Hard Deuteron Breakup into Delta-Isobar Channels

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Short Range Correlations in Nuclei and Hard QCD Phenomena

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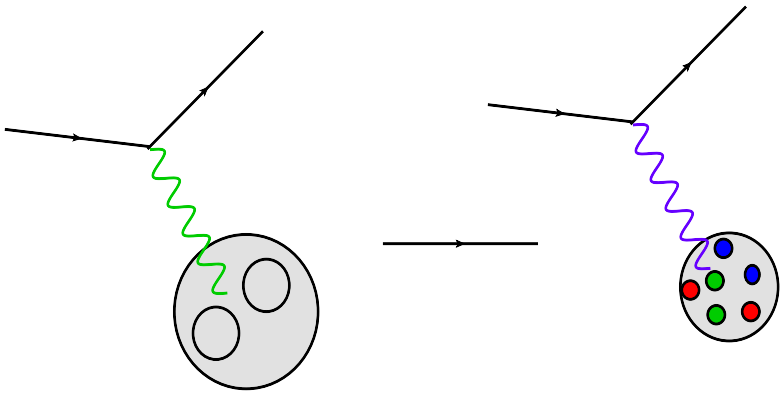
November 15, 2011

Δ-Isobar Production in Deuteron Breakup
Motivation I

The study of the deuteron photodisintegration into $\Delta\Delta$-isobars channels was proposed as a venue for investigating the evolution of a nucleon-nucleon system into a six quark system.



The onset of a six quark picture of the deuteron could then be marked by a large increase of the $\gamma d \to \Delta\Delta$ cross section. This prediction assumes that such cross section is small for a nucleon dominated deuteron wave function because of its suppressed $\Delta\Delta$ components.

In contrast, for a six quark deuteron, NN and $\Delta\Delta$ components contribute with comparable strength to the deuteron wave function (roughly 10% and 8% respectively) while more than 80% is contributed by CC (hiden color) components for which unlike N or Δ, C has a color charge (Brodsky,Ji,Lepage,1983).

$$\psi_d = \sqrt{\frac{1}{9}}\psi_{NN} + \sqrt{\frac{4}{45}}\psi_{\Delta\Delta} + \sqrt{\frac{4}{5}}\psi_{cc}$$

Δ-Isobar Production in Deuteron Breakup
Motivation II

High energy $\gamma d \to \Delta\Delta$ *with* $\Delta\Delta$ *emerging at large transverse momentum is thought to probe the onset of hidden color components in the deuteron. Assuming that the high* p_T $\Delta\Delta$ *system was created in the initial state of the interaction, in the asymptotic limit we have that*

$$\frac{d\sigma^{\gamma d \to \Delta\Delta}}{dt} \sim \frac{d\sigma^{\gamma d \to pn}}{dt}.$$

Under this same assumption we also have that,

$$\frac{d\sigma^{\gamma d \to \Delta^{++}\Delta^{-}}}{dt} = \frac{d\sigma^{\gamma d \to \Delta^{+}\Delta^{0}}}{dt},$$

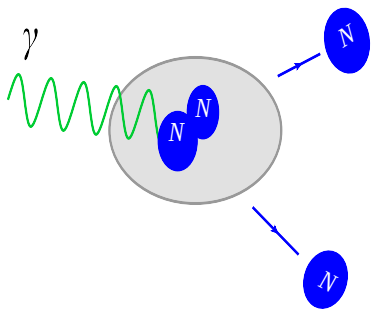
since both $\Delta\Delta$ *channels in general contribute with the same strength to the spin-isospin wave function of the deuteron. In the QCD hard rescattering model (HRM), the high* p_T $\Delta\Delta$ *system is created in a final state interaction mainly through a* pn *rescattering reaction. As it'll be shown in what follows under this scenario, the hard rescattering model predicts a dominance of* $\frac{d\sigma^{\gamma d \to \Delta^{++}\Delta^{-}}}{dt}$ *over* $\frac{d\sigma^{\gamma d \to \Delta^{+}\Delta^{0}}}{dt}$ *which contrasts the picture described above.*

Δ-Isobar Production in Deuteron Breakup
Motivation III

For comparison, we also estimate within the HRM the strength of $\Delta\Delta$ *channels in deuteron breakup relative to that of* $\gamma d \to pn$

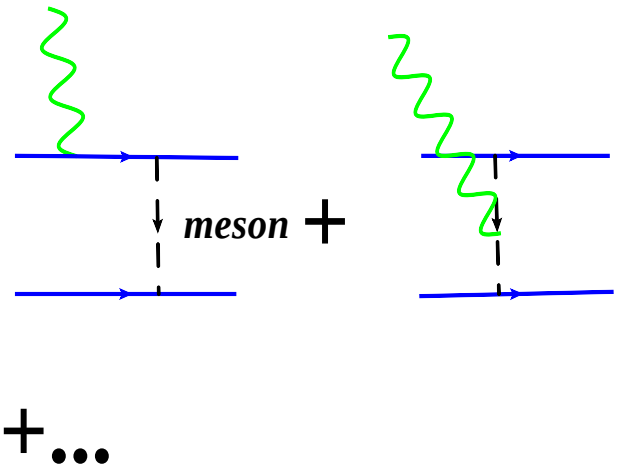
Hard Photodisintegration of a *NN* System in Nuclei I

Understanding of the strong force can be improved by studies of $\gamma + NN \to N + N$ *reactions*.

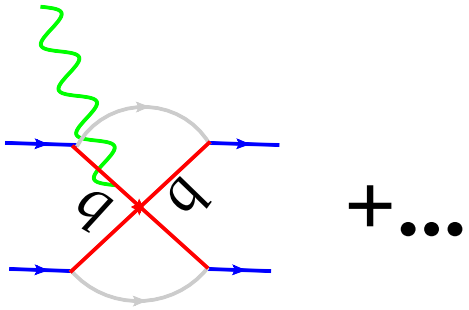


Hadronic and QCD descriptions of the NN force can be studied through such reactions.

Hadron picture



QCD picture



Hard Processes

To probe short distance structure, introduce the study of hard processes.

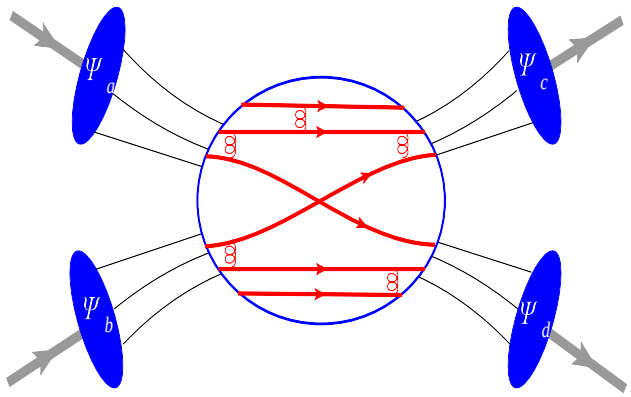
Kinematic Regime

- *Large kinematic variables,* $s >> |t|, |u|$ *and* $|t|, |u| > m_N^2$,
- *Hadron wave functions dominated by minimal Fock components,* $|N\rangle \sim |qqq\rangle$
- *Facilitate factorization schemes.*

$a + b \to c + d$

- *Mandelstam variables*
  $s = (p_a + p_b)^2 = (p_c + p_d)^2$
  $t = (p_c - p_a)^2 = (p_d - p_b)^2$
  $u = (p_d - p_a)^2 = (p_c - p_b)^2$

Explore the role of quarks and gluons through Hard NN scattering



Kinematic Advantages of Hard $\gamma + NN \to N + N$

$N + N$ emerge at large transverse momentum from a short range NN interaction.

$$-t \sim_{s\to\infty} \frac{s}{2}(1 - cos(\theta_{c.m.}))$$

Large center of mass energy of the emerging nucleons $\sqrt{s}/2$ more efficiently reached in $\gamma + NN \to N + N$,

$$s_{\gamma NN} \approx 4m_N^2 + 2E_\gamma \cdot 2m_N,$$

compared to

$$s_{NN} = 2m_N^2 + 2E \cdot m_N$$

in $N_{beam} + N_{target} \longrightarrow N + N$ processes (with E being the energy of the nucleon beam).

Then, at moderate $E_\gamma \sim$GeV and at large $\theta_{c.m.}$, $\gamma + NN \to N + N$ reaches the hard kinematic regime.

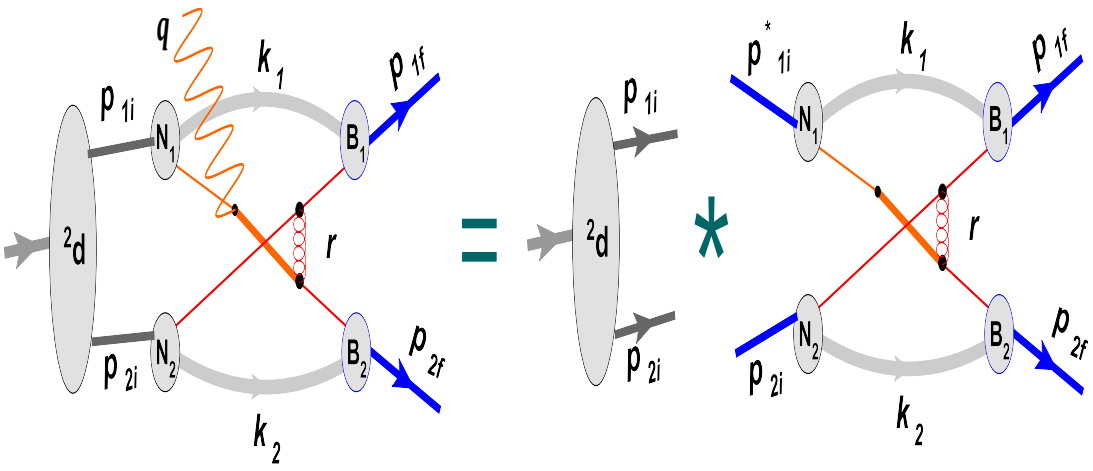
Producing High p_T $N + N$ in $\gamma + NN \to N + N$

There are two basic approaches,

- *Photon probes a preexisting compact NN system (large p_T component in initial state) in nucleus . This is assumed for instance in approaches such as the reduced nuclear amplitude (RNA)(Brodsky,Hiller,1983).*
  Or,
- *the energetic photon is absorbed by a low p_T NN system. The energy transfer triggers a final state interaction from which the nucleons emerge at large relative transverse.*

The latter is the scenario for the hard rescattering model (HRM) (Frankfurt et al. 2000) for which we develop applications in what follows.

