TESTS OF QUANTUM CHROMODYNAMICS IN EXCLUSIVE e^+e^- and $\gamma\gamma$ PROCESSES*

STANLEY J. BRODSKY
Stanford Linear Accelerator Center,
Stanford University, Stanford, California 94309, USA

1. INTRODUCTION

One of the most important areas of investigation in quantum chromodynamics are few-body exclusive reactions initiated by electromagnetic initial states, such as $e^+e^- \to H\bar{H}$, $e^+e^- \to \gamma H$, and the two-photon processes $\gamma\gamma \to H\bar{H}$ shown in Fig. 1. The simplicity of the photon's couplings to the quark currents and the absence of initial state hadronic interactions allows one to study the process of quark hadronization at its most basic level—the conversion of quarks into just one or two hadrons. In the low energy threshold regime the quarks interact strongly at low relative velocity to form ordinary or exotic resonances: $q\bar{q}$, $q\bar{q}g$, $qq\bar{q}\bar{q}$, ggg, etc. At high energies, where the quarks must interact at high momentum transfer, a perturbative expansion in powers of the QCD running coupling constant becomes applicable,[1] leading to simple and elegant PQCD predictions. In this domain one tests not only the scaling and form of elementary quark-gluon processes, but also the structure of the hadronic wavefunctions themselves, specifically, the "distribution amplitudes" $\phi_H(x_i, Q^2)$, which describe the binding of quarks and gluons into hadrons. Physically, $\phi_H(x_i, Q)$ is the probability amplitude for finding the valence quarks which carry fractional momenta x_i at impact separation $b_i \sim 1/Q$. The valence Fock state of a hadron is defined at a fixed light-cone time and in light-cone gauge. The $x_i = (k^0 + k^z)/(P^0 + P^z)$ are the boost-invariant momentum fractions which satisfy $\sum_i x_i = 1$. Such wavefunction information is critical not only for understanding QCD from first principles, but also for a fundamental understanding of jet hadronization at the amplitude rather than probabilistic level.

* Work supported by the Department of Energy, contract DE-AC03-76SF00515.

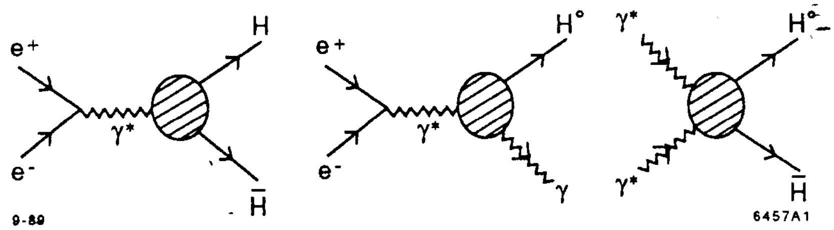

Figure 1. Exclusive processes from e^+e^- and $\gamma\gamma$ annihilation.

At large momentum transfer all exclusive scattering reactions in QCD are characterized by the fixed angle scaling law:

$$\frac{d\sigma(AB \to CD)}{dt} \simeq \frac{F(\theta_{\rm cm})}{s^N}.$$

To first approximation the leading power is set by the sum of the minimum number of fields entering the exclusive amplitude: $N = n_A + n_B + n_C + n_D - 2$, where $n = 3$ for baryons, $n = 2$ for mesons, and $n = 1$ for leptons and photons. This is the dimensional counting law[2] for the leading twist or power-law contribution. The nominal power N is modified by logarithmic corrections from the QCD running coupling constant, the logarithmic evolution of the hadronic distribution amplitudes, and in the case of hadron-hadron scattering, so-called "pinch" or multiple-scattering contributions, which lead to a small fractional change in the leading power behavior.[3] The recent analysis of Botts and Sterman[3] shows that hard subprocesses dominate large momentum transfer exclusive reactions, even when pinch contributions dominate. The functional form of $F(\theta_{\rm cm})$ depends on the structure of the contributing quark-gluon subprocess and the shape of the hadron distribution amplitudes.

Large momentum transfer exclusive amplitudes generally involve the $L_z = 0$ projection of the hadron's valence Fock state wavefunction. Thus in QCD, quark helicity conservation leads to a general rule concerning the spin structure of exclusive amplitudes: the leading twist contribution to any exclusive amplitude conserves hadron helicity—the sum of the hadron helicity in the initial state equals that of the final state.

The study of time-like hadronic form factors using e^+e^- colliding beams can provide very sensitive tests of the QCD helicity selection rule. This follows because the virtual photon in $e^+e^- \to \gamma^* \to h_A \bar{h}_B$ always has spin ± 1 along the beam axis at high energies. Angular-momentum conservation implies that the virtual photon can "decay" with one of only two possible angular distributions in the center-of-momentum frame: $(1+\cos^2\theta)$ for $|\lambda_A - \lambda_B| = 1$, and $\sin^2\theta$ for $|\lambda_A - \lambda_B| = 0$, where $\lambda_{A,B}$ are the helicities of hadron $h_{A,B}$. Hadronic-helicity conservation, as required by QCD, greatly restricts the possibilities. It implies that $\lambda_A + \lambda_B = 2\lambda_A = -2\lambda_B$. Consequently, angular-momentum conservation requires $|\lambda_A| = |\lambda_B| = \frac{1}{2}$ for baryons and $|\lambda_A| = |\lambda_B| = 0$ for mesons; and the angular distributions are now completely determined:

$$\frac{d\sigma}{d\cos\theta}(e^+e^- \to B\bar{B}) \propto 1 + \cos^2\theta(\text{baryons}),$$
$$\frac{d\sigma}{d\cos\theta}(e^+e^- \to M\bar{M}) \propto \sin^2\theta(\text{mesons}).$$

It should be emphasized that these predictions are far from trivial for vector mesons and for all baryons. For example, one expects distributions like $\sin^2\theta$ for baryon pairs in theories with a scalar or tensor gluon. Simply verifying these angular distributions would give strong evidence in favor of a vector gluon.

In the case of $e^+e^- \to H\bar{H}$, time-like form factors which conserve hadron helicity satisfy the dimensional counting rule:

$$F_H(Q^2) \sim 1/(Q^2)^{N_H - 1}.$$

Thus at large $s = Q^2$, QCD predicts, modulo computable logarithms,

$$\lambda_B = -\bar{\lambda}_{\overline{B}} = \pm\tfrac{1}{2}, \quad Q^4 F_1^B(Q^2) \to const,$$

for baryon pairs, and

$$\lambda_M = \bar{\lambda}_{\overline{M}} = 0, \quad Q^2 F^M(Q^2) \to const$$

for mesons. Other form factors, such as the Pauli form factor which do not conserve hadron helicity, are suppressed by additional powers of $1/Q^2$. Similarly, form factors for processes in which either hadron of the pair is produced with helicity other than 1/2 or 0 are non-leading at high Q^2.

In the case of e^+e^- annihilation into vector plus pseudoscalar mesons, such as $e^+e^- \to \rho\pi, \pi\omega$, and KK^*, Lorentz invariance requires that the vector meson will be produced transversely polarized. Since this amplitude does not conserve hadron helicity, PQCD predicts that it will be dynamically suppressed at high momentum transfer.

We can see this in more detail as follows: The $\gamma - \pi - \rho$ can couple through only a single form factor – $\epsilon^{\mu\nu\tau\sigma}\epsilon_\mu^{(\gamma)}\epsilon_\nu^{(\rho)}p_\tau^{(\pi)}p_\sigma^{(\rho)}F_{\pi\rho}(s)$ — and this requires $|\lambda_\rho| = 1$ in e^+e^- collisions. Hadronic-helicity conservation requires $\lambda = 0$ for mesons, and thus these amplitudes are suppressed in QCD (although, not in scalar or tensor theories). Notice however that the processes $e^+e^- \to \gamma\pi, \gamma\eta, \gamma\eta'$ are allowed by the helicity selection rule; helicity conservation applies only to the hadrons. The form factors governing these such processes are not expected to be large, *e.g.* $F_{\pi\gamma}(s) \sim 2f_\pi/s$.

The hadron helicity conservation rule has also been used to explain the observed strong suppression of ψ' decay to $\rho\pi$ and KK^*. However, a puzzle then arises why the corresponding J/ψ decays are not suppressed. I will review this problem in section 7.

The predictions of PQCD for the leading power behavior of exclusive amplitudes are rigorous in the asymptotic limit. Analytically, this places important constraints on the form of the amplitude even at low momentum transfer. For example, Dubnicka and Etim[4] have made detailed predictions for meson and baryon form factors based on vector meson dominance considerations at low energies, and the PQCD constraints in the large space-like and time-like Q^2 domains. (See Fig. 2.)

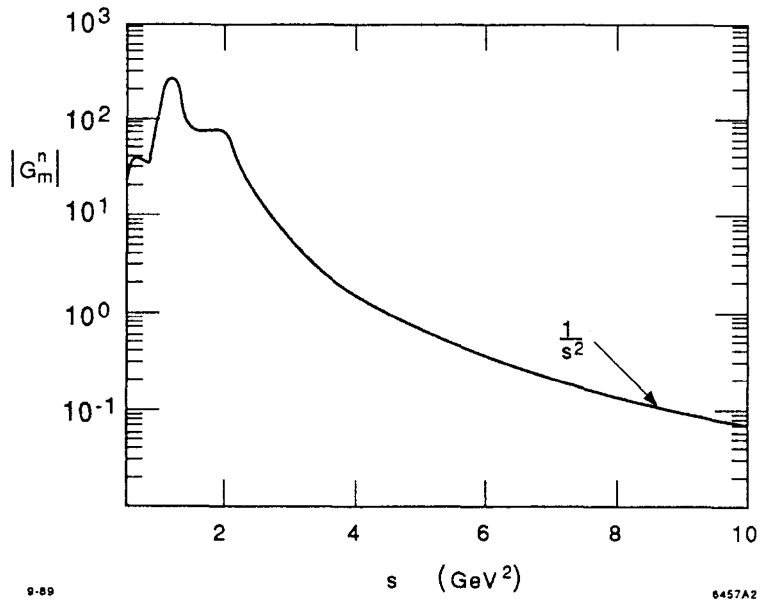

Figure 2. Prediction for the time-like magnetic form factor of the neutron using vector meson dominance and asymptotic PQCD constraints. From Ref. 4.

A central question for the phenomenology of exclusive reactions is the regime of applicability of the leading power-law predictions and the relative size of higher-twist higher power-law contributions. Thus far dimensional counting rules are all in general agreement with experiment at momentum transfers beyond a few GeV. This appears reasonable since, ignoring heavy quark production, the natural expansion scales of QCD are $\Lambda_{\overline{MS}}$, the light quark masses, and the intrinsic transverse momentum in the hadronic wavefunctions. An extensive review of the data is given in Ref.1.

The proposed high-luminosity Tau-Charm factory would allow the exploration of a large array of exclusive channels such as e^+e^- or $\gamma\gamma \to p\bar{p},\ n\bar{n},\ \Lambda\bar{\Lambda},\ \pi^+\pi^-, K\bar{K}, N\bar{N}^*,\ \pi\rho,\ \gamma\pi^0$, etc., both on and off the charmonium resonances. Many of these channels have not yet been studied experimentally, and measurements will only become practical at luminosities of $10^{33}\text{cm}^{-2}\text{sec}^{-1}$ or greater. At such intensities, corresponding to approximately $10^8\mu^+\mu^-$ per year, one can also study *nuclear* final states such as $e^+e^- \to \bar{d}np$. It is very important to measure the ratio of the neutron and proton form factors to high precision, and to check the angular distribution of the baryon pairs to test the predicted dominance of the helicity conserving Dirac form factor F_1 over the Pauli form factor at large time-like Q^2.

Since exclusive channels have highly constrained final states of minimal complexity, they are generally distinctive and background-free. In each exclusive channel one tests not only the scaling and helicity structure of the quark and gluon processes, but also features of the distribution amplitude, the most basic measure of a hadron in terms of its valence quark degrees of freedom.

