the highest PETRA and TRISTAN energies one finds: $\mathbf{M_Z = 89.4 \pm 1.3}$ GeV/c^2 (using $\mathbf{M_{top} = 120 \pm 35 \mp 20}$ GeV/c^2, as obtained from a fit to M_Z and $\sin^2\theta_W$ from deep inelastic scattering), which is about 2 standard deviations below the direct measurement from the Z^0 resonance, at least if one considers the hadronic data from PETRA and TRISTAN. If one includes both hadronic and leptonic data, the value becomes 90.4 GeV. This implies that the hadronic data is above the expectation of the SM, while the leptonic data is below, which is the possible signature for a second neutral heavy gauge boson[48]. Clearly the effect is statistically not yet compelling, but it will be interesting to watch future data from TRISTAN at somewhat higher energies.

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