Hard rescattering model, HRM



L. Frankfurt et al.,Phys. Rev. Lett. 84, 3045 (2000)

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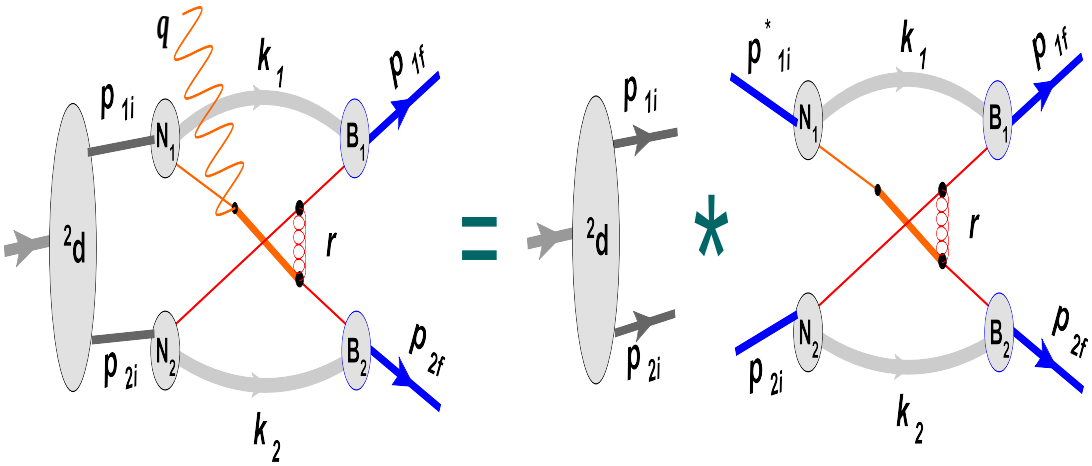
Hard rescattering model, HRM

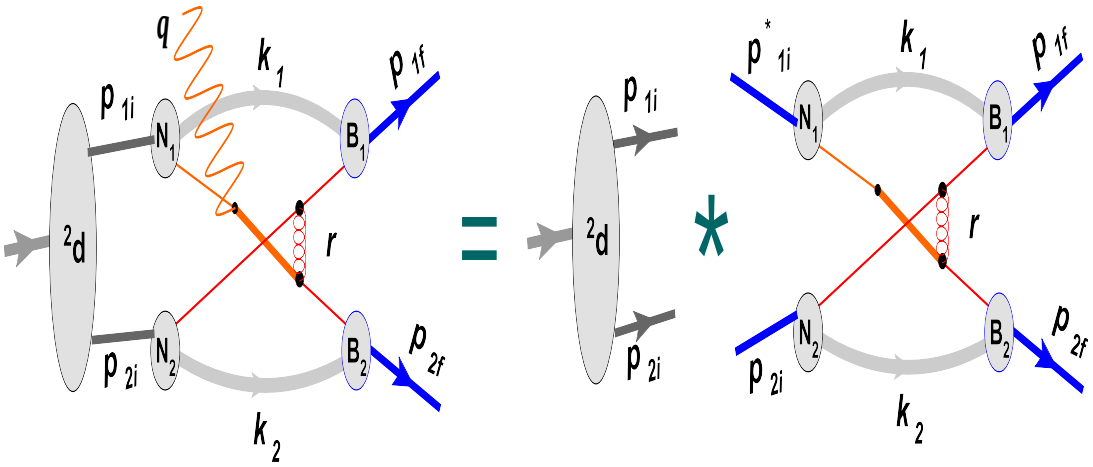


L. Frankfurt et al.,Phys. Rev. Lett. 84, 3045 (2000)

- $\gamma + d \to p + n$ *HRM amplitude results in a convolution of a hard process amplitude and a nuclear wave function,*

$$\langle \lambda_{1f}, \lambda_{2f} \mid \mathcal{M} \mid \lambda_\gamma, \lambda_d \rangle = -\frac{i[\lambda_\gamma] e Q_f \sqrt{2}(2\pi)^3}{\sqrt{2S'_{NN}}} \times \sum_{\lambda_{2i}} \int \langle \lambda_{2f}; \lambda_{1f} \mid T^{QIM}_{NN}(s_{NN}, t_N) \mid \lambda_\gamma; \lambda_{2i} \rangle$$
$$\times \quad \Psi^{\lambda_d \lambda_\gamma; \lambda_{2i}}_{d,NR}(\vec{p}_1, \vec{p}_2) m_N \frac{d^2 p_\perp}{(2\pi)^2},$$

Hard rescattering model, HRM



L. Frankfurt et al.,Phys. Rev. Lett. 84, 3045 (2000)

- *A parameter free cross section calculable through the input of experimental data,*

$$\frac{d\sigma^{\gamma d\to pn}}{dt}(s,\theta_{c.m.}) = \frac{8\alpha}{9}\pi^4\frac{1}{s'}C(\frac{\tilde{t}}{s}) \times \frac{\mathbf{d\sigma^{pn\to pn}}}{\mathbf{dt}}(\mathbf{s},\theta^{\mathbf{N}}_{\mathbf{c.m.}})$$
$$\times \left|\int \Psi_d^{NR}(p_z=0,p_t)\sqrt{m_n}\frac{d^2p_t}{(2\pi)^2}\right|^2$$

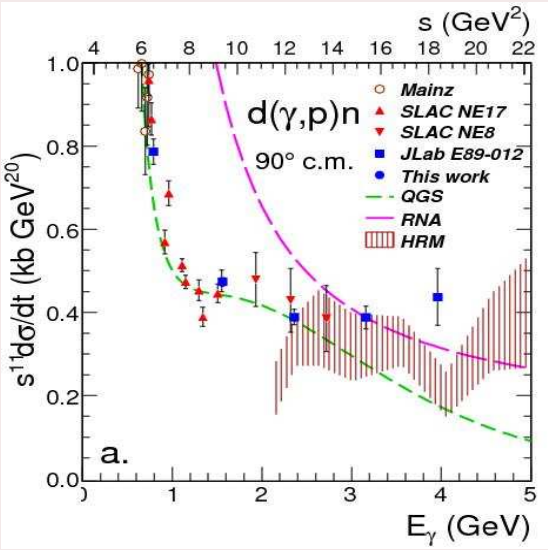
- *Energy and Angular distributions in pn elastic scattering data should be reflected by* $\gamma + d \longrightarrow p + n$ *data.*

Energy Distribution

$$\frac{d\sigma^{\gamma d\to pn}}{dt}(s,\theta_{c.m.}) = \frac{8\alpha}{9}\pi^4\frac{1}{s'}C(\frac{\tilde{t}}{s}) \times \frac{d\sigma^{pn\to pn}}{dt}(s,\theta^{N}_{c.m.})$$

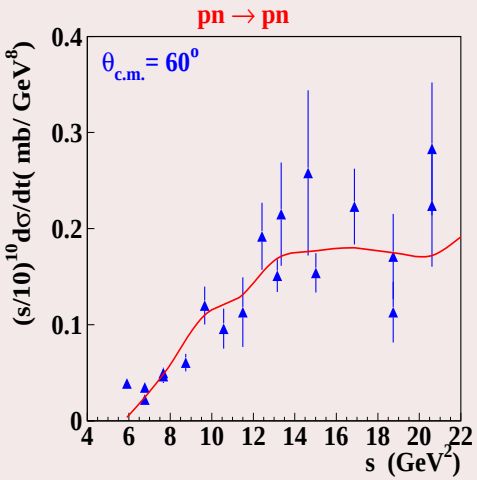
$$\times \left|\int \Psi_d^{NR}(p_z=0,p_t)\sqrt{m_n}\frac{d^2p_t}{(2\pi)^2}\right|^2$$

