Resolution of a long standing discrepancy in R with spin zero quarks

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Abstract

A previously successful dispersive method has been applied to understand different values for $R(\sqrt{s} = 5 \div 7.5 \, \text{GeV})$ obtained by MARK I and Crystal Ball collaborations. We compute $R$ in the disputed region with data from outside this region and asymptotic behavior given by the standard model with 5 quark flavors, and find agreement with the Crystal Ball result. On the other hand, the MARK I data are reproduced if we augment the asymptotic behavior with contributions from a single spin zero quark of charge ($-1/3$). The visible hadronic fragments from such scalar quarks are not likely to produce predominantly pure 2-jet events at such low energies. Hence, such decay modes may have been removed by the Crystal Ball energy imbalance cuts in their definition of hadronic events but not in MARK I events, thus accounting for the discrepancy in the two results. Upper bounds on spin zero quark production at LEP through Z decay data are used to estimate the mixing angle between $T_3 = -1/2$ and $T_3 = 0$ scalar quarks. Recent negative results about spin zero quarks from CLEO are critically examined. We briefly discuss diquark production hypothesis and find it very unlikely to explain the discrepancy.

The purpose of this work is to shed some light on an ancient and as yet unresolved discrepancy in the measured values of $R(s)$ in the energy range $\sqrt{s} = 5 \div 7.5 \, \text{GeV}$, both measured at SLAC. The MARK I & II collaborations [1,2] found a higher value $R_{\text{av}} = 4.3 \pm 0.4$, whereas the Crystal Ball collaboration [3] found a lower value $R_{\text{av}} = 3.44 \pm 0.03 \pm 0.18$. There are a few data points in this region from PLUTO [4] and DASP [5] collaborations as well which are in agreement with the MARK I collaboration.

For the problem at hand, we consider the (Lorentz scalar) vacuum polarization function $\Pi_\gamma(s)$ which is defined via the EM polarization tensor $\Pi_{\mu\nu}(s)$ as:

$$ie^2 \int d^4x e^{iqx} \langle 0| T j^\mu_{\text{em}}(x) j^\nu_{\text{em}}(0) |0 \rangle = -(q^2 g^{\mu\nu} - q^\mu q^\nu) \Pi_\gamma(q^2),$$
where \( j_{\mu m}(x) \) is the electro-magnetic current. \( \Pi_\gamma(s) \) is an analytic function in the complex \( s \) plane with a branch cut on the real positive axis for \( s > s_0 \), where \( s_0 = 4m_\pi^2 \). We use the dispersion relation:

\[
\Pi_\gamma(t) - \Pi_\gamma(0) = \frac{t}{\pi} \int_{s_0}^{\infty} ds \frac{Im\Pi_\gamma(s)}{s(s - t - i\epsilon)}.
\] (1)

subtracted at \( t = 0 \), with the renormalized photon vacuum polarization: \( \Pi(t) = \Pi_\gamma(t) - \Pi_\gamma(0) \).

The optical theorem relates the imaginary part of the vacuum polarization \( \Pi(s) \) to the function \( R(s) \) by means of the identity:

\[
Im\Pi(s) = \left( \frac{\alpha}{3} \right) R(s).
\] (2)

One obtains, therefore, the dispersion integral [9]:

\[
\Pi(t) = \frac{t\alpha}{3\pi} \int_{s_0}^{\infty} ds \frac{R(s)}{s(s - t)}.
\] (3)

through which \( \Pi \) for space-like \( t \) gets related to \( R \) for time-like \( s \).

The idea is to consider this relation as an integral equation to compute \( R(s) \) in the interval \([5 \div 7.5 \text{ GeV}]\), using the following input: (i) experimental data in the time-like region from threshold up to 5 GeV and from 7.5 \( \div 100 \) GeV, (ii) PQCD asymptotic behavior for the rest of the time like region and (iii) in the space-like region only the asymptotic value of \( \Pi(t) \) calculated for values of \( t \) in the interval \([-205 \text{ GeV}^2 \leq t \leq -200 \text{ GeV}^2]\) through PQCD. Then we have:

\[
\Pi(t) - I_{01}(t) - I_{23}(t) - I_a(t) = \frac{t\alpha}{3\pi} \int_{s_1}^{s_2} ds \frac{R(s)}{s(s - t)}.
\] (4)

where:

\[
I_{01} = \frac{t\alpha}{3\pi} \int_{s_0}^{s_1} ds \frac{R_{\text{exp}}(s)}{s(s - t)} , \quad I_{23} = \frac{t\alpha}{3\pi} \int_{s_2}^{s_3} ds \frac{R_{\text{exp}}(s)}{s(s - t)} , \quad I_a(t) = \frac{t\alpha}{3\pi} \int_{s_3}^{\infty} ds \frac{R_a(s)}{s(s - t)}.
\] (5)

with: \( s_1 = (5 \text{ GeV})^2 \), \( s_2 = (7.5 \text{ GeV})^2 \) and \( s_3 = (100 \text{ GeV})^2 \). The function \( R_{\text{exp}}(s) \) used in the first and second integral (5) is a fit of the experimental data, the function \( R_a(s) \) in the third integral is the PQCD asymptotic behavior at leading order:

\[
R_a = N_c \sum_{q} Q_q^2 \left[ 1 + \frac{\alpha_s}{\pi} \right],
\] (6)

where \( N_c \) is the colour factor and \( Q_q \) is the charge of the quark \( q \). We note that higher order QCD correction terms (that is, beyond the \( \alpha_s/\pi \) term) are indeed “corrections to a correction” and thus play a negligible role in our analysis of the discrepancy.

The five light quark \((u, d, s, c, b)\) contributions to \( \Pi(t) \) can be safely calculated only at large \(-t\), since at low energies the quark interactions are modified considerably by strong interactions. At high energies, by virtue of asymptotic freedom inherent in QCD, we can treat the quark contribution similar to that for the leptons and, at leading order, we obtain:

\[
\Pi(t) = -N_c \frac{\alpha}{3\pi} \sum_{q} Q_q^2 \left[ \ln \left( \frac{-t}{m_q^2} \right) - \frac{5}{3} + O \left( \frac{m_q^2}{t} \right) \right].
\] (7)

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The difficulty in using this function comes from the masses, which is particularly severe for the light quarks $u, d$ and $s$, that are not unambiguously defined. To avoid this problem, but without any further approximation, we consider the derivative in $t$ of the dispersion integral (3) so as to obtain a new integral equation:

$$\frac{d\Pi(t)}{dt} - \frac{dI_{01}(t)}{dt} - \frac{dI(t)}{dt} - \frac{dI_a(t)}{dt} = \frac{\alpha}{3\pi} \int_{s_1}^{s_2} ds \frac{R(s)}{(s-t)^2}. \quad (8)$$

The asymptotic behavior of $\frac{d\Pi(t)}{dt}$ for large $-t$ is:

$$\left[\frac{d\Pi(t)}{dt}\right]_A = -N_c \frac{\alpha}{3\pi} \sum_q Q_q^2 \frac{1}{t}. \quad (9)$$

It has not just the virtue of being independent of the quark masses: it requires only a knowledge of the quark charges and the colour factor. Upon solving the integral equation (8) [6-8], we can calculate the values of $R(s)$ in the disputed region $[s_1, s_2]$ using as input the time-like experimental data and the quark charges.

Including only the five standard (spin 1/2) light quarks $u, d, s, c, b$, as shown in Fig.(1a), we find a value for $R(s)$ in good agreement with the Crystal Ball data [10].

But, if we include an additional contribution from a single species of spin zero quark of charge (-1/3), which requires simply adding in the summations (6) and (9) another charge (-1/3) squared, multiplied by a factor 1/4 (due to zero spin of the quark), we reproduce the MARK I data, shown in Fig.(1b).
Figure 1b. Values of $R(s)$ (light grey band) obtained from integral equation (8) with five spin 1/2 quarks and one scalar quark.

In Fig.(2a) and Fig.(2b), we show a similar comparison but imposing continuity at the boundaries.

Figure 2a. Values of $R(s)$ (light grey band) obtained from eq. (8) with five spin 1/2 quarks. Continuity imposed at 5 GeV and 7.5 GeV.

Figure 2b. Values of $R(s)$ (light grey band) obtained from eq. (8) with five spin 1/2 quarks and one scalar quark. Continuity imposed at 5 GeV and 7.5 GeV.