Although two-photon exclusive channels are usually measured at higher energy storage rings, the high luminosity at a Tau-Charm factory makes possible detailed study of the $\gamma\gamma \to H\bar{H}$ channels in the few GeV region where perturbative QCD predictions begin to become applicable. A more detailed review of the two-photon predictions applicable to the Tau-Charm factory are given in section 11 and Ref. 5.

2. Factorization Theorem For Exclusive Processes

The predictions of QCD for the leading twist contribution to exclusive e^+e^- and $\gamma\gamma$ annihilation amplitudes have the general form:

$$M(e^+e^- \to H\bar{H}) = \int_0^1 \Pi dx_i\ T_H(x_i, \alpha_s(Q^2))\ \phi_H(x,Q)\ \phi_{\bar{H}}(x,Q).$$

The hard-scattering amplitude $T_H(e^+e^- \to q\bar{q}q\bar{q})$ is computed by replacing each hadron with its collinear valence quarks. By definition, the internal integrations in T_H are restricted to transverse momentum greater than an intermediate scale $\tilde{Q}$; it is thus free of infrared or collinear divergences and it can be expanded systematically in powers of $\alpha_s(\tilde{Q}^2)$. The distribution amplitudes are gauge-invariant wavefunctions obtained by integrating the valence Fock State wavefunctions over transverse momentum up to the scale $\tilde{Q}$. As in the case of the factorization theorem for inclusive reactions, it is convenient to choose the intermediate renormalization scale $\tilde{Q}$ to be of order Q in order to minimize large higher order terms.

The distribution amplitude $\phi_H(x, Q)$ satisfies an evolution equation in $log\ Q^2$ which sums all logarithms from the collinear integration regime. The solution has the form

$$\phi_H(x_i, Q) = \sum_n a_n^H C_n(x_i) log^{-\gamma_n} Q^2$$

where the C_n are known polynomials, the fractional numbers γ_n are computed anomalous dimensions, and the a_n^H are determined from an initial condition or non-perturbative input for $\phi_H(x_i, Q_0)$. The results for meson pair production are rigorous in the sense that they are proved to all orders in perturbation theory. In the case of baryon pair production, one can use an all-orders resumation to show that the soft region of integration where $x \sim 1$ is, in fact, Sudakov suppressed.

3. Electromagnetic Form Factors of Baryons

Applying factorization, any helicity-conserving baryon form factor at large space-like or time-like Q^2 has the form: (see Fig. 3)

$$F_B(Q^2) = \int_0^1 [dy] \int_0^1 [dx]\ \phi_B^\dagger(y_j, Q) T_H(x_i, y_j, Q) \phi_B(x_i, Q) \ ,$$

where to leading order in $\alpha_s(Q^2)$, T_H is computed from $3q + \gamma^* \to 3q$ tree graph amplitudes:

$$T_H = \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^2 f(x_i, y_j)$$

and

$$\phi_B(x_i, Q) = \int [d^2k_\perp]\ \psi_V(x_i, \vec{k}_{\perp i}) \theta(k_{\perp i}^2 < Q^2)$$

is the valence three-quark wavefunction evaluated at quark impact separation $b_\perp \sim \mathcal{O}(Q^{-1})$. Since ϕ_B only depends logarithmically on Q^2 in QCD, the main

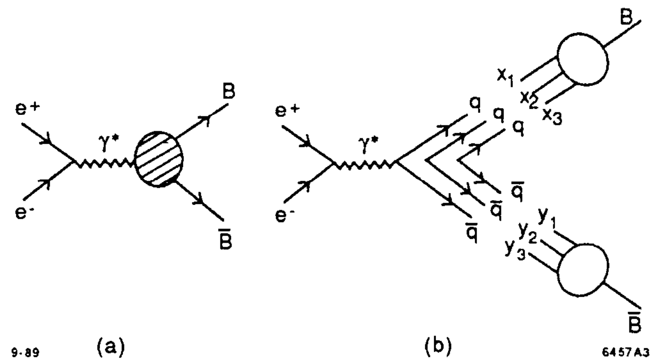

Figure 3. Calculation of the time-like baryon form factor from PQCD factorization.

dynamical dependence of $F_B(Q^2)$ is the power behavior $(Q^2)^{-2}$ derived from the scaling behavior of the elementary propagators in T_H.

More explicitly, the proton's magnetic form factor has the form:[6]

$$G_M(Q^2) = \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^2 \sum_{n,m} a_{nm} \left(\log \frac{Q^2}{\Lambda^2}\right)^{-\gamma_n-\gamma_m}$$

$$\times \left[1 + \mathcal{O}(\alpha_s(Q)) + \mathcal{O}\left(\frac{1}{Q}\right)\right] .$$

The first factor, in agreement with the quark counting rule, is due to the hard scattering of the three valence quarks from the initial to final nucleon direction. Higher Fock states lead to form factor contributions of successively higher order in $1/Q^2$. The logarithmic corrections derive from an evolution equation for the nucleon distribution amplitude. The γ_n are the computed anomalous dimensions, reflecting the short distance scaling of three-quark composite operators. The results hold for any baryon to baryon vector or axial vector transition amplitude that conserves the baryon helicity. Helicity non-conserving form factors should fall as an additional power of $1/Q^2$.[8] Measurements[7] of the transition form factor to the $J = 3/2$ $N(1520)$ nucleon resonance are consistent with $J_z = \pm 1/2$ dominance, as predicted by the helicity conservation rule.[8] A review of the data on spin effects in electron nucleon scattering in the resonance region is given in Ref. 7. A Tau-Charm factory could provide measurements on the whole range of baryon pair production processes, including hyperon production, isobar production, etc.

An essential question for the interpretation of such experiments is the scale of momentum transfer where leading-twist PQCD contributions dominate exclusive amplitudes.

The perturbative scaling regime of the meson form factor and $\gamma\gamma \to M\bar{M}$ amplitudes is primarily controlled by the virtuality of the hardest quark propagator– if the quark is far off-shell, multiple gluon exchange contributions involving soft gluon insertions are suppressed by inverse powers of the quark propagator. Thus non-leading twist contributions are suppressed by powers of $\mu^2/ < (1-x)Q^2 >$, where μ^2 is a typical hadronic scale. Physically, there is not sufficient time to exchange soft gluons or gluonium. Thus the perturbative analysis is valid as long as the single gluon exchange propagator can be approximated by inverse power behavior $D(k^2) \propto 1/k^2$. The gluon virtuality $< (1-x)(1-y)Q^2 >$ thus needs to be larger than a small multiple of $\Lambda^2_{\overline{MS}}$. This allows the PQCD predictions to start to be valid at Q^2 of order a few GeV^2, which is consistent with data.

However, the normalization of the leading twist predictions may be strongly affected by higher corrections in $\alpha_s(Q^2)$. A similar situation occurs in time-like inclusive reactions, such as massive pair production, where large K factors occur. Thus st this time normalization predictions for exclusive amplitudes cannot be considered decisive tests of PQCD.

The predictions for the leading twist contributions to the magnitude of the proton form factor are sensitive to the $x \sim 1$ dependence of the proton distribution amplitude,[9] particularly if one assumes the validity of the strongly asymmetric QCD sum rule forms for distribution amplitude. Chernyak, *et al.*[11] have found, however, that their QCD sum rule predictions are not significantly changed when higher moments of the distribution amplitude are included. In the analysis of Ref. 12 it was argued that only a small fraction of the proton and pion form factor normalization at experimentally accessible momentum transfer comes from regions of integration in which all the propagators are hard. However, a new analysis by Dziembowski, *et al.*[13] shows that the QCD sum rule distribution amplitudes of Chernyak, *et al.*[10] together with the perturbative QCD prediction gives contributions to the form factors which agree with the measured normalization of the pion form factor at $Q^2 > 4\ GeV^2$ and proton form factor $Q^2 > 20\ GeV^2$ to within a factor of two. In this calculation the virtuality of the exchanged gluon is restricted to $|k^2| > 0.25\ GeV^2$. The authors assume $\alpha_s = 0.3$ and that the underlying wavefunctions fall off exponentially at the $x \simeq 1$ endpoints. Another model of the proton distribution amplitude with di-quark clustering[14] chosen to satisfy the QCD sum rule moments come even closer. Considering the uncertainty in the magnitude of the higher order corrections, one cannot expect better agreement between the

QCD predictions and experiment.

Measurements of rare exclusive processes are essential for testing the PQCD predictions and for placing constraints on hadron wavefunctions. However, the relative importance of non-perturbative contributions to form factors clearly remains an important issue. Models can be constructed in which non-perturbative effects persist to high Q.[12] In other models, which are explicitly rotationally invariant,[15] such effects vanish rapidly as Q increases.[16,17,18,19] The resolution of such uncertainties will require better understanding of the non-perturbative wavefunction and the role played by Sudakov form factors in the end-point region. In the case of elastic hadron-hadron scattering amplitudes, the recent analysis of Botts and Sterman[3] shows that, because of Sudakov suppression, even pinch contributions are dominated by hard gluon exchange subprocesses.

If the QCD sum rule results are correct, then hadrons have highly structured momentum-space valence wavefunctions. In the case of mesons, the results from both the lattice calculations and QCD sum rules show that the pion and other pseudo-scalar mesons have a dip structure at zero relative velocity their distribution amplitude– the light quarks in hadrons are highly relativistic. This gives further indication that while nonrelativistic potential models are useful for enumerating the spectrum of hadrons (because they express the relevant degrees of freedom), they may not be reliable in predicting wavefunction structure.

4. Suppression of Final State Interactions

In general, one expects exclusive amplitudes to be complicated by strong hadronic final state interactions. For example, the intermediate process $e^+e^- \rightarrow p\bar{p}$ shown in Fig. 4 leads by charged pion exchange to a contribution to neutron pair production $e^+e^- \rightarrow n\bar{n}$. Such final-state interactions corrections to the time-like neutron form factor correspond to higher Fock contributions of the neutron wavefunction. By dimensional power counting, such terms are suppressed at large Q^2 by at least two powers of $1/Q^2$. Thus final state interactions are dynamically suppressed in the high momentum transfer domain.

Because of the absence of meson exchange and other final state interactions, the perturbative QCD predictions for the time-like baryon form factors are relatively uncomplicated, and directly reflect the coupling of the virtual photon to the quark current. For example, in the case of the ratio of nucleon magnetic form factors $G_M^n(Q^2)/G_M^p(Q^2)$, the ratio of quark charges $e_d/e_u = -1/2$ is the controlling factor. Various model wave functions have been proposed to describe the nucleon distribution amplitudes. In the case of the QCD sum rule wavefunction calculated by Chernyak, Ogloblin, and Zhitnitskii, the neutron to proton form factor ratio is

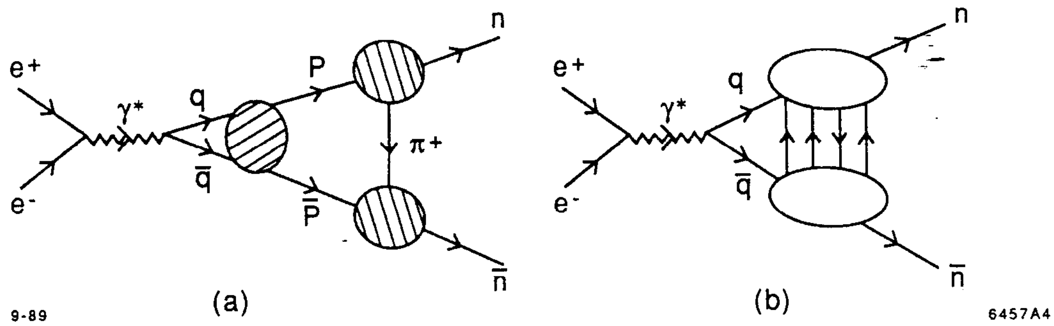

Figure 4. Illustration of a final-state interaction correction to the time-like neutron form factor. As shown in (b), the meson exchange contributions correspond to higher Fock components in the neutron wavefunction and are suppressed at high Q^2.

predicted to be -0.47 because of the strong dominance at large light-cone momentum fraction x of the u quark which has its helicity aligned with of the helicity of the proton. An alternative model given by Gary and Stefanis gives a much smaller ratio: -0.10. Both the COZ and GS model forms for $\phi_p(x_i, Q)$ taken together with the PQCD factorization formula can account for the magnitude and sign of the proton form factor at large space-like Q^2 : $Q^4G_M^p(Q^2) = 0.95\ GeV^4$ for COZ and $1.18\ GeV^4$ for GS. (See Fig. 5.) Experimentally $Q^4G_M^p(Q^2) \approx 1.0\ GeV^4$ for $10 < Q^2 < 30\ GeV^2$. These QCD sum rule predictions assume a constant value for the effective running coupling constant, $\alpha_s(\tilde{Q}^2) = 0.3$. The validity of such predictions for the absolute normalization of form factors is thus in considerable doubt, particularly because of the many uncertainties from higher order corrections. Still it should be noted that the predictions of the general magnitude and sign is non-trivial. For example, a "non-relativistic" nucleon distribution amplitude proportional to $\delta(x_1 - 1/3)\delta(x_2 - 1/3)$ gives $Q^4G_M^p(Q^2) = -0.3 \times 10^{-2}$.