$\gamma + d \to p + n$



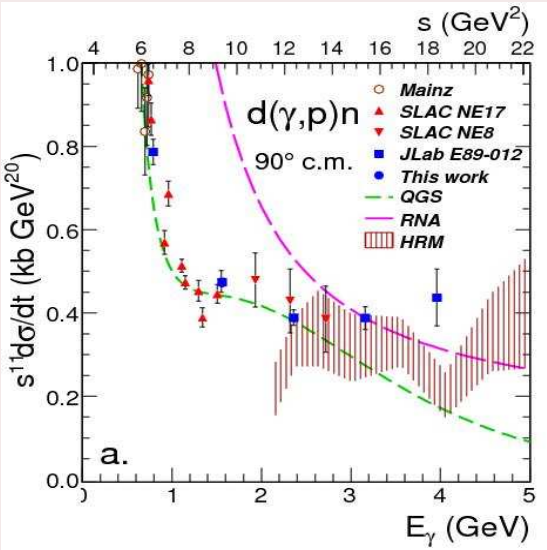
Bochna et al. 1998

$p + n \to p + n$



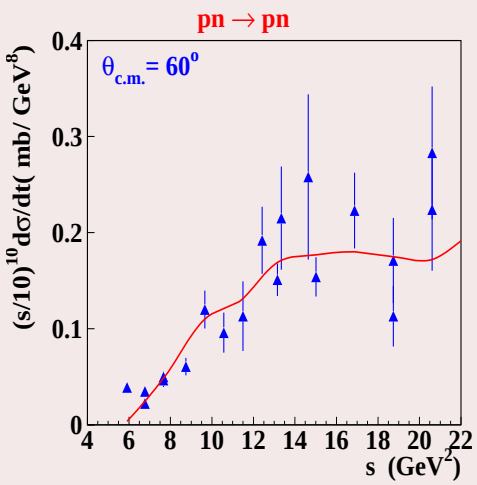
Energy Distribution

$\gamma + d \rightarrow p + n$



Bochna et al. 1998

$p + n \rightarrow p + n$

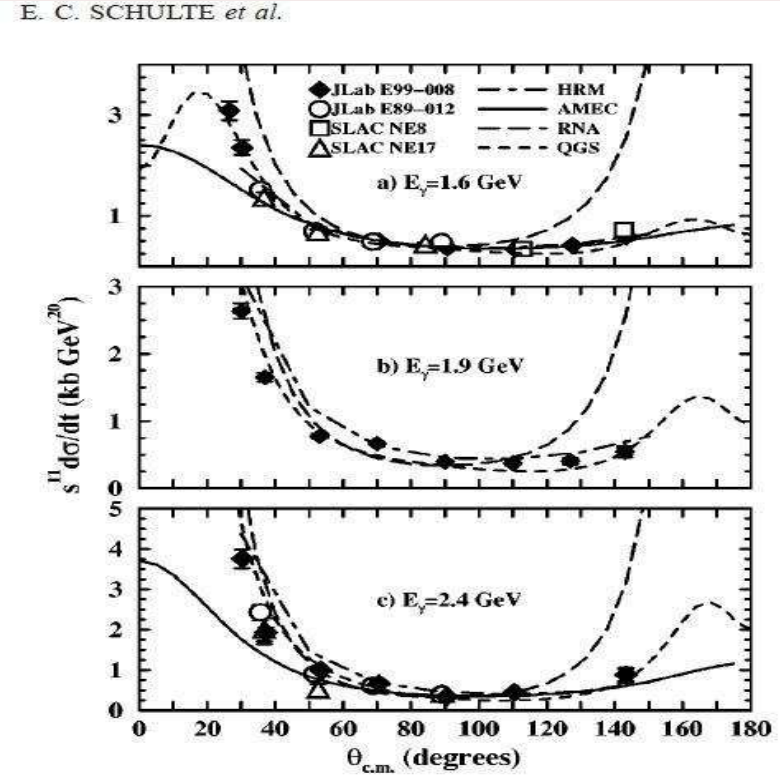


$\gamma + d \rightarrow p + n \sim s^{-11}$ ***for*** $E_\gamma > 2GeV$, $s > 10\boldsymbol{GeV}^2$

However, pn elastic scattering data is not good enough to clearly confirm HRM prediction

Angular Distributions

$\gamma + d \rightarrow p + n$



Schulte et al., Phys.Rev.C66:042201,2002.

$p + n \rightarrow p + n$

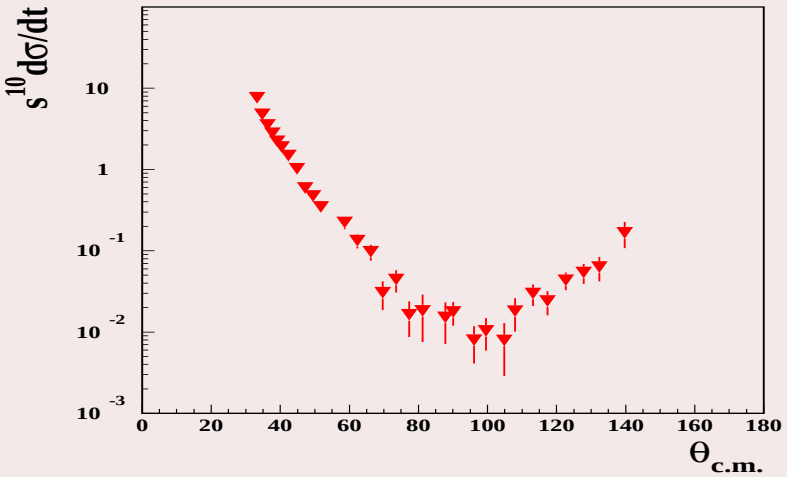
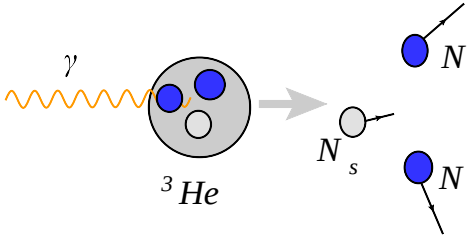
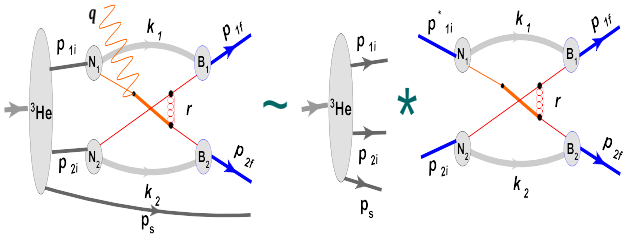


Figure: pn elastic scattering for p_{Lab}=8GeV/c

Summary: Hard rescattering model of $\gamma^3He \to N + N$



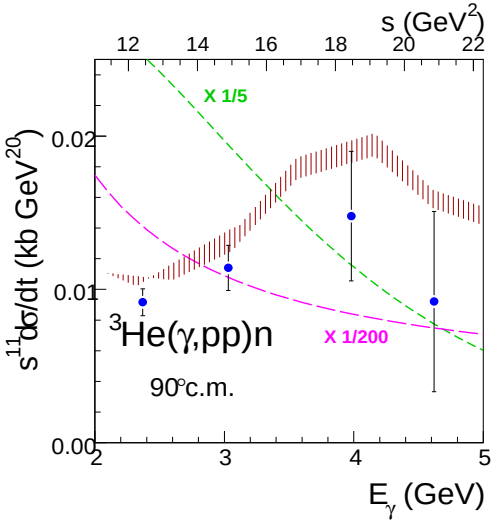


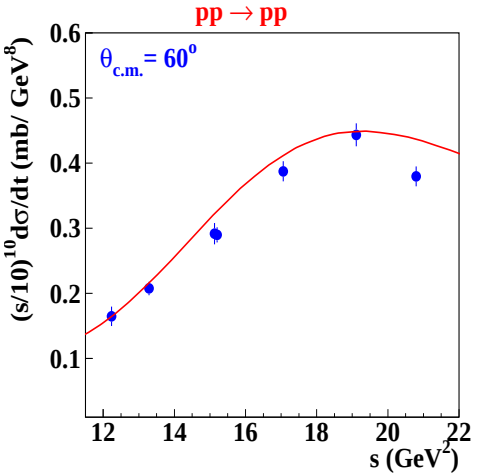
$$\frac{d\sigma^{\gamma^3He\to(pp)n}}{dt\,\frac{d^3p_s}{E_s}} = \alpha Q^2_{F,pp} 16\pi^4 S^{pp}_{34}(\alpha,\vec{p}_s)\,\frac{2\beta^2}{1+2C^2}\,\frac{s_{NN}(s_{NN}-4m_N^2)}{(s_{NN}-p^2_{NN})^2(s-M^2_{^3\mathrm{He}})} \times \frac{d\sigma^{pp\to pp}(s_{NN},t_N)}{dt},$$

M.Sargsian, CG,PRC 80,(2009)

By introducing quark d.o.f,

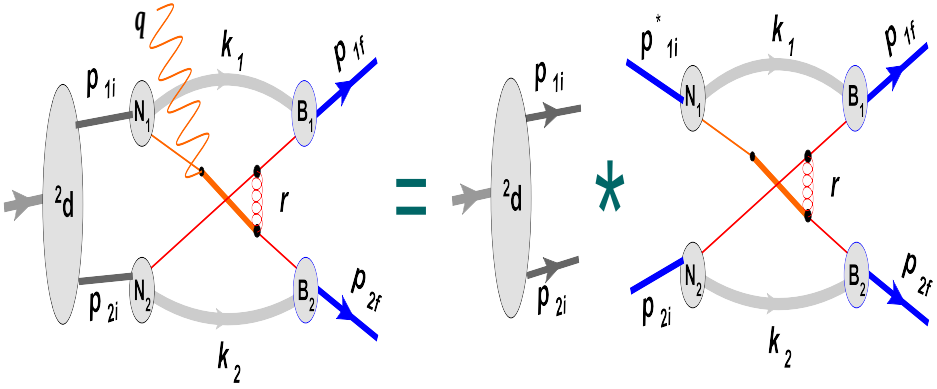
- *a scattering amplitude is obtained from a convolution of an electromagnetic amplitude, a nuclear wave function and a hard NN scattering amplitude.*
- *Cross section is computed from corresponding data of NN cross section.*
- *Good agreement with experiment is achieved. No need of fitting parameter. Energy dependency also in agreement with constituent counting rules as predicted.*





JLab Hall A.(Pomerantz et al.