As a consistency check, we write a dispersion relation for $\Delta R(s)$ in the above region whose asymptotic value contains contribution from one species of colour triplet, charge $(-1/3)$, spin zero quark only. Phenomenologically, in the range $s_3 \geq s \geq s_2$, the standard 5 quark model gives a rather good description of the data. Hence, we obtain the approximate relation

$$
\int_{s_1}^{s_2} \frac{\Delta R(s) ds}{(s-t)^2} \approx -\frac{1}{12t} \left[ 1 + \frac{\bar{t}}{(s_3-t)} \right]
$$

This gives an average value $\Delta R \approx 0.75$. By way of comparison, the difference $\Delta R(s)$ between
the two curves from Fig.(1) are plotted in Fig.(3a). Fig.(3b) shows the difference $\Delta R(s)$ for Fig.(2).

These average values are mutually consistent. Such a “locally averaged” $\Delta R$ should not be confused with $\Delta R_{\text{asymptotic}}$. The above has been obtained under the assumption that beyond the disputed region, experimental data are adequately described by the standard $5$ spin $1/2$ quark model. Thus, roughly ($\Delta R_{\text{asy}} \times (-t) \approx (\Delta R) \times (s_2 - s_1)$).

The two experiments can be reconciled only under the hypothesis that Crystal Ball experimental cuts may have eliminated signals due to the spin zero quarks. The bound states and resonances with $J^{PC} = 1^{--}$ coupling to $e^+e^-$ are in relative $P$-wave for spin zero quarks. They would be relatively close in energy and more importantly, are expected to decay copiously into a low energy photon ($E_\gamma \approx 300$ MeV) recoiling against $S$-wave states which would subsequently decay into hadrons via 2 gluons [11]. Such radiative decay events would have a very asymmetric energy pattern. For the model discussed below, a similar pattern may also follow (for the non resonant) scalar quark production. In their definition of hadronic events contributing to $R$, Crystal Ball imposed the following kinematic cuts for the energy imbalance between left-right, top-bottom and front-back hemispheres. If any of these fractional energy differences, $A_{\text{left-right}}$, $A_{\text{top-bottom}}$ or $A_{\text{front-back}}$ was less than 0.8, the event was classified as being due to beam-gas interactions, $\gamma\gamma$ collisions or large missing energy $\tau$ decays. Thus, such events were not included in $R$ by Crystal Ball.
In Fig.(4), we show a plot of the ratio between \(N_{bg}\), which is the number of the background events from beam-gas and beam-wall interactions, and the total collisions \(N_{coll}\) passing the hadron selection criteria, versus \(\sqrt{s}\) from Table II of ref.(3). That the rejected events have peaks and structures resembling the MARK I data appears significant and supports our hypothesis.

Now we turn to a rough estimate of the mass of such a scalar quark \((Y)\). Taking our cue from the steep rise in the CB rejected events between \(7.2 \div 7.4\) GeV, we may estimate \(m_Y \approx 3.6 \div 3.7\) GeV from the production threshold for a pair \((Y^+Y)\) of such quarks. The local rise in MARK I data, prior to this threshold, beginning around 5 GeV would however require that a single scalar quark be produced along with a couple of other light (spin \(1/2\) \(u, d, s, c\)) quarks. One possibility (out of many others) would be a charge \((-1/3)\) color triplet scalar quark carrying baryon number \((-2/3)\) to coincide with the quantum numbers of a standard \(\bar{U}\bar{D}\) (where \(U = u, c\) and \(D = d, s\)) quark pair. Then, the final state \(YUD\) would be allowed in \(e^+e^-\) reaction, with a threshold around 5 GeV. We expect no sharp structures associated with it and the level of production should decrease rapidly with \(s\) so as to reach its asymptotic parton level. Two points are worthy of note here. First, such a scalar quark would be an “elementary” particle and not a composite “diquark” state of two standard \(\bar{U}\bar{D}\) quarks. (For this case, there might be non trivial mixing and interference terms). While for the narrow window of pair produced resonances \((Y^+Y)\) states), non-relativistic potential models may be a reasonable rough guide, the production dynamics and level of cross section for multi-particle relativistic systems (with exotic quantum numbers such as \(YUD\)) are not easy to compute or estimate directly since non-relativistic potential models can not be employed for this purpose. The hadronization of scalar quarks would also be very different from those for the other quarks. Not knowing the nature of the beast, our modest aim here has been to estimate these contributions to \(R\) through unitarity, dispersion relations and asymptotic behavior which are consistent with data from outside this region. Secondly, if one were to study the mass distribution of an \(e^+e^-\) or \(\mu^+\mu^-\) in the final state produced say from a \(p\bar{p}\) initial state, we would expect to see an enhancement visible only around \(7.2 \div 7.4\) GeV (due to pair production of scalar quarks) since the probability to find \(YUD\) in the initial state would be negligible.