In the case of the inverse process, $\bar{p}p \to e^+e^-$, initial state interactions are suppressed. It is interesting to consider the consequences of this PQCD prediction if the $\bar{p}p$ annihilation occurs inside a nucleus, as in the quasi-elastic reaction $\bar{p}A \to e^+e^-(A-1)$. The absence of initial state interactions implies that the reaction rate for exclusive annihilation in the nucleus will be additive in the number of protons Z. This is the prediction of "color transparency."[21] In general, this novel feature of large momentum quasi-elastic processes in nuclei is a consequence of the small color dipole moment of the hadronic state entering the exclusive amplitude. Even in the case of hadronic scattering such as $pp \to p\bar{p}$ where pinch contributions are important, one can show[3] that the impact separation of the quarks entering the subprocess is small, almost of order $1/Q$ so that color transparency is a universal feature of the PQCD predictions.

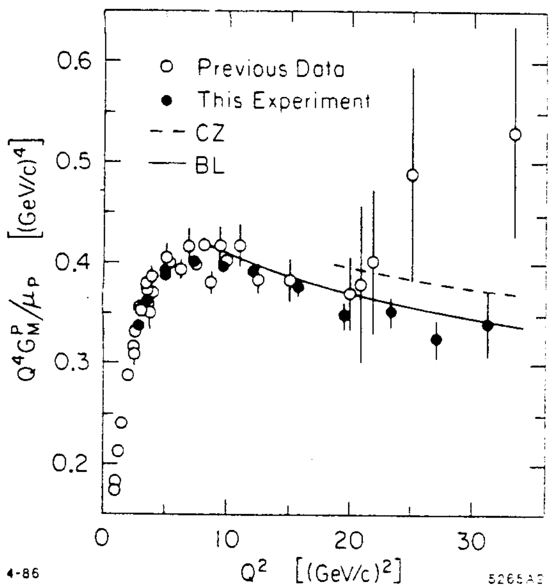

Figure 5. Comparison of the scaling behavior of the proton magnetic form factor with the theoretical predictions of Refs. 6 and 10. The slow fall-off is mainly due to the QCD running coupling constant. The CZ predictions[10] are normalized in sign and magnitude. The data are from Ref. 20.

An important test of color transparency was recently made at BNL through measurements of the nuclear dependence of quasi-elastic large angle pp scattering in nuclei. Conventional analysis of the absorptive initial and final state interactions predict that only $\sim 15\%$ of the protons are effective scatters in large nuclei. The results for various energies up to $E_{\rm cm} = 5\ GeV$ show that the effective fraction of protons Z_{eff}/Z rises monotonically with momentum transfer to about 0.5, as predicted by PQCD color transparency, contrary to the conventional Glauber analyses. However, at $E_{\rm cm} \sim 5\ GeV$, normal absorption was observed, contrary to the PQCD predictions. This unexpected and anomalous behavior, as well as the sharp features observed in the spin correlation A_{NN} seen in large angle pp scattering at the same energy could be due to a resonance or threshold enhancement at the threshold for open charm production. Further discussion is given in Ref. 22.

5. THE $\gamma\pi_0$ TRANSITION FORM FACTOR

The most elementary exclusive amplitude in QCD is the photon-meson transition form factor $F_{\gamma\pi^0}(Q^2)$, since it involves only one hadronic state. As seen from the structure of the diagram in Fig. 6 that the leading behavior of $F\gamma\pi^0(Q^2)$ at large Q^2 is simply $1/Q^2$, reflecting the elementary scaling of the quark propagator at large virtuality. This scaling tests PQCD in exclusive processes in as basic a way as Bjorken scaling in deep inelastic lepton-hadron scattering tests the short distance behavior of QCD in inclusive reactions.

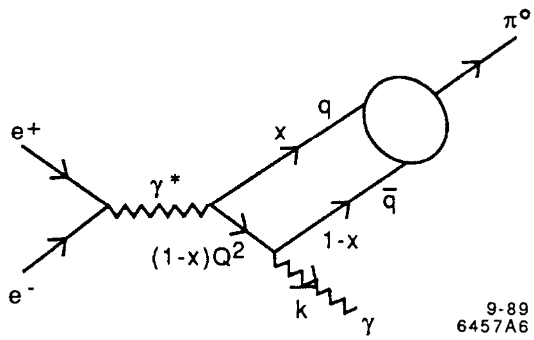

Figure 6. Illustration of the leading PQCD contribution to the $\gamma^* \to \pi^0\gamma$ time-like form factor.

One can easily show that the asymptotic behavior of the transition form factor has the simple form

$$F_{\gamma\pi^0} \propto \frac{1}{Q^2} \int_0^1 \frac{dx}{1-x} \phi_\pi(x,Q).$$

Thus

$$R(e^+e^- \to \gamma\pi^0) \propto \alpha | \int_0^1 \frac{dx}{1-x} \phi_\pi(x,Q)|^2 \sim 10^{-4}$$

at $Q^2 = 10\ GeV^2$. Detailed predictions are given in Ref. 1. Furthermore, the ratio of the pion form factor to the square of the $F_\gamma\pi^0$ transition form factor is directly proportional to $\alpha_s(Q^2)$, independent of the pion distribution amplitude. Thus measurements of this ratio at time-like Q^2 will give a new rigorous measure of the running QCD coupling constant.

Higher order corrections to $F_{\gamma\pi^0}$ from diagrams in which the quark propagator is interrupted by soft gluons are power-law suppressed. If the gluon carries high momentum of order Q, the corrections are higher order in $\alpha_s(Q^2)$. Unlike the meson and baryon form factors, there are no potentially soft gluon propagators in T_H for this process.

The scaling behavior of the PQCD prediction has recently been checked for the space-like $\gamma\eta$ and $\gamma\eta'$ transition form factors. This amplitude was obtained from measurements of tagged two-photon processes $\gamma^*\gamma \to \eta$ and η' by the $TPC/\gamma\gamma$ collaboration at PEP. The results, shown in figure 13, in section 11, provide a highly significant test of the PQCD analysis. Similarly, the time-like $\gamma^* \to \gamma\pi^0$ measurement would be one of the most fundamental measurements possible at the Tau-Charm factory.

6. Exclusive Charmonium Decays

The J/ψ decays into isospin-zero final states through the intermediate three-gluon channel. If PQCD is applicable, then the leading contributions to the decay amplitudes preserve hadron helicity. Thus as in the continuum decays, baryon pairs are predicted to have a $1 + v^2 cos^2\theta_{\rm cm}$ distribution with opposite helicities $\lambda = -\bar{\lambda} = \pm\frac{1}{2}$, and mesons with a $sin^2\theta_{\rm cm}$ distribution and helicity zero.

The calculation of the decay of the J/ψ to baryon pairs is obtained simply by (1) constructing the hard scattering amplitude T_H for $c\bar{c} \to ggg \to (q\bar{q})(q\bar{q})(q\bar{q})$ where the final qqq and $\bar{q}\bar{q}\bar{q}$ are collinear with the produced baryon and anti-baryon respectively, and (2) convoluting T_H with $\phi_B(x_i, \hat{Q})$ and $\phi_{\bar{B}}(y_i, \hat{Q})$. (See Fig. 7.) The scale $\hat{Q}$ is set by the characteristic momentum transfers in the decay. The J/ψ itself enters through its wavefunction at the origin which is fixed by its leptonic decay. Assuming a mean value $\alpha_s = 0.3$, one predicts $\Gamma(J/\psi \to p\bar{p}) = 0.34$ KeV for the recent QCD sum rule distribution amplitude proposed by Chernyak, Ogloblin, and Zhitnitskii. The QCD sum rule form obtained by King and Sachrajda predicts $\Gamma(J/\psi \to p\bar{p}) = 0.73$ KeV. Both models for the distribution amplitude together with the PQCD factorization for exclusive amplitudes can account for the magnitude and sign as well as the scaling of the proton form factor at large space-like Q^2. In contrast a non-relativistic ansatz for the distribution amplitude centered at $x_i = 1/3$ gives a much smaller rate: $\Gamma = 0.4 \times 10^{-3}$ KeV. The measured rate is 0.15 KeV. (Note that the PQCD prediction depends on α_s to the sixth power. Thus if the mean value of $\alpha_s = 0.26$, one finds agreement with the calculated rate for $J/\psi \to p\bar{p}$ using the COZ proton distribution amplitude.) The predicted angular distribution $1+ \cos^2\theta$ is consistent with published data.[23] This is important evidence favoring a vector gluon, since scalar- or tensor-gluon theories would predict a distribution of $\sin^2\theta + O(\alpha_s)$.

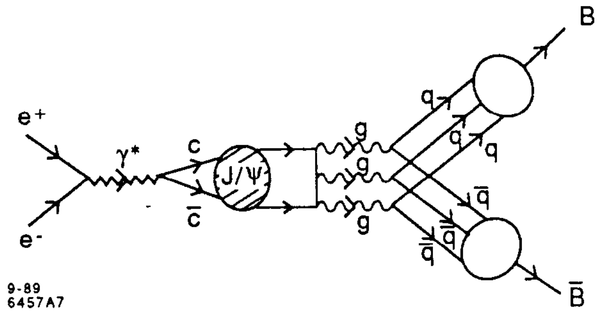

Figure 7. Illustration of the leading PQCD contribution for J/ψ decay to baryon pairs.

Dimensional-counting rules can also be checked by comparing the ψ and ψ' rates into $p\bar{p}$, normalized by the total rates into light-quark hadrons so as to remove dependence upon the heavy-quark wave functions. Theory predicts that the ratio of branching fractions for the $p\bar{p}$ decays of the ψ and ψ' is

$$\frac{B(\psi' \to p\bar{p})}{B(\psi \to p\bar{p})} \sim Q_{e^+e^-} \left(\frac{M_\psi}{M'_\psi}\right)^8 ,$$

where $Q_{e^+e^-}$ is the ratio of branching fractions into e^+e^-:

$$Q_{e^+e^-} \equiv \frac{B(\psi' \to e^+e^-)}{B(J/\psi \to e^+e^-)} = 0.135 \pm 0.023 \quad .$$

Existing data suggest a ratio $(M_{\psi'}/M_\psi)^n$ with $n = 6 \pm 3$, in good agreement with QCD. One can also use the data for $\psi \to p\bar{p}, \Lambda\bar{\Lambda}, \Xi\bar{\Xi}$, etc. to estimate the relative magnitudes of the quark distribution amplitudes for baryons. Correcting for phase space, one obtains $\phi_p \sim 1.04(13)\ \phi_n \sim 0.82(5)\ \phi_\Xi \sim 1.08(8)\ \phi_\Sigma \sim 1.14(5)\ \phi_\Lambda$ by assuming similar functional dependence on the quark momentum fractions x_i for each case.