Δ-Isobar Production in Deuteron Breakup
$\gamma d \to \Delta\Delta$ in the HRM



As it was the case for $\gamma d \to pn$, through the HRM for $\gamma d \to \Delta\Delta$ we find that,

$$\frac{d\sigma^{\gamma d \to \Delta\Delta}(s, \theta_{c.m.})}{dt} = \frac{\alpha Q^2_{F,\Delta\Delta} 8\pi^4}{s'} \frac{d\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt} \bar{S}_{0,NR},$$

But $\frac{\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt}$ *has not been well measured in the required kinematic regime.*

C. G. Granados and M. M. Sargsian, Phys.Rev. C83 (2011)

Δ-Isobar Production in Deuteron Breakup $\gamma d \to \Delta\Delta$ in the HRM

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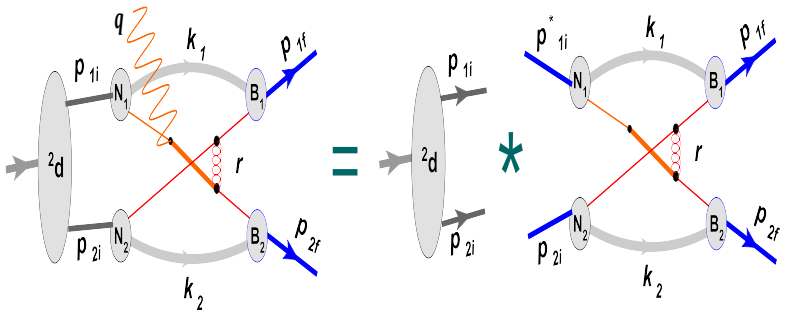
$$\frac{d\sigma^{\gamma d \to \Delta\Delta}(s, \theta_{c.m.})}{dt} = \frac{\alpha Q^2_{F,\Delta\Delta} 8\pi^4}{s'} \frac{d\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt} \bar{S}_{0,NR},$$

But $\frac{\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt}$ has not been well measured in the required kinematic regime. Then model dependent $pn \to \Delta\Delta$ amplitudes ϕ are input in

$$\begin{aligned}
|\bar{\mathcal{M}}|^2{}_{\gamma d \to \Delta\Delta} &= \frac{1}{2}\frac{1}{3}\frac{e^2}{2s'} \\
&\times \Big[S_{12} \Big\{ |(\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_1|^2 + |(\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_6|^2 + |(\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_7|^2 \Big\} \\
&+ S_{34} \Big\{ |\hat{Q}^{N_1}\phi_3 + \hat{Q}^{N_2}\phi_4|^2 + |\hat{Q}^{N_1}\phi_4 + \hat{Q}^{N_2}\phi_3|^2 \\
&+ |\hat{Q}^{N_1}\phi_8 + \hat{Q}^{N_2}\phi_9|^2 + |\hat{Q}^{N_1}\phi_9 + \hat{Q}^{N_2}\phi_8|^2 \Big\} \Big],
\end{aligned}$$

C. G. Granados and M. M. Sargsian, Phys.Rev. C83 (2011)

Summary: HRM angular distributions in $\gamma d \to \Delta\Delta$



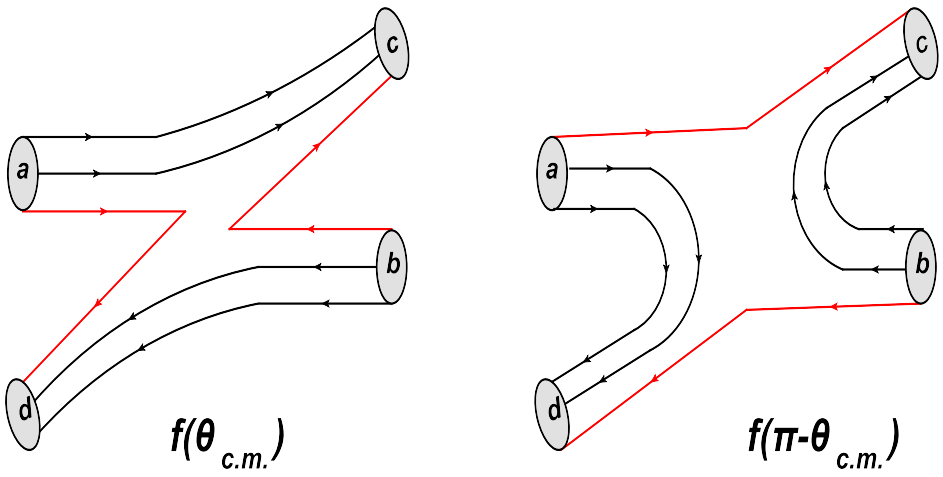
But $\frac{\sigma^{pn\to\Delta\Delta}(s,\theta^N_{c.m.})}{dt}$ *has not been well measured in the required kinematic regime. Then model dependent* $pn \to \Delta\Delta$ *amplitudes* ϕ *are input in*

$$
\begin{aligned}
|\bar{\mathcal{M}}|^2{}_{\gamma d\to\Delta\Delta} &= \frac{1}{2}\frac{1}{3}\frac{e^2}{2s'} \\
&\times \Big[S_{12}\Big\{ |(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_1|^2 + |(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_6|^2 \\
&\quad + |(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_7|^2 \Big\} \\
&\quad + S_{34}\Big\{ |\hat{Q}^{N_1}\phi_3 + \hat{Q}^{N_2}\phi_4|^2 + |\hat{Q}^{N_1}\phi_4 + \hat{Q}^{N_2}\phi_3|^2 \\
&\quad + |\hat{Q}^{N_1}\phi_8 + \hat{Q}^{N_2}\phi_9|^2 + |\hat{Q}^{N_1}\phi_9 + \hat{Q}^{N_2}\phi_8|^2 \Big\} \Big] ,
\end{aligned}
$$

The amplitudes ϕ *are estimated within the quark interchange model.*

Quark Interchange Mechanism (QIM)

See e. g., Farrar et al., PRD20(79)
and Brodsky et al. PRD20(79)



Experimentally shown to be dominant for 90° c.m. hard exclusive scattering of baryons containing quarks of common flavor, (White et al. PRD49(94)).

Quark Interchange Mechanism (QIM)

For an exclusive baryon baryon scattering process $(B_a B_b \to B_c B_d)$ *in the hard kinematic regime,*

$$\begin{aligned}\langle \psi_c \psi_d \mid \mathcal{M} \mid \psi_a \psi_b \rangle \quad = \quad & \sum_{\alpha,\beta,\gamma} \langle \psi_c \mid \alpha_2', \beta_1', \gamma_1' \rangle \langle \psi_d \mid \alpha_1', \beta_2', \gamma_2' \rangle \\ & \times \langle \alpha_2', \beta_2', \gamma_2', \alpha_1' \beta_1' \gamma_1' \mid H \mid \alpha_1, \beta_1, \gamma_1, \alpha_2, \beta_2, \gamma_2 \rangle \\ & \times \langle \alpha_1, \beta_1, \gamma_1 \mid \psi_a \rangle \langle \alpha_2, \beta_2, \gamma_2 \mid \psi_b \rangle,\end{aligned}$$

in which the expansion coefficients

$$C^j_{\alpha,\beta,\gamma} \equiv \langle \alpha, \beta, \gamma \mid \psi_j \rangle,$$

are independent of the hard kinematic variables. In the quark interchange model (QIM), the hard factor takes the form

$$H_{QIM} \approx \delta_{\alpha_1 \alpha_1'} \delta_{\alpha_2 \alpha_2'} \delta_{\beta_1, \beta_1'} \delta_{\gamma_1, \gamma_1'} \delta_{\beta_2, \beta_2'} \delta_{\gamma_2, \gamma_2'} \frac{f(\theta_{c.m})}{s^4}.$$

Quark Interchange Mechanism (QIM)

Introduce a parameter ρ to break SU(6) in scalar [qq], and vector (qq) diquark components,

$[qq] = |h = 0, I = 0\rangle$

$(qq) = |h = 1, I = 1\rangle$

$$Baryon \sim q[qq] + \rho \times q(qq)$$

There is not scalar diquark component for Δ-isobar, while for the nucleon wave function

$$\begin{aligned}
\psi^{i_N^3, h_N} \quad = \quad & \frac{1}{\sqrt{2}} \Big\{ \Phi_{0,0}(k_1, k_2, k_3)(\chi^{(23)}_{0,0} \chi^{(1)}_{\frac{1}{2}, h_N}) \cdot (\tau^{(23)}_{0,0} \tau^{(1)}_{\frac{1}{2}, i_N^3}) + \Phi_{1,1}(k_1, k_2, k_3) \times \\
& \sum_{i_{23}^3 = -1}^{1} \sum_{h_{23}^3 = -1}^{1} \langle 1, h_{23}; \frac{1}{2}, h_N - h_{23} \mid \frac{1}{2}, h_N \rangle \langle 1, i_{23}^3; \frac{1}{2}, i_N^3 - i_{23}^3 \mid \frac{1}{2}, i_N^3 \rangle \\
& \times (\chi^{(23)}_{1, h_{23}} \chi^{(1)}_{\frac{1}{2}, h_N - h_{23}}) \cdot (\tau^{(23)}_{1, i_{23}^3} \tau^{(1)}_{\frac{1}{2}, i_N^3 - i_{23}^3}) \Big\}
\end{aligned}$$