An alternative explanation of this rather substantial difference between MARK I and CB data has been offered by Eidelman and Jegerlehner [12]. According to them, this difference arises from QED and QCD radiative corrections. In the final analysis, our explanation of this difference also involves radiative events but its source is different. Ours is generated from an extra scalar quark and MARK I data are “physical”, whereas in the other explanation MARK I data should be divided (“corrected”) by the radiative effects. It should be firmly kept in mind that the \((g-2)\) results discussed in ref.(12) can not be used to discriminate between MARK I and CB data since the results are hardly changed whether one or the other data are used.

The non resonant part of events from scalar quark decays would also not be of the 2-jet type at low energies. Both in the ALEPH [13] and the CLEO [14] collaboration searches for scalar quarks, it has been assumed that each scalar quark decays weakly, viz., into \(c\bar{\nu}_s\) and \(e\bar{\nu}_s\), where \(\bar{\nu}_s\) is a scalar neutrino. For \(\sqrt{s} = 5 \div 7.5\) GeV and scalar quark mass in the \(3 \div 4\) GeV range, the decay events would not be of the 2-jet type. On the other hand, at much higher energies, for example \(\sqrt{s} = 160 \div 205\) GeV as in the ALEPH data, the decay products would be confined to the forward and backward directions similar to a 2-jet profile.

Evidence supporting a low mass spin zero quark must now be confronted with the extremely accurate decay systematics for the \(Z^0\) for various channels available from LEP [15-18]. A computation of the coupling of observable spin zero quarks to the \(Z\) boson requires a knowledge
of their weak iso-spin. Let \( Y_L \) and \( Y_R \) denote spin zero quarks with \( T_3 = -1/2 \) and \( T_3 = 0 \) respectively. After mixing, let \( Y^- \) and \( Y^+ \) denote the low and high mass eigenstates. If the higher of these masses happens to be larger than the \( Z^0 \) mass, kinematically \( Z^0 \) can decay only into \( Y^+ Y^- \), with the decay amplitude given by

\[
\text{Amplitude}(\ Z^0(P) \rightarrow Y^+(p_1)Y^-(p_2)) = \left( \frac{e}{\sin \vartheta_W \cos \vartheta_W} \right) K \epsilon \mu(p_1 - p_2)^\mu,
\]

where \( K \) is given by

\[
K = \left( \frac{1}{3} \sin^2 \vartheta_W \right) |U_{R-}|^2 + \left( -\frac{1}{2} + \frac{1}{3} \sin^2 \vartheta_W \right) |U_{L-}|^2.
\]

Here \( U_{R-} = \cos \delta \) and \( U_{L-} = \sin \delta \) are the mixing matrix elements to the lower mass state \((-)\). Thus, the branching ratio of \( Z^0 \) into a pair of low mass spin zero quarks can be arbitrarily small, even zero for \( \delta^* \)

\[
\delta^* = \arcsin \left( \sqrt{\frac{2\sin^2 \vartheta_W}{3}} \right) \approx 23^\circ.
\]

Present measurements of the \( Z^0 \) width and observed branching ratios therefore can serve to place limits on the mixing angle \( \delta \). For example, if we assume that the branching ratio for the spin zero quark decay channel is less than one part per mille, approximately \( \delta = (23 \pm 15)^\circ \) (for \( \delta \) chosen in the first quadrant).

The recent CLEO experiment has given quite stringent bounds regarding a low mass squark partner of the \( b \) quark. They have looked for and not found an expected high level \( D \) and \( D^* \) signal. Of course, it could be that the low mass spin zero quark is not the partner of the \( b \) quark or even if it were so, the squark mixing matrix may not be similar to the quark mixing matrix. Also, the estimate of the weak decay signal expected in the CLEO experiment is based on phase space, i.e., a structureless matrix element and an essentially zero mass scalar neutrino. Neither of these hypotheses can be justified theoretically. For example, even a 300 \( MeV \) neutrino would lower the signal to practically its background value. We are not aware of any argument which can so restrict the scalar neutrino mass.