As is well known, the decay $\psi \to \pi^+\pi^-$ must be electromagnetic if G-parity is conserved by the strong interactions. To leading order in α_s, the decay is through a virtual photon (*i.e.* $\psi \to \gamma^* \to \pi^+\pi^-$) and the rate is determined by the pion's electromagnetic form factor:

$$\frac{\Gamma(\psi \to \pi^+\pi^-)}{\Gamma(\psi \to \mu^+\mu^-)} = \frac{1}{4}[F_\pi(s)]^2[1 + O(\alpha_s(s))],$$

where $s = (3.1GeV)^2$. Taking $F_\pi(s) \simeq (1 - s/m_\rho^2)^{-1}$ gives a rate $\Gamma(\psi \to \pi^+\pi^-) \sim 0.0011\ \Gamma(\psi \to \mu^+\mu^-)$, which compares well with the measured ratio 0.0015(7). This indicates that there is indeed little asymmetry in the pion's wave function.

The same analysis applied to $\psi \to K^+K^-$ suggests that the kaon's wave function is nearly symmetric about $x = \frac{1}{2}$. The ratio $\Gamma(\psi \to K^+K^-)/\Gamma(\psi \to \pi^+\pi^-)$ is 2 ± 1, which agrees with the ratio $(f_K/f_\pi)^4 \sim 2$ expected if π and K have similar quark distribution amplitudes. This conclusion is further supported by measurements of $\psi \to K_L K_S$ which vanishes completely if the K distribution amplitudes are symmetric; experimentally the limit is $\Gamma(\psi \to K_L K_S)/\Gamma(\psi \to K^+K^-) \lesssim \frac{1}{2}$.

It is important to test these PQCD and QCD sum rule predictions for the whole array of baryon pairs at both the J/ψ and ψ'. These decays give a direct measurement on the relative normalization of moments of the baryon distribution amplitudes. A particularly interesting quantity is the ratio $\Gamma(J/\psi \to p\bar{p})/\Gamma(J/\psi \to n\bar{n})$. Including the electromagnetic one-photon intermediate state contribution, one then obtains the prediction $\Gamma(J/\psi \to p\bar{p})/\Gamma(J/\psi \to n\bar{n}) = 1.16$. The present measurements[24] give $BR(J/\psi \to p\bar{p}) = 0.22 \pm 0.02\%$ and $BR(J/\psi \to n\bar{n}) = 0.18 \pm 0.09\%$. An important part of the QCD prediction is the electromagnetic decay amplitude controlled by the ratio of time-like form factors near the J/ψ. Using the QCD sum rule distribution amplitudes obtained by Chernyak and Zhitnitskii, one predicts

$$M^4_{J/\psi} G^p_M (Q^2 = M^2_{J/\psi}) = 1.1 \ GeV^4$$

$$M^4_{J/\psi} G^n_M (Q^2 = M^2_{J/\psi}) = -0.55 \ GeV^4,$$

which can be directly checked by measurements off resonance.

7. The π-ρ Puzzle

We have emphasized that a central prediction of perturbative QCD for exclusive processes is hadron helicity conservation: to leading order in $1/Q$, the total helicity of hadrons in the initial state must equal the total helicity of hadrons in the final state. This selection rule is independent of any photon or lepton spin appearing in the process. The result follows from (a) neglecting quark mass terms, (b) the vector coupling of gauge particles, and (c) the dominance of valence Fock states with zero angular momentum projection.[8] The result is true in each order of perturbation theory in α_s.

Hadron helicity conservation appears relevant to a puzzling anomaly in the exclusive decays J/ψ and $\psi' \to \rho\pi, K^*\overline{K}$ and possibly other Vector-Pseudoscalar (VP) combinations. One expects the J/ψ and ψ' mesons to decay to hadrons via three gluons or, occasionally, via a single direct photon. In either case the decay proceeds via $|\Psi(0)|^2$, where $\Psi(0)$ is the wave function at the origin in the nonrelativistic quark model for $c\bar{c}$. Thus it is reasonable to expect on the basis of perturbative QCD that for any final hadronic state h that the branching fractions scale like the branching fractions into e^+e^-:

$$Q_h \equiv \frac{B(\psi' \to h)}{B(J/\psi \to h)} \cong Q_{e^+e^-}$$

Usually this is true, as is well documented in Ref. 27 for $p\bar{p}\pi^0$, $2\pi^+2\pi^-\pi^0$, $\pi^+\pi^-\omega$, and $3\pi^+3\pi^-\pi^0$, hadronic channels. The startling exceptions occur for $\rho\pi$ and $K^*\overline{K}$ where the present experimental limits[27] are $Q_{\rho\pi} < 0.0063$ and $Q_{K^*\overline{K}} < 0.0027$.

Perturbative QCD quark helicity conservation implies[8] $Q_{\rho\pi} \equiv [B(\psi' \to \rho\pi)/B(J/\psi \to \rho\pi)] \leq Q_{e^+e^-}[M_{J/\psi}/M_{\psi'}]^6./t$ This result includes a form factor suppression proportional to $[M_{J/\psi}/M_{\psi'}]^4$ and an additional two powers of the mass ratio due to helicity flip. However, this suppression is not nearly large enough to account for the data.

From the standpoint of perturbative QCD, the observed suppression of $\psi' \to V\ P$ is to be expected; it is the J/ψ that is anomalous.[28] The ψ' obeys the perturbative QCD theorem that total hadron helicity is conserved in high-momentum transfer exclusive processes. The general validity of the QCD helicity conservation theorem at charmonium energies is of course open to question. An alternative model[29] based on nonperturbative exponential vertex functions, has recently been proposed to account for the anomalous exclusive decays of the J/ψ. However, helicity conservation has received important confirmation in $J/\psi \to p\bar{p}$ where the angular distribution is known experimentally to follow $[1+\cos^2\theta]$ rather than $\sin^2\theta$ for helicity flip, so the decays $J/\psi \to \pi\rho$, and $K\bar{K}$ seem truly exceptional.

The helicity conservation theorem follows from the assumption of short-range point-like interactions among the constituents in a hard subprocess. One way in which the theorem might fail for $J/\psi \to$ gluons $\to \pi\rho$ is if the intermediate gluons resonate to form a gluonium state $\mathcal{O}$. (See Fig. 8.) If such a state exists, has a mass near that of the J/ψ, and is relatively stable, then the subprocess for $J/\psi \to \pi\rho$ occurs over large distances and the helicity conservation theorem need no longer apply. This would also explain why the J/ψ decays into $\pi\rho$ and not the ψ'.

Tuan, Lepage, and I[28] have thus proposed, following Hou and Soni,[30] that the enhancement of $J/\psi \to K^*\overline{K}$ and $J/\psi \to \rho\pi$ decay modes is caused by a quantum mechanical mixing of the J/ψ with a $J^{PC} = 1^{--}$ vector gluonium state $\mathcal{O}$ which causes the breakdown of the QCD helicity theorem. The decay width for $J/\psi \to \rho\pi(K^*\overline{K})$ via the sequence $J/\psi \to \mathcal{O} \to \rho\pi(K^*\overline{K})$ must be substantially larger

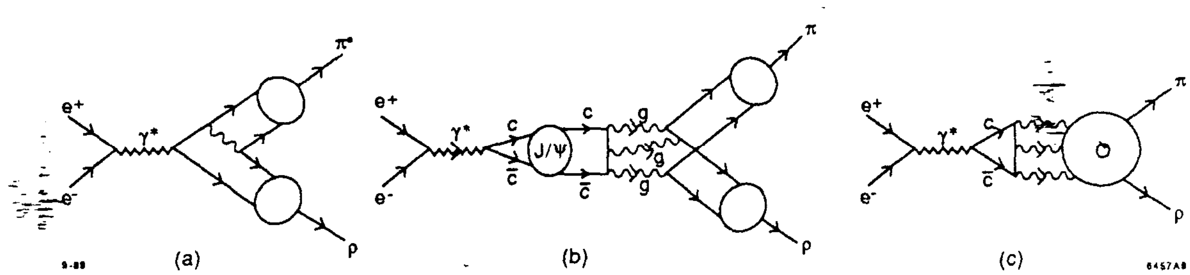

Figure 8. Illustration of QCD contributions for $J/\psi \to \rho\pi$. A non-perturbative contribution due to a gluonium resonance is shown in (c).

than the decay width for the (non-pole) continuum process $J/\psi \to 3$ gluons $\to \rho\pi(K^*\overline{K})$. In the other channels (such as $p\overline{p}, p\overline{p}\pi^0, 2\pi^+2\pi^-\pi^0$, etc.), the branching ratios of the $\mathcal{O}$ must be so small that the continuum contribution governed by the QCD theorem dominates over that of the $\mathcal{O}$ pole. For the case of the ψ' the contribution of the $\mathcal{O}$ pole must always be inappreciable in comparison with the continuum process where the QCD theorem holds. The experimental limits on $Q_{\rho\pi}$ and $Q_{K^*\overline{K}}$ are now substantially more stringent than when Hou and Soni made their estimates of $M_{\mathcal{O}}$, $\Gamma_{\mathcal{O}\to\rho\pi}$ and $\Gamma_{\mathcal{O}\to K^*\overline{K}}$ in 1982.

A gluonium state of this type was first postulated by Freund and Nambu[31] based on OZI dynamics soon after the discovery of the J/ψ and ψ' mesons. In fact, Freund and Nambu predicted that the $\mathcal{O}$ would decay primarily into $\rho\pi$ and $K^*\overline{K}$, with severe suppression of decays into other modes like e^+e^- as required for the solution of the puzzle.

Branching fractions for final states h which can proceed only through the intermediate gluonium state have the ratio:

$$Q_h = Q_{e^+e^-}\frac{(M_{J/\psi} - M_{\mathcal{O}})^2 + \frac{1}{4}\,\Gamma_0^2}{(M_{\psi'} - M_{\mathcal{O}})^2 + \frac{1}{4}\,\Gamma_0^2} \; .$$

It is assumed that the coupling of the J/ψ and ψ' to the gluonium state scales as the e^+e^- coupling. The value of Q_h is small if the $\mathcal{O}$ is close in mass to the J/ψ. Thus one requires $(M_{J/\psi} - M_{\mathcal{O}})^2 + \frac{1}{4}\,\Gamma_{\mathcal{O}}^2 \lesssim 2.6\; Q_h\;\mathrm{GeV}^2$. The experimental limit for $Q_{K^*\overline{K}}$ then implies $\left[(M_{J/\psi} - M_{\mathcal{O}})^2 + \frac{1}{4}\,\Gamma_{\mathcal{O}}^2\right]^{1/2} \lesssim 80$ MeV. This implies $\mid M_{J/\psi} - M_{\mathcal{O}} \mid < 80$ MeV and $\Gamma_{\mathcal{O}} < 160$ MeV. Typical allowed values are $M_{\mathcal{O}} = 3.0$ GeV, $\Gamma_{\mathcal{O}} = 140$ MeV or $M_{\mathcal{O}} = 3.15$ GeV, $\Gamma_{\mathcal{O}} = 140$ MeV. Notice that the gluonium state could be either lighter or heavier than the J/ψ. The branching ratio of the $\mathcal{O}$ into a given channel must exceed that of the J/ψ.

It is not necessarily obvious that a $J^{PC} = 1^{--}$ gluonium state with these parameters would necessarily have been found in experiments to date. One must remember that though $\mathcal{O} \to \rho\pi$ and $\mathcal{O} \to K^*\overline{K}$ are important modes of decay, at a mass of order 3.1 GeV many other modes (albeit less important) are available. Hence, a total width $\Gamma_{\mathcal{O}} \cong 100$ to 150 MeV is quite conceivable. Because of the proximity of $M_{\mathcal{O}}$ to $M_{J/\psi}$, the most important signatures for an $\mathcal{O}$ search via exclusive modes $J/\psi \to K^*\overline{K}h$, $J/\psi \to \rho\pi h$; $h = \pi\pi, \eta, \eta'$, are no longer available by phase-space considerations. However, the search could still be carried out using $\psi' \to K^*\overline{K}h$, $\psi' \to \rho\pi h$; with $h = \pi\pi$, and η. Another way to search for $\mathcal{O}$ in particular, and the three-gluon bound states in general, is via the inclusive reaction $\psi' \to (\pi\pi) + X$, where the $\pi\pi$ pair is an iso-singlet. The three-gluon bound states such as $\mathcal{O}$ should show up as peaks in the missing mass (*i.e.* mass of X) distribution.