Then,

$$\rho = \frac{\langle \Phi_{1,1} \rangle}{\langle \Phi_{0,0} \rangle}$$

C. Granados, M. Sargsian, Phys. Rev. Lett. 103, 212001 (2009).

Quark Interchange Mechanism (QIM)

Working explicitly a nucleon wave function in quark (qqq) components yields,

$$\begin{aligned}\mathbf{N}(\rho)|\mathbf{p}(+)\rangle \quad = \quad & \frac{3+\rho}{2}|u(+)u(+)d(-)\rangle - \frac{3-\rho}{2}|u(+)d(+)u(-)\rangle \\ & -\rho|d(+)u(+)u(-)\rangle + \frac{3+\rho}{2}|u(+)d(-)u(+)\rangle \\ & -\frac{3-\rho}{2}|u(+)u(-)d(+)\rangle - \rho|d(+)u(-)u(+)\rangle \\ & +2\rho|d(-)u(+)u(+)\rangle - \rho|u(-)u(+)d(+)\rangle \\ & -\rho|u(-)d(+)u(+),\rangle\end{aligned}$$

which results in a ρ parameterization of the scattering amplitude and consequently of the cross section,

$$\sigma(s,\theta_{c.m},\rho) = \sigma(s)\left[C_t(\rho)F^2(\theta_{c.m.}) + C_{tu}(\rho)F(\theta_{c.m})F(\pi-\theta_{c.m.}) + C_u(\rho)F^2(\pi -\right.$$

choosing

$$F(\theta_{c.m.}) \propto \frac{1}{(sin(\theta_{c.m})(1-cos(\theta_{c.m.})))^2}$$

C. Granados, M. Sargsian, Phys. Rev. Lett. 103, 212001 (2009).

Quark Interchange Mechanism (QIM)

Independent amplitudes in NN elastic scattering

$$\begin{aligned}
\left\langle +,+|T_{NN}^{QIM}|+,+\right\rangle &= \phi_1 \\
\left\langle +,+|T_{NN}^{QIM}|+,-\right\rangle &= \phi_5 \\
\left\langle +,+|T_{NN}^{QIM}|-,-\right\rangle &= \phi_2 \\
\left\langle +,-|T_{NN}^{QIM}|+,-\right\rangle &= \phi_3 \\
\left\langle +,-|T_{NN}^{QIM}|-,+\right\rangle &= -\phi_4 .
\end{aligned}$$

np elastic scattering amplitudes are derived in the QIM yielding,

$$\begin{aligned}
\phi_1(\theta_{c.m.}) &= (2-y)f(\theta_{c.m.}) + (1+2y)f(\pi-\theta_{c.m.}) \\
\phi_2(\theta_{c.m.}) &= 0 \\
\phi_3(\theta_{c.m.}) &= (2+y)f(\theta_{c.m.}) + (1+4y)f(\pi-\theta_{c.m.}) \\
\phi_4(\theta_{c.m.}) &= 2yf(\theta_{c.m.}) + 2yf(\pi-\theta_{c.m.}) \\
\phi_5(\theta_{c.m.}) &= 0
\end{aligned} \tag{1}$$

where

$$\begin{aligned}
y &= x(1+x) \\
&= \frac{2}{3}\frac{\rho}{1+\rho^2}\left(1+\frac{2}{3}\frac{\rho}{1+\rho^2}\right)
\end{aligned}$$

C. Granados, M. Sargsian, Phys. Rev. Lett. 103, 212001 (2009).

QIM scattering amplitudes $pn \to \Delta\Delta$

$$\phi(\theta^N_{c.m.}) = \phi_t(\theta^N_{c.m.}) + \phi_u(\theta^N_{c.m.})$$

Helicity conserving $pn \to \Delta\Delta$ *amplitudes*

$$\begin{aligned}
\left\langle +\frac{1}{2}, +\frac{1}{2} \,|T|\, +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_1 \\
\left\langle +\frac{1}{2}, -\frac{1}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_3 \\
\left\langle -\frac{1}{2}, +\frac{1}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_4 \\
\left\langle +\frac{3}{2}, -\frac{1}{2} \,|T|\, +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_6 \\
\left\langle -\frac{1}{2}, +\frac{3}{2} \,|T|\, +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_7 \\
\left\langle +\frac{3}{2}, -\frac{3}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_8 \\
\left\langle -\frac{3}{2}, +\frac{3}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_9,
\end{aligned}$$

QIM scattering amplitudes $pn \to \Delta\Delta$

$$\phi(\theta^N_{c.m.}) = \phi_t(\theta^N_{c.m.}) + \phi_u(\theta^N_{c.m.})$$

Helicity conserving $pn \to \Delta\Delta$ *amplitudes* *Obtained in the quark exchange model*

$pn \to \Delta^+\Delta^0$

$$\begin{aligned}
\phi_1 &= \frac{2}{9} N_{\Delta\Delta}(2f(\theta^N_{c.m.}) - f(\pi - \theta^N_{c.m.})) \\
\phi_3 &= \frac{1}{9} N_{\Delta\Delta}(4f(\theta^N_{c.m.}) + f(\pi - \theta^N_{c.m.})) \\
\phi_4 &= \frac{2}{9} N_{\Delta\Delta}(f(\theta^N_{c.m.})) + f(\pi - \theta^N_{c.m.}) \\
\phi_6 &= \frac{N_{\Delta\Delta}}{3\sqrt{3}}(2f(\theta^N_{c.m.}) - f(\pi - \theta^N_{c.m.})) \\
\phi_7 &= \frac{N_{\Delta\Delta}}{3\sqrt{3}}(2f(\theta^N_{c.m.}) - f(\pi - \theta^N_{c.m.})) \\
\phi_8 &= \frac{2}{9} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_9 &= \frac{1}{3} N_{\Delta\Delta} f(\pi - \theta^N_{c.m.}),
\end{aligned}$$

$pn \to \Delta^{++}\Delta^-$

$$\begin{aligned}
\phi_1 &= -\frac{2}{3} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_3 &= -\frac{2}{3} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_4 &= -\frac{1}{3} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_6 &= \frac{-N_{\Delta\Delta}}{\sqrt{3}} f(\theta^N_{c.m.}) \\
\phi_7 &= \frac{-N_{\Delta\Delta}}{\sqrt{3}} f(\theta^N_{c.m.}) \\
\phi_8 &= -N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_9 &= 0,
\end{aligned}$$

$\gamma d \to \Delta\Delta$ in the HRM

$$\langle \lambda_{1f}, \lambda_{2f} \mid \mathcal{M} \mid \lambda_\gamma, \lambda_d \rangle = -\frac{i[\lambda_\gamma] e Q_f \sqrt{2}(2\pi)^3}{\sqrt{2S'_{NN}}} \times \sum_{\lambda_{2i}} \int \langle \lambda_{2f}; \lambda_{1f} \mid T^{QIM}_{BB}(s, t_N) \mid \lambda_\gamma; \lambda_{2i} \rangle$$

$$\times \quad \Psi^{\lambda_d \lambda_\gamma; \lambda_{2i}}_{d,NR}(\vec{p}_1, \vec{p}_2) m_N \frac{d^2 p_\perp}{(2\pi)^2},$$

$\gamma d \to \Delta\Delta$ in the HRM

$$
\begin{aligned}
\left\langle +\frac{1}{2}, +\frac{1}{2} \,|\mathcal{M}|\, +, \lambda_d \right\rangle &= B\left((\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_1 \int \Psi_{\mathrm{d}}^{\lambda_d}(p_1, +; p_2, +) \frac{d^2 p_\perp}{(2\pi)^2} \right) \\
\left\langle +\frac{1}{2}, -\frac{1}{2} \,|\mathcal{M}|\, +, \lambda_d \right\rangle &= B\left((\hat{Q}^{N_1}\phi_3 + \hat{Q}^{N_2}\phi_4) \int \Psi_{\mathrm{d}}^{\lambda_d}(p_1, +; p_2, -) \frac{d^2 p_\perp}{(2\pi)^2} \right) \\
\left\langle -\frac{1}{2}, +\frac{1}{2} \,|\mathcal{M}|\, +, \lambda_d \right\rangle &= B\left((\hat{Q}^{N_1}\phi_4 + \hat{Q}^{N_2}\phi_3) \int \Psi_{\mathrm{d}}^{\lambda_d}(p_1, +; p_2, -) \frac{d^2 p_\perp}{(2\pi)^2} \right) \\
\left\langle +\frac{3}{2}, -\frac{1}{2} \,|\mathcal{M}|\, +, \lambda_d \right\rangle &= B\left((\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_6 \int \Psi_{\mathrm{d}}^{\lambda_d}(p_1, +; p_2, +) \frac{d^2 p_\perp}{(2\pi)^2} \right) \\
\left\langle -\frac{1}{2}, +\frac{1}{2} \,|\mathcal{M}|\, +, \lambda_d \right\rangle &= B\left((\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_7 \int \Psi_{\mathrm{d}}^{\lambda_d}(p_1, +; p_2, +) \frac{d^2 p_\perp}{(2\pi)^2} \right) \\
\left\langle +\frac{3}{2}, -\frac{3}{2} \,|\mathcal{M}|\, +, \lambda_d \right\rangle &= B\left((\hat{Q}^{N_1}\phi_8 + \hat{Q}^{N_2}\phi_9) \int \Psi_{\mathrm{d}}^{\lambda_d}(p_1, +; p_2, -) \frac{d^2 p_\perp}{(2\pi)^2} \right) \\
\left\langle -\frac{3}{2}, +\frac{3}{2} \,|\mathcal{M}|\, +, \lambda_d \right\rangle &= B\left((\hat{Q}^{N_1}\phi_9 + \hat{Q}^{N_2}\phi_8) \int \Psi_{\mathrm{d}}^{\lambda_d}(p_1, +; p_2, -) \frac{d^2 p_\perp}{(2\pi)^2} \right),
\end{aligned}
\tag{2}
$$