Among other presently available beams, high luminosity asymmetric B-factories offer a good avenue for new and independent check of data on \( R \) in the \( 5 \div 8 \) \( GeV \) region, by observing a bremsstrahlung photon of energy \( 1.3 \div 3.8 \) \( GeV \). An independent check of scalar quark production at B-factories would be through their weak decays by looking at changes in the number of opposite sign leptons (in the opposite hemispheres) as the beam energy passes through the \( bb \) threshold. Signals from spin zero quarks should be looked for also in \( \mu^+ \mu^- \) production and in the hadronic machines via the gluon induced process \( g + \bar{D} \rightarrow Y + U \) (for the mechanism discussed earlier). In addition, the next Tevatron run offers the possibility of looking for \( Y \) jets as well as a means to probe for \( \tilde{w} \) (any generic spin 1/2 object recoiling against the \( Y \)) up-to a mass of about 170 \( GeV \) through the top quark decay \( t \rightarrow Y \ + \tilde{w}^+ \).

We now discuss a completely different hypothesis to resolve the discrepancy. If one accepts that the difference between the two experiments is due to one experiment being essentially sensitive only to 2-jet events and the other not so biased, this may be accounted for through the production of a diquark-antidiquark pair, since they would produce predominantly 4 quark final
states. As the diquark photon vertex includes a form factor their production would disappear at high energies, e.g., at LEP.

A diquark diagram (inclusive of a form factor) to the self energy of the photon is shown in Fig.(5).

![Diquark Diagram](image)

**Figure 5.** Diquarks photon self energy.

It contributes to the imaginary part

$$\text{Im}(\pi_{DQ}(s)) = \frac{1}{4} \frac{N_c \alpha Q^2}{3} |F(s)|^2 \left(1 - \frac{4m^2}{s}\right)^{\frac{3}{2}} \quad (14)$$

and thus to the dispersion relation

$$\hat{\pi}_{DQ}(t) = \frac{1}{4} \frac{N_c \alpha Q^2}{3} \frac{t}{\pi} \int_{4m^2}^\infty \left[1 - \frac{4m^2}{s}\right]^{\frac{3}{2}} \frac{|F(s)|^2}{s(s-t)} \, ds. \quad (15)$$

$F(s)$ is normalized to 1 at $s = 0$ and is expected to go asymptotically as

$$|F(s)| \sim \frac{Q^2_0}{s}, \quad (16)$$

with $Q^2_0 \simeq 10 \text{ GeV}^2$ [19]. For a diquark of charge $q$ to saturate the discrepancy in $\Delta R = 0.75$ (as discussed previously), we require

$$|F(s)| = \begin{cases} \frac{3}{N_c q} & \text{if } s \in [s_1, s_2] \\ \frac{N_c q}{s} Q^2_0 & \text{if } s > s_2 \end{cases}, \quad (17)$$

to be fed in

$$\frac{d\hat{\pi}_{DQ}(t)}{dt} = \frac{\alpha}{3\pi} \frac{1}{12} \int_{s_1}^\infty \left(1 - \frac{s_1}{s}\right)^{\frac{3}{2}} \frac{|F(s)|^2}{(s-t)^2} \, ds. \quad (18)$$

Unfortunately, the contribution of such diquarks in the space like region falls short by a factor 4. Hence, the hypothesis of diquarks being responsible for the discrepancy does not appear to be internally self consistent.

In conclusion, our dispersive analysis of the MARK I data (the only set included in PDG, the particle data tables [20]) appears to require the existence of a charge ($-1/3$) scalar quark with a low mass. While the level of the reported cross section appears to be consistent with our analysis, it remains an open dynamical problem how to generate such a large value. Preliminary estimates of contributions to $R$ from bound states computed through static potentials appear to fall short by a factor of 3 or more. This problem obviously deserves further study and the expected data from BABAR on a new measurement of $R$ in the disputed region suggested here
would be most useful. On the other hand, lack of evidence in the present LEP data about the
direct production of such low mass scalar quarks in $Z^0$ decay only serves to indicate that it is
predominantly an iso-singlet. On the other hand, a $b$ jet from the $Z^0$ decay may sequentially
produce such a scalar quark whose signature might be revealed through an analysis of the final
$\mu^+\mu^-$ mass distribution in the $7.2 \div 7.4$ GeV region.

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