The most direct way to search for the $\mathcal{O}$ is to scan $\overline{p}p$ or e^+e^- annihilation at $\sqrt{s}$ within ~ 100 MeV of the J/ψ, triggering on vector/pseudoscalar decays such as $\pi\rho$ or $\overline{K}K^*$.

The fact that the $\rho\pi$ and $K^*\overline{K}$ channels are strongly suppressed in ψ' decays but not in J/ψ decays clearly implies dynamics beyond the standard charmonium analysis. The hypothesis of a three-gluon state $\mathcal{O}$ with mass within $\cong 100$ MeV of the J/ψ mass provides a natural, perhaps even compelling, explanation of this anomaly. If this description is correct, then the ψ' and J/ψ hadronic decays not only confirm hadron helicity conservation (at the ψ' momentum scale), but they also provide a signal for bound gluonic matter in QCD.

A major problem, however, for the gluonium explanation of the $\rho\pi$ puzzle, is the relatively large decay rate recently reported for $J/\psi \to \omega\pi^0$. The published branching ratio is $0.048 \pm 0.007\%$ approximately three times larger than the $\pi^+\pi^-$ rate. Both of these $I = 1$ decays are evidently due to electromagnetic decays, but there is no sign of suppression due to hadron helicity conservation. One possibility is that there are additional $q\bar{q}g$ $I = 1$ resonances in the 3 GeV mass range which contribute to the $\omega\pi$ channel. In any event it will be very important to compare these branching ratios at the ψ' and off resonance.

8. Time-like Compton Processes

The high luminosity of a Tau Charm factory can allow the study of the basic Compton amplitude $M(\gamma^* \to \pi^+\pi^-\gamma)$ and the related Compton processes. The interference of this amplitude with contributions from diagrams where the photon is emitted from the initial electron or positron will produce a large front-back asymmetry in the $e^+e^- \to \pi^+\pi^-\gamma$ process. (See Fig. 9.) We can estimate

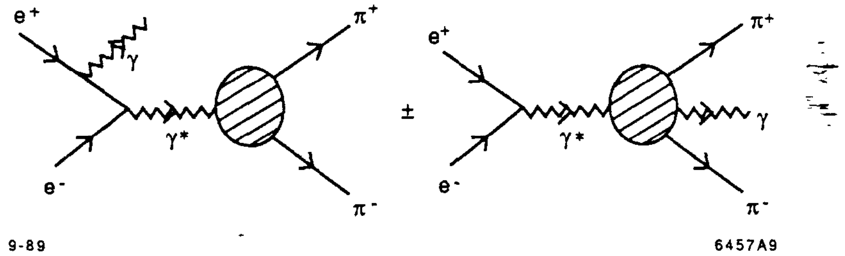

Figure 9. Interfering coherent amplitudes contributing to $e^+e^- \rightarrow \gamma\pi^+\pi^-$. This process measures the crossed pion Compton amplitude.

the event rate from $R(e^+e^- \rightarrow \pi^+\pi^-\gamma) \sim (\alpha/\pi)F_\pi^2(Q^2) \sim 10^{-4}$ to 10^{-5} which corresponds to 10^4 to 10^3 events per year at $10^{33} cm^{-2} sec^{-1}$ luminosity.

The Compton amplitude on a pion has thus far been studied only in the $\gamma\gamma \rightarrow \pi^+\pi^-$ reaction. The available Mark II and TPC/$\gamma\gamma$ data is in reasonable agreement with the leading twist QCD predictions. The QCD analysis predicts simple crossing of the large-angle $\gamma\gamma \rightarrow \pi^+\pi^-$ amplitude to the $\gamma^* \rightarrow \pi^+\pi^-\gamma$ amplitude. Extensive predictions are also now available for off-shell photons using PQCD factorization. A critical feature of the predictions is the presence of local two-photon couplings which lead to a dependence on photon mass Q much less severe than that predicted by vector meson dominance.

9. Multi-Hadron Production

A high luminosity e^+e^- facility could be used for the study of four-baryon exclusive final states and the search for new types of di-baryon states such as the H, the postulated $\Lambda\Lambda$ resonance suggested by Jaffe and others. (See Fig. 10.)

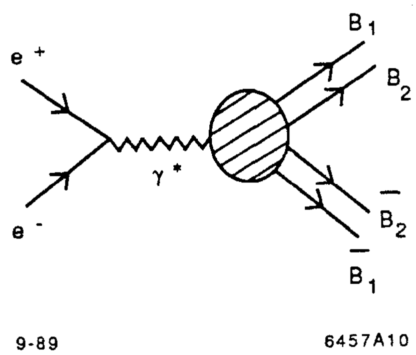

Figure 10. Production of four-baryon states in e^+e^- annihilation.

Dimensional counting predicts that the cross section for the production of N_M mesons, N_B baryons and $N_{\overline{B}}$ baryons at different fixed center of mass solid angle $\Delta\Omega$ scales as $\Delta\sigma \propto s^{-2-N_M-2N_B-2N_{\overline{B}}}$. Thus we can estimate $R_{e^+e^-\to B_1B_2\bar{B}_1\bar{B}_2} \sim |F_{B_1}(Q^2/4)F_{B_2}(Q^2/4)|^2$. The argument of the baryon form factor is $Q^2/4$ since each baryon is produced with half the available momentum. At $s = Q^2 = 16\ GeV^2$, this corresponds to an annihilation ratio $R \sim 10^{-4}$. The production of the $np\bar{d}$ nuclear final state is further reduced by the probability that the nucleons fuse in a restricted phase space, and thus is suppressed by an additional power of $1/Q^2$.

The above estimates are consistent with the "reduced amplitude" formalism for exclusive nuclear processes which has been successful predicting the scaling behavior of the deuteron form factor and the deuteron photo-disintegration cross section at fixed $\theta_{\rm cm}$.

One can thus envision having sufficient luminosity at a Tau-Charm factory to search for the H di-lambda in the missing mass distribution in the reaction $e^+e^- \to \bar{\Lambda}\bar{\Lambda}X$. This method can be extended to search for exotic resonances in the Λp, Σp di-baryon systems. The rate for four-meson exclusive channels is considerably larger, and affords the possibility of studying the interactions of di-meson systems such as K^+K^+. In each case the study of multi-hadron exclusive channels can allow the study of the scattering length and range of hadron-hadron final state interactions.

10. Heavy Quark Exclusive States and Form Factor Zeros in QCD

The exclusive pair production of heavy hadrons $|Q_1\overline{Q}_2\rangle$, $|Q_1Q_2Q_3\rangle$ consisting of higher generation quarks ($Q_i = t, b, c$, and possibly s) can be reliably predicted within the framework of perturbative QCD, since the required wavefunction input is essentially determined from nonrelativistic considerations.[32] The results can be applied to e^+e^- annihilation, $\gamma\gamma$ annihilation, and W and Z decay into higher generation pairs. The normalization, angular dependence and helicity structure can be predicted away from threshold, allowing a detailed study of the basic elements of heavy quark hadronization.

In the case of the Tau-Charm factory, it is interesting to test the predictions of QCD factorization for time-like meson form factors for the production of heavy meson pairs, such as $e^+e^- \to D\bar{D}$ and $e^+e^- \to D_s\bar{D}_s$.

A particularly striking feature of the QCD predictions is the existence of a zero in the form factor and e^+e^- annihilation cross section for zero-helicity hadron pair production close to the specific time-like value $q^2/4M_H^2 = m_h/2m_\ell$ where m_h and m_ℓ are the heavier and lighter quark masses, respectively. This zero reflects the destructive interference between the spin-dependent and spin-independent (Coulomb exchange) couplings of the exchanged gluon shown in Fig. 11; it is thus a novel feature of the gauge theory. In fact, all pseudoscalar meson form factors are predicted in QCD to reverse sign from space-like to time-like asymptotic momentum transfer because of their essentially monopole form. For $m_h > 2m_\ell$ the form factor zero occurs in the physical region.

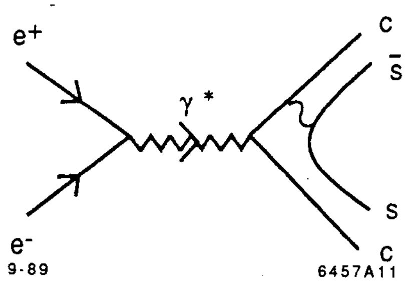

Figure 11. Illustration of the dominant hard scattering diagram for $D_s \bar{D}_s$ pair production in QCD.

In the case of $e^+e^- \rightarrow D_s\bar{D}_s$ the amplitude vanishes and changes sign at $Q^2/4M^2_{D_s} \approx m_c/2m_s$. Since background terms are expected to be monotonic, an amplitude zero must occur somewhere above threshold in $e^+e^- \rightarrow D_s\bar{D}_s$. (See Fig. 12.) The absolute rate near threshold for this process depends on the wavefunction parameters, particularly the mean square relative velocity of the constituents. We estimate $R(D_s\bar{D}_s) < 10^{-4}$, a rate measurable at a high-luminosity Tau-Charm factory.

To leading order in $1/q^2$, the production amplitude for hadron pair production is given by the factorized form

$$M_{H\overline{H}} = \int [dx_i] \int [dy_j]\, \phi^\dagger_H(x_i, \widetilde{q}^2)\, \phi^\dagger_{\overline{H}}(y_j, \widetilde{q}^2)\, T_H(x_i, y_j; \widetilde{q}^2, \theta_{CM})$$

where $[dx_i] = \delta\left(\sum_{k=1}^n x_k - 1\right) \prod_{k=1}^n dx_k$ and $n = 2, 3$ is the number of quarks in the valence Fock state. The scale $\widetilde{q}^2$ is set from higher order calculations, but it reflects the minimum momentum transfer in the process. The main dynamical dependence of the form factor is controlled by the hard scattering amplitude T_H which is computed by replacing each hadron by collinear constituents $P_i^\mu = x_i P_H^\mu$. Since the collinear divergences are summed in ϕ_H, T_H can be systematically computed as a perturbation expansion in $\alpha_s(q^2)$.

The distribution amplitude required for heavy hadron production $\phi_H(x_i, q^2)$ is computed as an integral of the valence light-cone Fock wavefunction up to the

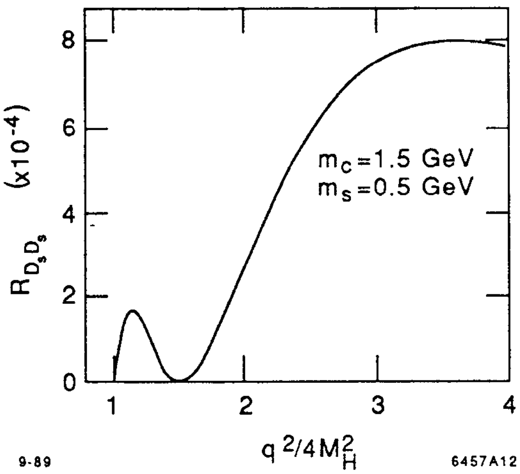

Figure 12. Perturbative QCD prediction[32] for $R(e^+e^- \to D_s\bar{D}_s)$. The normalization depends on assumptions for the D_s wavefunction.

scale Q^2. For the case of heavy quark bound states, one can assume that the constituents are sufficiently non-relativistic that gluon emission, higher Fock states, and retardation of the effective potential can be neglected. The analysis of Section 2 is thus relevant. The quark distributions are then controlled by a simple nonrelativistic wavefunction, which can be taken in the model form:

$$\psi_M(x_i, \vec{k}_{\perp i}) = \frac{C}{x_1^2 x_2^2 \left[M_H^2 - \frac{\vec{k}_{\perp 1}^2 + m_1^2}{x_1} - \frac{\vec{k}_{\perp 2}^2 + m_2^2}{x_2} \right]^2}$$

This form is chosen since it coincides with the usual Schrödinger- Coulomb wavefunction in the nonrelativistic limit for hydrogenic atoms and has the correct large momentum behavior induced from the spin- independent gluon couplings. The wavefunction is peaked at the mass ratio $x_i = m_i/M_H$:

$$\left(x_i - \frac{m_i}{M_H} \right)^2 \sim \frac{\langle k_z^2 \rangle}{M_H^2}$$

where $\langle k_z^2 \rangle$ is evaluated in the rest frame. Normalizing the wavefunction to unit probability gives

$$C^2 = 128\pi \left(\langle v^2 \rangle \right)^{5/2} m_r^5 (m_1 + m_2)$$

where $\langle v^2 \rangle$ is the mean square relative velocity and $m_r = m_1 m_2/(m_1 + m_2)$ is the reduced mass. The corresponding distribution amplitude is

$$\phi(x_i) = \frac{C}{16\pi^2} \frac{1}{[x_1 x_2 M_H^2 - x_2 m_1^2 - x_1 m_2^2]}$$

$$\cong \frac{1}{\sqrt{2\pi}} \frac{\gamma^{3/2}}{M_H^{1/2}} \delta\left(x_1 - \frac{m_1}{m_1 + m_2}\right) .$$

It is easy to see from the structure of T_H for $e^+e^- \to M\overline{M}$ that the spectator quark pair is produced with momentum transfer squared $q^2 x_s y_s = 4m_s^2$. Thus heavy hadron pair production is dominated by diagrams in which the primary coupling of the virtual photon is to the heavier quark pair. The perturbative predictions are thus expected to be accurate even near threshold to leading order in $\alpha_s(4m_\ell^2)$ where m_ℓ is the mass of lighter quark in the meson.