$\gamma d \to \Delta\Delta$ in the HRM

$$
\begin{aligned}
|\bar{\mathcal{M}}|^2_{\gamma d \longrightarrow \Delta^+\Delta^-} = \; & \frac{1}{6}\frac{e^2}{2s'} Q^2_{F,\Delta\Delta} \Big\{ S_{12} \left[|\phi_1|^2 + |\phi_6|^2 + |\phi_7|^2 \right] \\
& + S_{34} \left[\left(\frac{\phi_3^t + \phi_4^t}{2} + 2\phi_4^u - \phi_3^u \right)^2 + \left(\frac{\phi_4^t + \phi_3^t}{2} + 2\phi_3^u - \phi_4^u \right) \right. \\
& \left. + \left(\frac{\phi_8^t + \phi_9^t}{2} + 2\phi_9^u - \phi_8^u \right)^2 + \left(\frac{\phi_9^t + \phi_8^t}{2} + 2\phi_8^u - \phi_9^u \right)^2 \right] \Big\}
\end{aligned}
$$

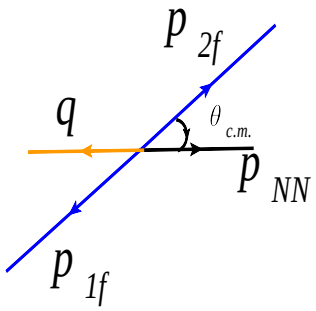
$$
\begin{aligned}
|\bar{\mathcal{M}}|^2_{\gamma d \longrightarrow \Delta^{++}\Delta^-} = & \frac{1}{6}\frac{e^2}{2s'} Q^2_{F,\Delta\Delta} \Big\{ S_{12} \left(|\phi_1|^2 + |\phi_6|^2 + |\phi_7|^2 \right) \\
& + S_{34} \left[(2\phi_3 - \phi_4)^2 + (2\phi_4 - \phi_3)^2 + 5|\phi_8|^2 \right] \Big\} .
\end{aligned}
$$

γNN and $B + B$ Kinematics I

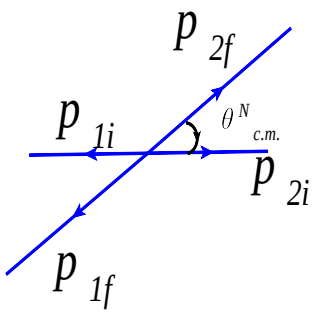
Calculating $\frac{d\sigma^{\gamma D \to BB}}{dt}(s_{NN}, \theta_{c.m.})$ *within HRM, using* $\phi^{NN \to BB}(s_{NN}, \theta^{N}_{c.m.})$ *requires using the right correspondence between* $\theta_{c.m}$ *and* $\theta^{N}_{c.m.}$.

In the center of mass reference frame for,

$\gamma + NN \to B + B$



$N + N \to B + B$



$$t = (p_{NN} - p_{2f})^2 \qquad t^N = (p_{2i} - p_{2f})^2$$

which leads to,

$$t^N = \frac{t}{2} + \frac{m^2_{B_2}}{2} - \frac{M^2_{NN}}{4}$$

γNN and $B + B$ Kinematics II

that together with

$$t_N = -\frac{s}{2} + m_n^2 + m_B^2 + \frac{1}{2}\sqrt{\left(s - 4m_N^2\right)\left(s - 4m_B^2\right)}\, cos\theta_{c.m.}^N,$$

and

$$t = m_B^2 - \frac{s - M_{NN}^2}{2\sqrt{s}}\left(\sqrt{s} - \sqrt{s - 4m_B^2}\, cos\theta_{c.m.}\right),$$

is used to obtain that,

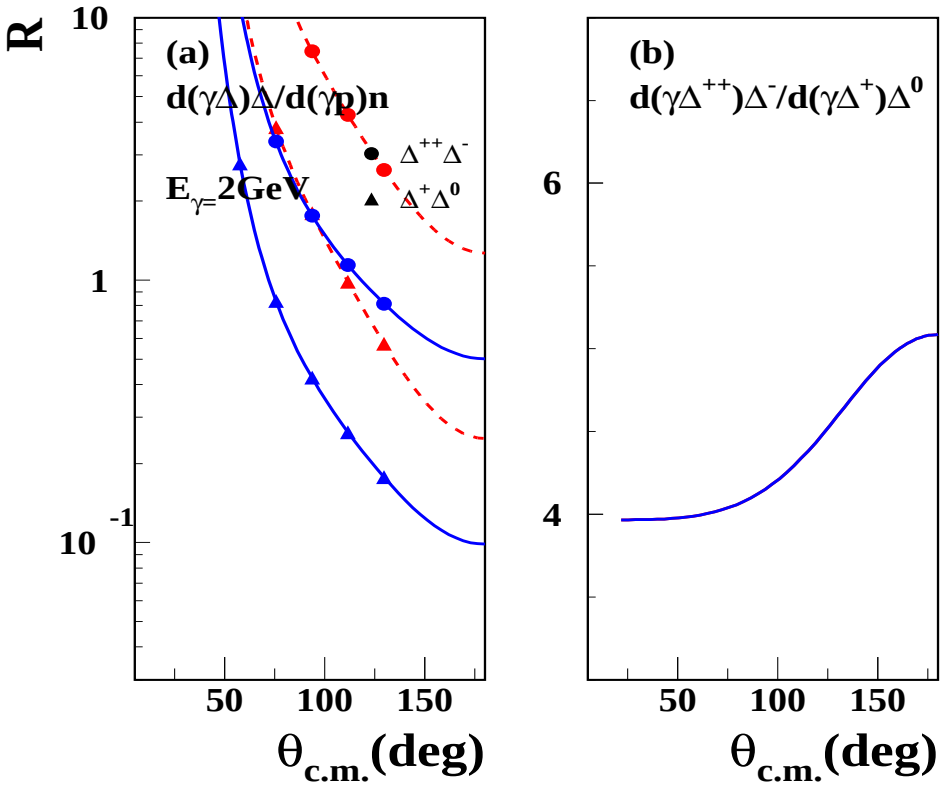
$$cos\theta_{c.m.}^N = \frac{1}{2\sqrt{\left(s - 4m_N^2\right)\left(s - 4m_\Delta^2\right)}}\left[\frac{s - M_d^2}{2\sqrt{s}}\sqrt{s - 4m_\Delta^2}\, cos\theta_{c.m.} + s - 4m_N^2\right].$$

For $s \to \infty$,

$$cos\theta_{c.m.}^N \sim \frac{cos\theta_{c.m.} + 1}{2}$$

Summary of results

HRM angular distributions in $\gamma d \to \Delta\Delta$



R in Fig. (a)

$$R = \frac{\sigma^{\gamma D \to \Delta\Delta}}{\sigma^{\gamma d \to pn}}.$$

- *is sensitive to model of baryon wavefunction*
- *shows a strong dependence on $\theta_{c.m}$ contrasting R$\sim$ 1 expected from the onset of hidden color components.*

In Fig. (b)

$$R = \frac{\sigma^{\gamma D \to \Delta^{++}\Delta^{-}}}{\sigma^{\gamma d \to \Delta^{+}\Delta^{0}}}.$$

- *R doesn't depend on the choice of wave function*
- *$\Delta^{++}\Delta^{-}$ channel is consistently larger than that of the $\Delta^{+}\Delta^{0}$ channel. If the isobars were produced from a $\Delta\Delta$ component of the initial deuteron wave function, both channels should have the same strength and R=1.*

C. G. Granados and M. M. Sargsian, Phys. Rev. C 83, 054606 (2011).

Summary of results

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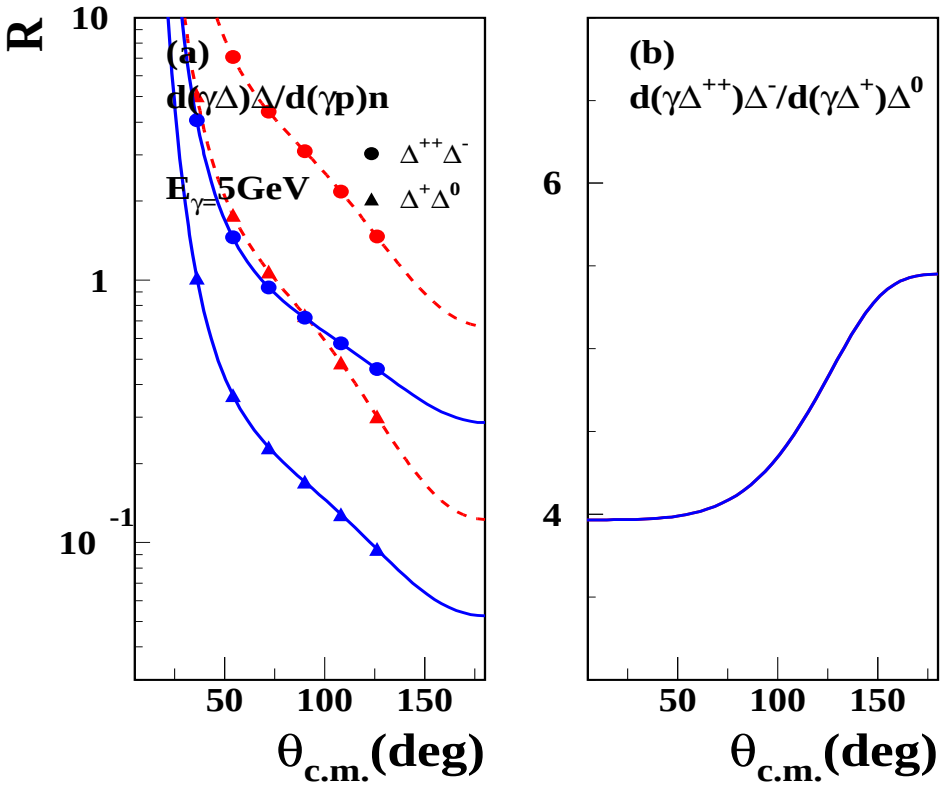

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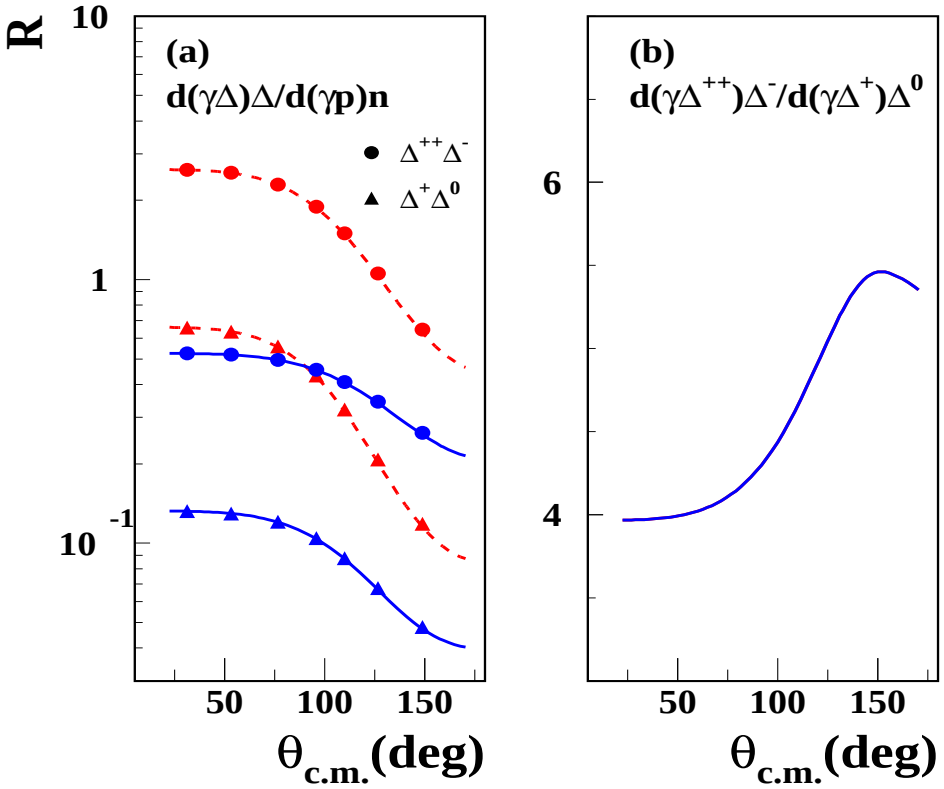

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Thank you

Δ-Isobar Production in Deuteron Breakup
Hard Processes Involving Nuclei
Hard Rescatrering Model of $\gamma + NN \to NN$

Introduction
Hard Processes
Hard Processes Involving Nuclei
Hard Rescatrering Model of $\gamma + NN \to NN$

^3He Photodisintegration within the HRM
Scattering Amplitude
pp and pn breakup
Polarization Transfer
Summary

Δ-Isobar Production in Deuteron Breakup
Angular Distributions

Further Applications

Hard Processes

To probe short distance structure, introduce the study of hard processes.

Kinematic Regime

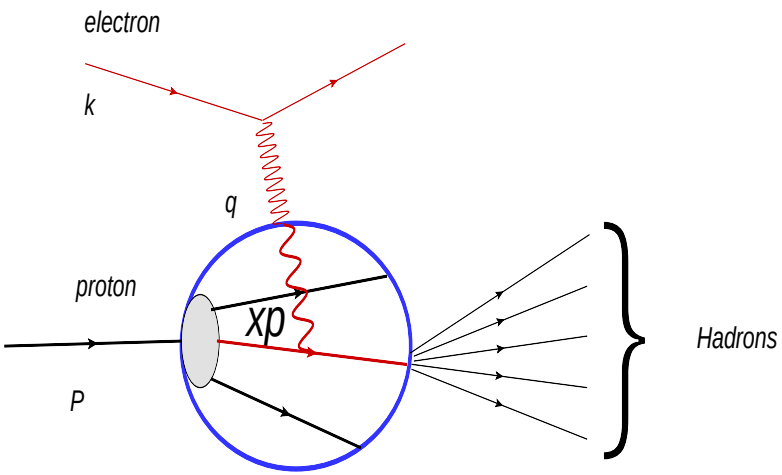
- *Large kinematic variables,* $s >> |t|, |u|$ *and* $|t|, |u| > m_N^2$,
- *Hadron wave functions dominated by minimal Fock components,* $|N\rangle \sim |qqq\rangle$
- *Facilitate factorization schemes.*

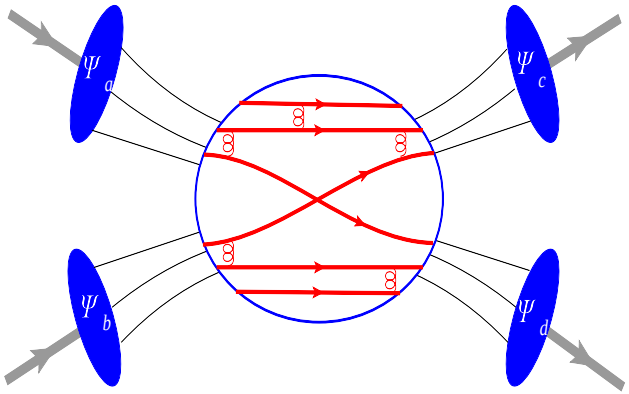
$a + b \to c + d$

- *Mandelstam variables*
  $s = (p_a + p_b)^2 = (p_c + p_d)^2$
  $t = (p_c - p_a)^2 = (p_d - p_b)^2$
  $u = (p_d - p_a)^2 = (p_c - p_b)^2$

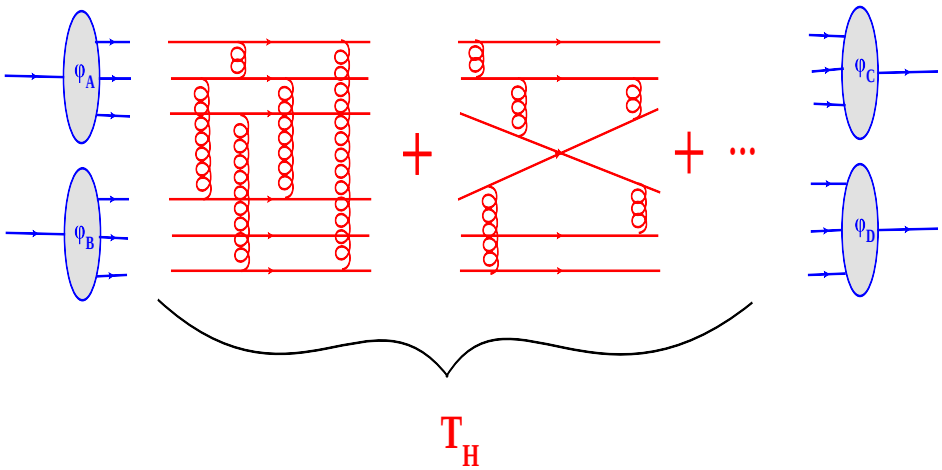
$ep \to eX$ *at large* Q^2 *revealed partons in nucleons*

Explore the role of quarks and gluons through Hard NN scattering





Factorization



$$
\begin{aligned}
\langle \psi_c \psi_d \mid T \mid \psi_a \psi_b \rangle \quad = \quad & \sum_{\alpha,\beta,\gamma} \langle \psi_c \mid \alpha_2', \beta_1', \gamma_1' \rangle \langle \psi_d \mid \alpha_1', \beta_2', \gamma_2' \rangle \\
& \times \langle \alpha_2', \beta_2', \gamma_2', \alpha_1' \beta_1' \gamma_1' \mid T_H(s,t) \mid \alpha_1, \beta_1, \gamma_1, \alpha_2, \beta_2, \gamma_2 \rangle \\
& \times \langle \alpha_1, \beta_1, \gamma_1 \mid \psi_a \rangle \langle \alpha_2, \beta_2, \gamma_2 \mid \psi_b \rangle,
\end{aligned}
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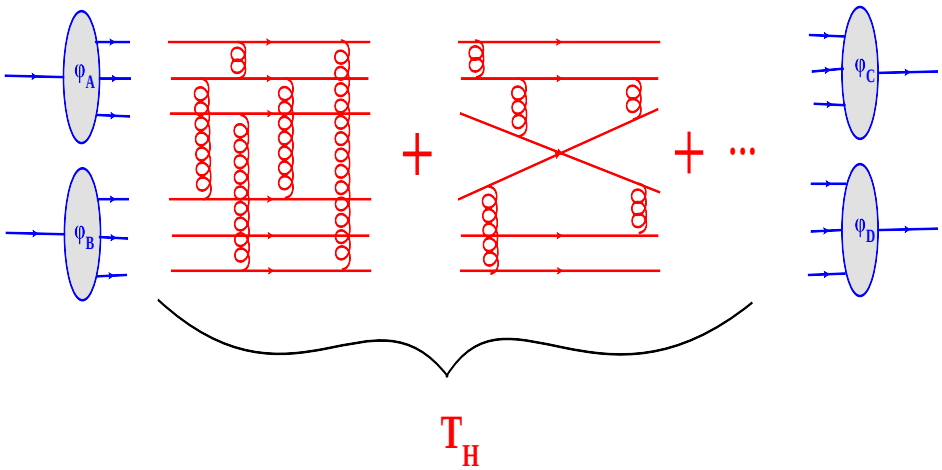
Factorization