The leading order e^+e^- production helicity amplitudes for higher generation meson ($\lambda = 0, \pm 1$) and baryon ($\lambda = \pm 1/2, \pm 3/2$) pairs are computed in Ref. 32 as a function of q^2 and the quark masses. The analysis is simplified by using the peaked form of the distribution amplitude, Eq. (6). In the case of meson pairs the (unpolarized) e^+e^- annihilation cross section has the general form*

$$4\pi \frac{d\sigma}{d\Omega}(e^+e^- \to M_\lambda \overline{M}_{\overline{\lambda}}) = \frac{3}{4} \beta \sigma_{e^+e^- \to \mu^+\mu^-} \left[\frac{1}{2} \beta^2 \sin^2\theta \right.$$

$$\times \left[|F_{0,0}(q^2)|^2 + \frac{1}{(1-\beta^2)^2} \left\{ (3 - 2\beta^2 + 3\beta^4) |F_{1,1}(q^2)|^2 \right. \right.$$

$$\left. \left. - 4(1+\beta^2)\, \mathrm{Re}\,(F_{1,1}(q^2) F_{0,1}^*(q^2)) + 4|F_{0,1}(q^2)|^2 \right\} \right]$$

$$\left. + \frac{3\beta^2}{2(1-\beta^2)} (1 + \cos^2\theta) |F_{0,1}(q^2)|^2 \right]$$

where $q^2 = s = 4M_H^2 \bar{q}^2$ and the meson velocity is $\beta = 1 - \frac{4M_H^2}{q^2}$. The production

---

* $F_{\lambda\overline{\lambda}}(q^2)$ is the form factor for the production of two mesons which have both spin and helicity (Z-component of spin) as λ and $\overline{\lambda}$ respectively. There are two Lorentz and gauge invariant form factors of vector pair production. However, one of them turns out to be the same as the form factor of pseudoscalar plus vector production multiplied by M_H. Therefore the differential cross section for the production of two mesons with spin 0 or 1 can be represented in terms of three independent form factors.

form factors have the general form

$$F_{\lambda\bar{\lambda}} = \frac{\langle v^2 \rangle^2}{(\bar{q}^2)^2} \left(A_{\lambda\bar{\lambda}} + \bar{q}^2 B_{\lambda\bar{\lambda}} \right)$$

where A and B reflect the Coulomb-like and transverse gluon couplings, respectively. The results to leading order in α_s are given in Ref. 32. In general A and B have a slow logarithmic dependence due to the q^2-evolution of the distribution amplitudes. The form factor zero for the case of pseudoscalar pair production reflects the numerator structure of the T_H amplitude.

$$\text{Numerator} \sim e_1 \left(\bar{q}^2 - \frac{m_1^2}{4M_H^2} \frac{1}{x_2 y_1} - \frac{m_2^2}{4M_H^2} \frac{x_1}{x_2^2 y_2} \right)$$

For the peaked wavefunction,

$$F_{0,0}^M(q^2) \propto \frac{1}{(\bar{q}^2)^2} \left\{ e_1 \left(\bar{q}^2 - \frac{m_1}{2m_2} \right) + e_2 \left(\bar{q}^2 - \frac{m_2}{2m_1} \right) \frac{m_2^2}{m_1^2} \right\}$$

If m_1 is much greater than m_2 then the e_1 is dominant and changes sign at $q^2/4M_H^2 = m_1/2m_2$. The contribution of the e_2 term and higher order contributions are small and nearly constant in the region where the e_1 term changes sign; such contributions can displace slightly but not remove the form factor zero. These results also hold in quantum electrodynamics; *e.g.* , pair production of muonium $(\mu - e)$ atoms in e_+e_- annihilation. Gauge theory predicts a zero at $\bar{q}^2 = m_\mu/2m_e$.

These explicit results for form factors also show that the onset of the leading power-law scaling of a form factor is controlled by the ratio of the A and B terms; *i.e.* , when the transverse contributions exceed the Coulomb mass-dominated contributions. The Coulomb contribution to the form factor can also be computed directly from the convolution of the initial and final wavefunctions. Thus, contrary to the claim of Ref. 12 there are no extra factors of $\alpha_s(q^2)$ which suppress the "hard" versus nonperturbative contributions.

The form factors for the heavy hadrons are normalized by the constraint that the Coulomb contribution to the form factor equals the total hadronic charge at $q^2 = 0$. Further, by the correspondence principle, the form factor should agree with the standard non-relativistic calculation at small momentum transfer. All of these constraints are satisfied by the form

$$F^M_{0,0}(q^2) = e_1 \, \frac{16\gamma^4}{(q^2+\gamma^2)^2} \left(\frac{M_H^2}{m_2^2}\right)^2 \left(1 - \frac{q^2}{4M_H^2}\,\frac{2m_2}{m_1}\right) + 1 \leftrightarrow 2 .$$

At large q^2 the form factor can also be written as

$$F^M_{(0,0)} = e_1 \, \frac{16\pi\alpha_s f_M^2}{9q^2} \left(\frac{M_H^2}{m_2^2}\right) + (1 \leftrightarrow 2) \, , \quad \frac{f_M}{2\sqrt{3}} = \int_0^1 dx \, \phi(x, Q)$$

where $f_M = (6\gamma^3/\pi M_H)^{1/2}$ is the meson decay constant. Detailed results for $F\bar{F}$ and $B_c\bar{B}_c$ production are give in Ref. 32.

At low relative velocity of the hadron pair one also expects resonance contributions to the form factors. For these heavy systems such resonances could be related to $qq\bar{q}\bar{q}$ bound states. From Watson's theorem, one expects any resonance structure to introduce a final-state phase factor, but not destroy the zero of the underlying QCD prediction.

Analogous calculations of the baryon form factor, retaining the constituent mass structure have also been done. The numerator structure for spin 1/2 baryons has the form

$$A + B\bar{q}^2 + c\bar{q}^4 \, .$$

Thus it is possible to have two form factor zeros; *e.g.* , at space-like and time-like values of q^2.

Although the measurements are difficult and require large luminosity, the observation of the striking zero structure predicted by QCD would provide a unique test of the theory and its applicability to exclusive processes. The onset of leading power behavior is controlled simply by the mass parameters of the theory.

11. Exclusive $\gamma\gamma$ Reactions

A number of interesting $\gamma\gamma$ annihilation processes could be studied advantageously at a Tau-Charm factory. Such two-photon reactions have a number of unique features which are important for testing QCD:[33]

1. Any even charge conjugation hadronic state can be created in the annihilation of two photons—an initial state of minimum complexity. Because $\gamma\gamma$ annihilation is complete, there are no spectator hadrons to confuse resonance analyses. Thus, one has a clean environment for identifying the exotic color-singlet even C composites of quarks and gluons $|q\bar{q}>$, $|gg>$, $|ggg>$, $|q\bar{q}g>$,

$|qq\overline{qq}>, \ldots$ which are expected to be present in the few GeV mass range. (Because of mixing, the actual mass eigenstates of QCD may be complicated admixtures of the various Fock components.)

2. The mass and polarization of each of the incident virtual photons can be continuously varied, allowing highly detailed tests of theory. Because a spin-one state cannot couple to two on-shell photons, a $J = 1$ resonance can be uniquely identified by the onset of its production with increasing photon mass.[34]

3. Two-photon physics plays an especially important role in probing dynamical mechanisms. In the low momentum transfer domain, $\gamma\gamma$ reactions such as the total annihilation cross section and exclusive vector meson pair production can give important insights into the nature of diffractive reactions in QCD. Photons in QCD couple directly to the quark currents at any resolution scale. Predictions for high momentum transfer $\gamma\gamma$ reactions, including the photon structure functions, $F_2^\gamma(x, Q^2)$ and $F_L^\gamma(x, Q^2)$, high p_T jet production, and exclusive channels are thus much more specific than corresponding hadron-induced reactions. The point-like coupling of the annihilating photons leads to a host of special features which differ markedly with predictions based on vector meson dominance models.

4. Exclusive $\gamma\gamma$ processes provide a window for viewing the wavefunctions of hadrons in terms of their quark and gluon degrees of freedom. In the case of $\gamma\gamma$ annihilation into hadron pairs, the angular distribution of the production cross section directly reflects the shape of the distribution amplitude (valence wavefunction) of each hadron.

A simple, but still very important example,[6] is the Q^2-dependence of the reaction $\gamma^*\gamma \to M$ where M is a pseudoscalar meson such as the η. The invariant amplitude contains only one form factor:

$$M_{\mu\nu} = \epsilon_{\mu\nu\sigma\tau} p_\eta^\sigma q^\tau F_{\gamma\eta}(Q^2) \ .$$

It is easy to see from power counting at large Q^2 that the dominant amplitude (in light-cone gauge) gives $F_{\gamma\eta}(Q^2) \sim 1/Q^2$ and arises from diagrams which have the minimum path carrying Q^2: *i.e.* , diagrams in which there is only a single quark propagator between the two photons. The coefficient of $1/Q^2$ involves only the two-particle $q\overline{q}$ distribution amplitude $\phi(x, Q)$, which evolves logarithmically on Q. Higher particle number Fock states give higher power-law falloff contributions to the exclusive amplitude.

The TPC/$\gamma\gamma$ data[36] shown in Fig. 13 are in striking agreement with the predicted QCD power: a fit to the data gives $F_{\gamma\eta}(Q^2) \sim (1/Q^2)^n$ with $n = 1.05 \pm$

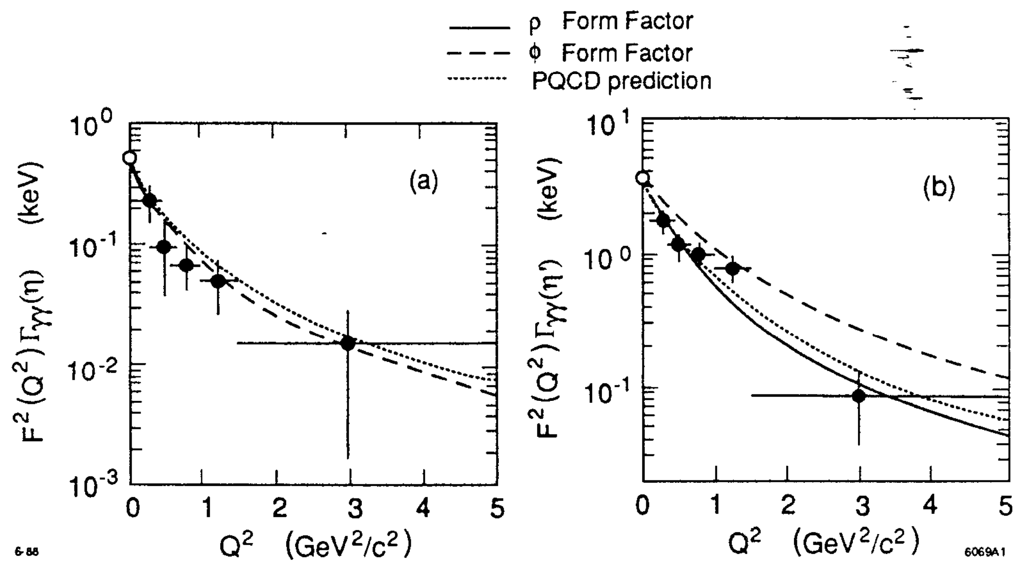

Figure 13. Comparison of TPC/$\gamma\gamma$ data[36] for the $\gamma-\eta$ and $\gamma-\eta'$ transition form factors with the QCD leading twist prediction of Ref. 35. The VMD predictions are also shown.

0.15. Data for the η' from Pluto and the TPC/$\gamma\gamma$ experiments give similar results, consistent with scale-free behavior of the QCD quark propagator and the point coupling to the quark current for both the real and virtual photons. In the case of deep inelastic lepton scattering, the observation of Bjorken scaling tests the same scaling of the quark Compton amplitude when both photons are virtual.

The QCD power law prediction, $F_{\gamma\eta}(Q^2) \sim 1/Q^2$, is consistent with dimensional counting[2] and also emerges from current algebra arguments (when both photons are very virtual).[37] On the other hand, the $1/Q^2$ falloff is also expected in vector meson dominance models. The QCD and VDM predictions can be readily discriminated by studying $\gamma^*\gamma^* \to \eta$. In VMD one expects a product of form factors; in QCD, the fall-off of the amplitude is still $1/Q^2$ where Q^2 is a linear combination of Q_1^2 and Q_2^2. It is clearly very important to test this essential feature of QCD.

We also note that photon-photon collisions provide a way to measure the running coupling constant in an exclusive channel, independent of the form of hadronic distribution amplitudes.[35] The photon-meson transition form factors $F_{\gamma-M}(Q^2)$, $M = \pi^0, \eta^0, f$, etc., are measurable in tagged $e\gamma \to e'M$ reactions. QCD predicts

$$\alpha_s(Q^2) = \frac{1}{4\pi} \frac{F_\pi(Q^2)}{Q^2|F_{\pi\gamma}(Q^2)|^2}$$

where to leading order the pion distribution amplitude enters both numerator and denominator in the same manner.

Exclusive two-body processes $\gamma\gamma \to H\overline{H}$ at large $s = W_{\gamma\gamma}^2 = (q_1 + q_2)^2$ and fixed $\theta_{\rm cm}^{\gamma\gamma}$ provide a particularly important laboratory for testing QCD, since the large momentum-transfer behavior, helicity structure, and often even the absolute normalization can be rigorously predicted.[35,38] The angular dependence of some of the $\gamma\gamma \to H\overline{H}$ cross sections reflects the shape of the hadron distribution amplitudes $\phi_H(x_i, Q)$. The $\gamma_\lambda\gamma_{\lambda'} \to H\overline{H}$ amplitude can be written as a factorized form

$$\mathcal{M}_{\lambda\lambda'}(W_{\gamma\gamma}, \theta_{\rm cm}) = \int_0^1 [dy_i]\, \phi_H^*(x_i, Q)\, \phi_{\overline{H}}^*(y_i, Q)\, T_{\lambda\lambda'}(x, y; W_{\gamma\gamma}, \theta_{\rm cm})$$

where $T_{\lambda\lambda'}$ is the hard scattering helicity amplitude. To leading order $T \propto \alpha(\alpha_S/W_{\gamma\gamma}^2)^n$ and $d\sigma/dt \sim W_{\gamma\gamma}^{-(2n+2)} f(\theta_{\rm cm})$ where $n = 1$ for meson and $n = 2$ for baryon pairs.

Lowest order predictions for pseudo-scalar and vector-meson pairs for each helicity amplitude are given in Ref. 35. In each case the helicities of the hadron pairs are equal and opposite to leading order in $1/W^2$. The normalization and angular dependence of the leading order predictions for $\gamma\gamma$ annihilation into charged meson pairs are almost model independent; *i.e.* , they are insensitive to the precise form of the meson distribution amplitude. If the meson distribution amplitudes is symmetric in x and $(1 - x)$, then the same quantity

$$\int_0^1 dx\, \frac{\phi_\pi(x, Q)}{(1 - x)}$$

controls the x-integration for both $F_\pi(Q^2)$ and to high accuracy $M(\gamma\gamma \to \pi^+\pi^-)$. Thus for charged pion pairs one obtains the relation:

$$\frac{\frac{d\sigma}{dt}(\gamma\gamma \to \pi^+\pi^-)}{\frac{d\sigma}{dt}(\gamma\gamma \to \mu^+\mu^-)} \cong \frac{4|F_\pi(s)|^2}{1 - \cos^4\theta_{\rm cm}} .$$

Note that in the case of charged kaon pairs, the asymmetry of the distribution amplitude may give a small correction to this relation.

The scaling behavior, angular behavior, and normalization of the $\gamma\gamma$ exclusive pair production reactions are nontrivial predictions of QCD. Mark II meson pair data and PEP4/PEP9 data[39] for separated $\pi^+\pi^-$ and K^+K^- production in the range $1.6 < W_{\gamma\gamma} < 3.2$ GeV near 90° are in satisfactory agreement with the normalization and energy dependence predicted by QCD (see Fig. 14). In the case of $\pi^0\pi^0$ production, the cos θ_{cm} dependence of the cross section can be inverted to determine the x-dependence of the pion distribution amplitude.

The wavefunction of hadrons containing light and heavy quarks such as the K, D-meson are likely to be asymmetric due to the disparity of the quark masses. In a gauge theory one expects that the wavefunction is maximum when the quarks have zero relative velocity; this corresponds to $x_i \propto m_{i\perp}$ where $m_\perp^2 = k_\perp^2 + m^2$. An explicit model for the skewing of the meson distribution amplitudes based on QCD sum rules is given by Benyayoun and Chernyak.[40] These authors also apply their model to two-photon exclusive processes such as $\gamma\gamma \to K^+K^-$ and obtain some modification compared to the strictly symmetric distribution amplitudes. If the same conventions are used to label the quark lines, the calculations of Benyayoun and Chernyak are in complete agreement with those of Ref. 35.

The one-loop corrections to the hard scattering amplitude for meson pairs have been calculated by Nizic.[41] The QCD predictions for mesons containing admixtures of the $|gg\rangle$ Fock state is given by Atkinson, Sucher, and Tsokos.[38]

The perturbative QCD analysis has been extended to baryon-pair production in comprehensive analyses by Farrar, *et al.*[42,38] and by Gunion, *et al.*[38] Predictions are given for the "sideways" Compton process $\gamma\gamma \to p\bar{p}$, $\Delta\bar{\Delta}$ pair production, and the entire decuplet set of baryon pair states. The arduous calculation of 280 $\gamma\gamma \to qqq\overline{qqq}$ diagrams in T_H required for calculating $\gamma\gamma \to B\overline{B}$ is greatly simplified by using two-component spinor techniques. The doubly charged Δ pair is predicted to have a fairly small normalization. Experimentally such resonance pairs may be difficult to identify under the continuum background.

The normalization and angular distribution of the QCD predictions for proton-antiproton production depend in detail on the form of the nucleon distribution amplitude, and thus provide severe tests of the model form derived by Chernyak, Ogloblin, and Zhitnitskii[11] from QCD sum rules.

The region of applicability of the leading power-law predictions for $\gamma\gamma \to p\bar{p}$ requires that one be beyond resonance or threshold effects. It presumably is set by the scale where $Q^4 G_M(Q^2)$ is roughly constant; *i.e.* , $Q^2 > 3$ GeV^2. Measurements of baryon pairs at a Tau-Charm factory are thus probably too close to threshold for meaningful tests of the PQCD predictions.[44]

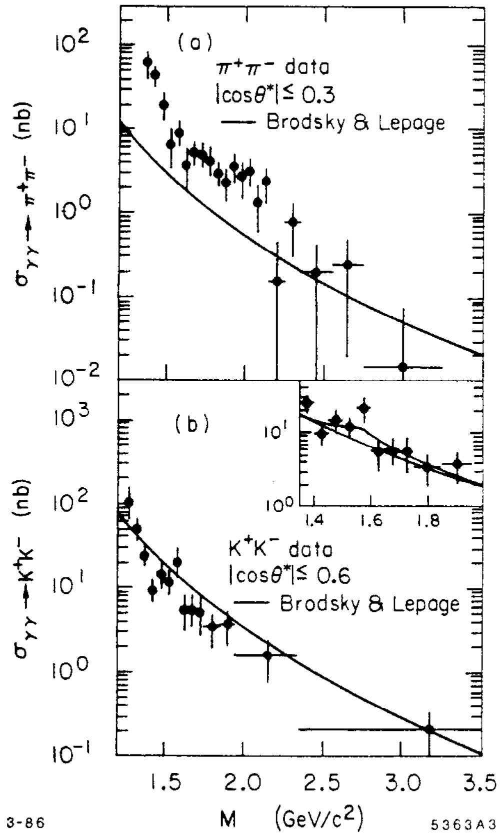

Figure 14. Comparison of $\gamma\gamma \to \pi^+\pi^-$ and $\gamma\gamma \to K^+K^-$ meson pair production data with the parameter-free perturbative QCD prediction of Ref. 35. The theory predicts the normalization and scaling of the cross sections. The data are from the TPC/$\gamma\gamma$ collaboration.[39]

The QCD predictions for $\gamma\gamma \to H\overline{H}$ can be extended to the case of one or two virtual photons, for measurements in which one or both electrons are tagged. Because of the direct coupling of the photons to the quarks, the Q_1^2 and Q_2^2 dependence of the $\gamma\gamma \to H\overline{H}$ amplitude for transversely polarized photons is minimal at W^2 large and fixed $\theta_{\rm cm}$, since the off-shell quark and gluon propagators in T_H already transfer hard momenta; *i.e.* , the 2γ coupling is effectively local for Q_1^2, $Q_2^2 \ll p_T^2$. The $\gamma^*\gamma^* \to \bar{B}B$ and $M\overline{M}$ amplitudes for off-shell photons have been calculated by Millers and Gunion.[38] In each case, the predictions show strong sensitivity to the form of the respective baryon and meson distribution amplitudes.

12. Higher Twist Effects

One of the most elusive topics in PQCD has been the unambiguous identification of higher-twist effects in inclusive reaction. A signal for a dynamical higher-twist amplitude has been seen in pion-induced Drell-Yan reactions, where a $1/Q^2$ component to the pion structure function $F_L^\pi(x_1, Q^2)$ coupling to longitudinal photons dominates the cross section at large x_1. In addition, a Rice–Fermilab experiment studying pion-induced di-jet production has found evidence for the directly-coupled pion higher-twist subprocess $\pi g \to q\bar{q}$ which has the unusual property that there is no jet of hadrons left in the beam direction.

In the case of inclusive quark jet fragmentation, $e^+e^- \to \pi X$, PQCD predicts analogous anomalous behavior in the jet distribution at large $z = E_\pi/Q$. In the analysis one must take into account the subprocess $\gamma^* \to \pi q\bar{q}$ illustrated in Fig. 15 where the pion is produced directly at short distances, in addition to the standard leading twist process where the pion is produced from jet fragmentation. The net result is a prediction at large z of the form

$$\frac{d\sigma(e^+e^- \to \pi X)}{dz\ dcos\theta} = A(1-z)^2(1+cos^2\theta) + B\frac{sin^2\theta}{Q^2}.$$

Although the corresponding B term has been observed in the Drell-Yan reaction, it has never been seen unambiguously in jet fragmentation. The lower energies of a Tau-Charm factory should be advantageous in identifying the $1/Q^2$ dependence of the direct pion contributions.

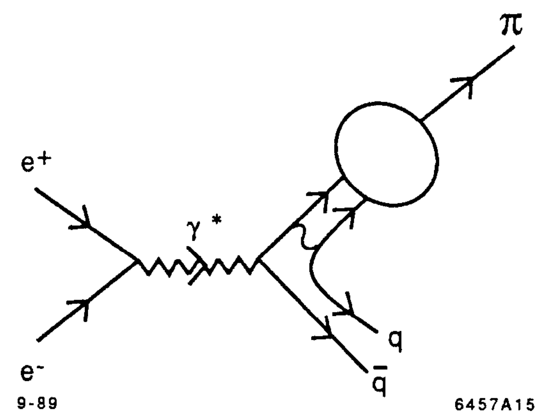

Figure 15. Higher-twist contribution to jet fragmentation in $e^+e^- \to \pi X$. The pion couples through its distribution amplitude $\phi_\pi(x, Q)$.

13. Tauonium and Threshold $\tau^+\tau^-$ Production

In principle, $J^P = 1^-$ QED bound states of $\tau^+\tau^-$ could be produced as very narrow resonances below threshold in e^+e^- annihilation.[45] Unfortunately the observation of even the lowest ortho-tauonium state at a measurable level would require much higher incident energy resolution then presently possible. The higher n excitations are suppressed by a factor $1/n^3$, so radiative decay signals would not be produced at a practical rate. Worse, the τ will decay weakly before radiative transitions can occur.

The continuum production of the $\tau^+\tau^-$ near threshold is strongly modified by final-state QED interactions.[46] The leading order correction to the Born term at threshold has the form $(1 + \alpha f(v))$ where $v = (1 - 4M_\tau^2/s)$ and

$$f(v) = \frac{\pi}{2v} - \frac{3+v}{4}\left(\frac{\pi}{2} - \frac{3}{4\pi}\right).$$

The singular factor in $1/v$ cancels the phase-space factor in the Born cross section, giving a non-zero rate for production at threshold. The analogous effect is well-known in QCD for threshold charm production, and has been taken into account in the duality formulas which relate charm hadron production to the mass of the charm quark.[47] It would be interesting to check the threshold production of $e^+e^- \to \tau^+\tau^-$ and verify this interesting feature of τ electrodynamics.

REFERENCES

1. For a general review of the PQCD analysis of exclusive processes and for further reference, see S. J. Brodsky and G. P. Lepage, SLAC-PUB-4947 (1989), to be published in "Perturbative Quantum Chromodynamics", ed. A. H. Mueller, World Scientific (1989).
2. S.J. Brodsky and G.R. Farrar, Phys. Rev. Lett. 31, 1153 (1973); Phys. Rev. D11, 1309 (1975). V. A. Matveev, R. M. Muradyan and A. V. Tavkheldize. Lett. Nuovo Cimento 7, 719 (1973).
3. A. H. Mueller, Phys. Rept. 73 237 (1981). J. Botts, and G. Sterman, Phys. Lett. B224 201 (1989); Stony Brook preprints ITP-SB-89-7, 44.
4. S. Dubnicka and E. Etim, Frascati preprint, LNF-89/013-PT, presented at Nucleon Structure Workshop: FENICE Experiment and Investigation of the Neutron Form-Factor, Frascati, Italy, October 1988.
5. S. J. Brodsky SLAC-PUB-4648, 8th Int. Workshop on Photon-Photon Collisions, Jerusalem, Israel, (1988.)
6. G. P. Lepage and S. J. Brodsky, Phys. Rev. D22, 2157 (1980); Phys. Lett. 87B, 359 (1979); Phys. Rev. Lett. 43, 545, 1625(E) (1979).
7. V. D. Burkert, CEBAF-PR-87-006.
8. S. J. Brodsky and G. P. Lepage, Phys. Rev. D24, 2848 (1981).
9. See also G. R. Farrar, presented to the Workshop on Quantum Chromodynamics at Santa Barbara, 1988.
10. V.L. Chernyak and A.R. Zhitnitskii, Phys. Rept. 112, 173 (1984). See also Xiao-Duang Xiang, Wang Xin-Nian, and Huang Tao, BIHEP-TH-84, 23 and 29 (1984).
11. V. L. Chernyak, A. A. Ogloblin and I. R. Zhitnitskii, Novosibirsk preprints INP 87–135,136, and references therein. See also Xiao-Duang Xiang, Wang Xin-Nian, and Huang Tao, BIHEP-TH-84, 23 and 29, 1984, and M. J. Lavelle, ICTP-84-85-12; Nucl. Phys. B260, 323 (1985).
12. N. Isgur and C.H. Llewellyn Smith, Phys. Rev. Lett. 52, 1080 (1984). G. P. Korchemskii, A. V. Radyushkin, Sov. J. Nucl. Phys. 45, 910 (1987) and refs. therein.
13. Z. Dziembowski, G. Farrar, H. Zhang, and L. Mankiewicz, contribution to the 12th Int. Conf. on Few Body Problems in Physics, Vancouver, 1989.
14. Z. Dziembowski and J. Franklin, contribution to the 12th Int. Conf. on Few Body Problems in Physics, Vancouver, 1989.

15. S. J. Brodsky, T. Huang, G. P. Lepage, published in the proceedings of the SLAC Summer Institute 1981:87.

16. C. Carlson and F. Gross, Phys. Rev. Lett. 53, 127 (1984); Phys. Rev. D36 2060 (1987).

17. O. C. Jacob and L. S. Kisslinger, Phys. Rev. Lett. 56, 225 (1986).

18. Z. Dziembowski and L. Mankiewicz, Phys. Rev. Lett. us 58 2175 (1987); Z. Dziembowski, Phys. Rev. D37 768, 778, 2030 (1988).

19. T. Huang, Q.-X. Shen, and P. Kroll, to be published.

20. R. G. Arnold *et al.* , Phys. Rev. Lett. 57, 174 (1986).

21. A. H. Mueller, Proc. XVII Recontre de Moriond (1982); S. J. Brodsky, Proc. XIII Int. Symp. on Multiparticle Dynamics, Volendam (1982). See also G. Bertsch, A. S. Goldhaber, and J. F. Gunion, Phys. Rev. Lett. 47, 297 (1981); G. R. Farrar, H. Liu, L. L. Frankfurt, M. J. Strikmann, Phys. Rev. Lett. 61, 686 (1988); A. H. Mueller, CU-TP-415, talk given at the DPF meeting, Stoors, Conn (1988), and CU-TP-412 talk given at the Workshop on Nuclear and Particle Physics on the Light-Cone, Los Alamos, (1988).

22. S. J. Brodsky and G. de Teramond, Phys. Rev. Lett. 60, 1924 (1988).

23. I. Peruzzi *et al.* , Phys. Rev. D17, 2901 (1978).

24. Unless otherwise noted, the data used here is from the compilation of the Particle Data Group, Phys. Lett. B, 204, 1988.

25. Arguments for the conservation of baryon chirality in large-momentum transfer processes have been given by B. L. Ioffe, Phys. Lett. 63B, 425 (1976). For some processes this rule leads to predictions which differ from the QCD results given here. The QCD helicity conservation rule also differs from the electroproduction helicity rules given in O. Nachtmann, Nucl. Phys. B115, 61 (1976).

26. See also A. I. Vainshtein and V. I. Zakharov, Phys. Lett. 72B, 368 (1978); G. R. Farrar and D. R. Jackson, Phys. Rev. Lett. 35, 1416 (1975); B. Ioffe, Ref. 25).

27. M. E. B. Franklin, Ph.D Thesis (1982), SLAC–254, UC–34d; M. E. B. Franklin *et al.* , Phys. Rev. Lett. 51, 963 (1983); G. Trilling, in the Proceedings of the Twenty-First International Conference on High Energy Physics, Paris, July 26–31, 1982; E. Bloom, ibid.

28. S. J. Brodsky, G. P. Lepage and San Fu Tuan, Phys. Rev. Lett. 59, 621 (1987).

29. M. Chaichian, N. A. Tornqvist HU-TFT-88-11 (1988).

30. Wei-Shou Hou and A. Soni, Phys. Rev. Lett. 50, 569 (1983).

31. P. G. O. Freund and Y. Nambu, Phys. Rev. Lett. 34, 1645 (1975).

32. S. J. Brodsky and C. R. Ji, Phys. Rev. Lett. 55, 2257 (1985).

33. For general discussions of $\gamma\gamma$ annihilation in $e^+e^- \to e^+e^-X$ reactions, see S. J. Brodsky, T. Kinoshita, and H. Terazawa, Phys. Rev. Lett. 25, 972 (1970), Phys. Rev. D4, 1532 (1971), V. E. Balakin, V. M. Budnev, and I. F. Ginzburg, JETP Lett. 11, 388 (1970), N. Arteaga-Romero, A. Jaccarini, and P. Kessler, Phys. Rev. D3, 1569 (1971), R. W. Brown and I. J. Muzinich, Phys. Rev. D4, 1496 (1971), and C. E. Carlson and W.-K. Tung, Phys. Rev. D4, 2873 (1971). Reviews and further references are given in H. Kolanoski and P. M. Zerwas, DESY 87-175 (1987), H. Kolanoski, in "Two-Photon Physics in e^+e^- Storage Rings," Springer Verlag (1984), and Ch. Berger and W. Wagner, Phys. Rep. 136 (1987); J. H. Field, University of Paris Preprint LPNHE 84-04 (1984).

34. G. Köpp, T. F. Walsh, and P. Zerwas, Nucl. Phys. B70, 461 (1974). F. M. Renard, Proc. of the Vth International Workshop on $\gamma\gamma$ Interactions, and Nuovo Cim. 80, 1 (1984). Backgrounds to the $C = +$, $J = 1$ signal can occur from tagged $e^+e^- \to e^+e^-X$ events which produce $C = -$ resonances.

35. S.J. Brodsky and G.P. Lepage, Phys. Rev. D24, 1808 (1981).

36. H. Aihara *et al.* , Phys. Rev. Lett. 57, 51, 404 (1986). Mark II data for combined charged meson pair production are also in good agreement with the PQCD predictions. See J. Boyer *et al.* , Phys. Rev. Lett. 56, 207 (1986).

37. H. Suura, T. F. Walsh, and B. L. Young, Lett. Nuovo Cimento 4, 505 (1972). See also M. K. Chase, Nucl. Phys. B167, 125 (1980).

38. G. W. Atkinson, J. Sucher, and K. Tsokos, Phys. Lett. 137B, 407 (1984); G. R. Farrar, E. Maina, and F. Neri, Nucl. Phys. B259, 702 (1985) Err.-ibid. B263, 746 (1986).; E. Maina, Rutgers Ph.D. Thesis (1985); J. F. Gunion, D. Millers, and K. Sparks, Phys. Rev. D33, 689 (1986); P. H. Damgaard, Nucl. Phys. B211, 435,(1983); D. Millers and J. F. Gunion, Phys. Rev. D34, 2657 (1986).

39. J. Boyer *et al.* , Ref. 36; TPC/Two Gamma Collaboration (H. Aihara *et al.*), Phys. Rev. Lett. 57, 404 (1986).

40. M. Benyayoun and V. L. Chernyak, College de France preprint LPC 89 10 (1989).

41. B. Nizic Phys. Rev. D35, 80 (1987)

42. G. R. Farrar RU-88-47, Invited talk given at Workshop on Particle and Nuclear Physics on the Light Cone, Los Alamos, New Mexico, 1988; G. R. Farrar, H. Zhang, A. A. Globlin and I. R. Zhitnitskii, Nucl. Phys. B311, 585 (1989); G. R. Farrar, E. Maina, and F. Neri, Phys. Rev. Lett. 53, 28 and, 742 (1984).

43. S. J. Brodsky, G. Kopp, and P. Zerwas Phys. Rev. Lett. 58, 443 (1987).

44. A simple method for estimating hadron pair production cross sections near threshold in $\gamma\gamma$ collisions is given in Ref. 43.

45. C. Avilez, E. Ley Koo, M. Moreno, Phys. Rev. D19 2214,(1979). C. Avilez, R. Montemayor, M. Moreno, Nuovo Cim. Lett.21 301,(1978). M. B. Voloshin, unpublished.

46. J. Schwinger, in "Particles, Sources and Fields," Addison-Wesley, New York, 1973. Vol II.

47. For a review and references, see R. M. Barnett, M. Dine, L. McLerran, Phys. Rev. D22 594, (1980).