- *Segment analysis in subprocesses according to scale.*
- *Hard subprocesses insensitive to large distance scale effects. Potentially workable through perturbative methods.*
- *Soft factors group long-distance scale effects. They're mostly approached phenomenologically*

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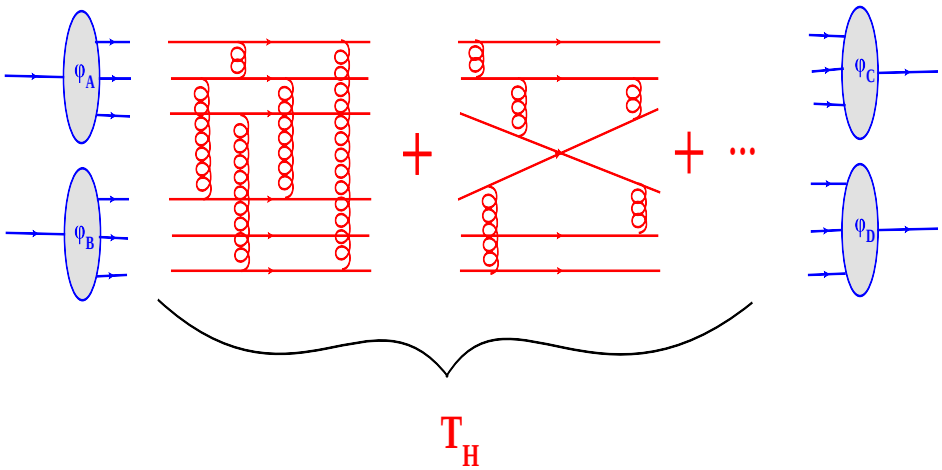

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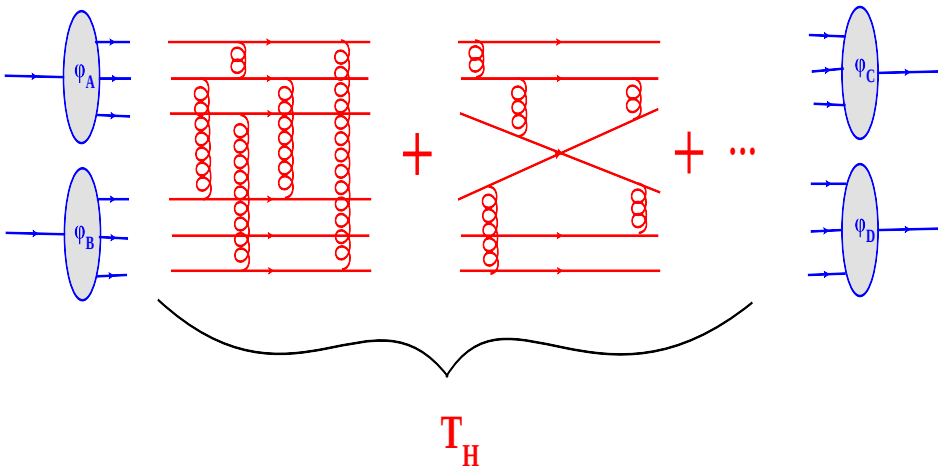

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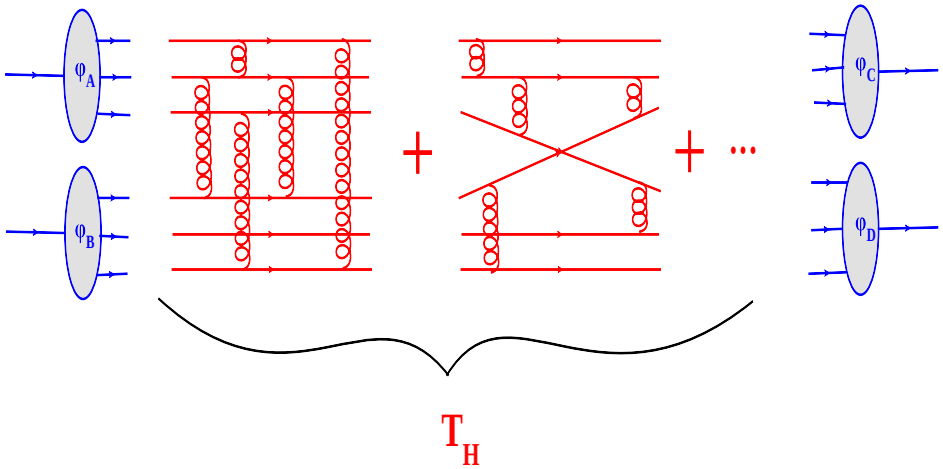
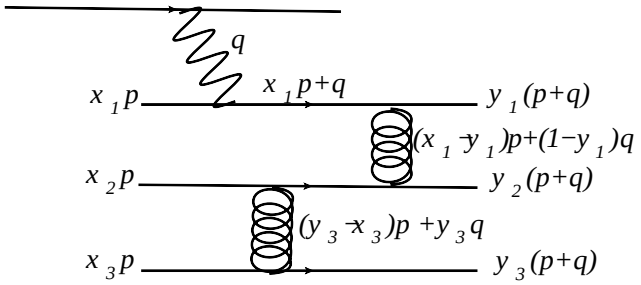
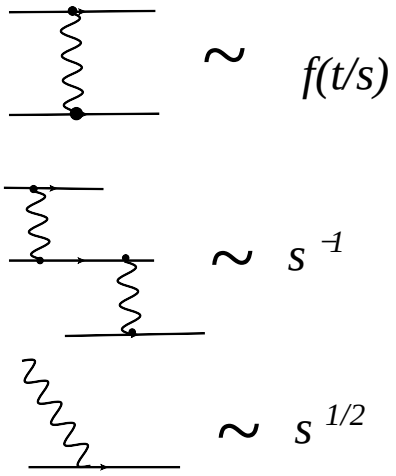
Dimensional Counting Rules





In exclusive reactions $a + b \to c + d$, *for hard subprocesses in the asymptotic limit* $(-t, s \to \infty$ *and fixed* $\frac{t}{s})$,



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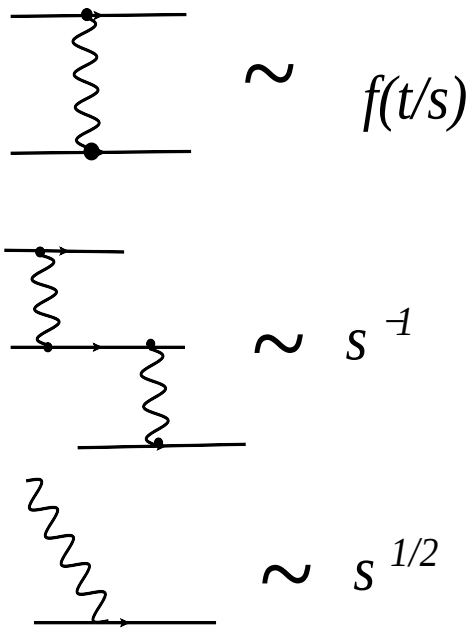

leading to,

$$\mathcal{M}_{ab\to cd} \sim s^{-\frac{n_a+n_b+n_c+n_d-4}{2}}$$

and consequently,

$$\frac{d\sigma}{dt}^{ab\to cd} \to \frac{|\mathcal{M}_{ab\to cd}|^2}{s^2}$$

$$\frac{d\sigma}{dt}^{ab\to cd} \to s^{-(n_a+n_b+n_c+n_d-2)}$$

with n_i *being the number of minimal constituents of particle* i *taking part in the hard subprocess.*

S.J. Brodsky and G.R. Farrar, Phys. Rev. Lett. 31, 1153 (1973)

V. Matveev, R.M. Muradyan and A.N. Tavkhelidze, Lett. Nuovo Cimento 7, 719 (1973).

Dimensional Counting Rules

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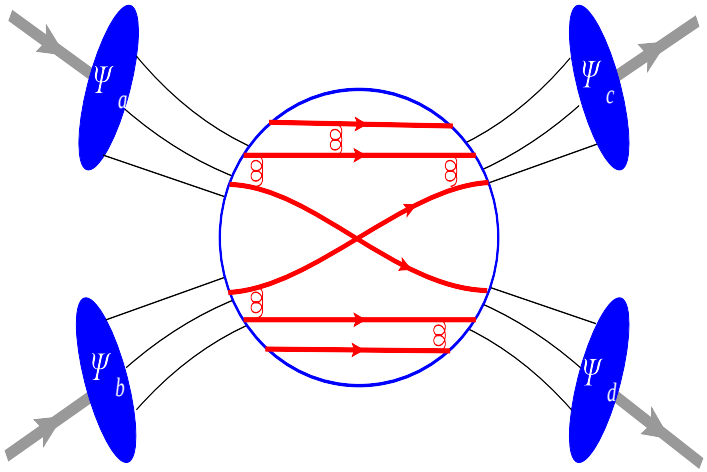
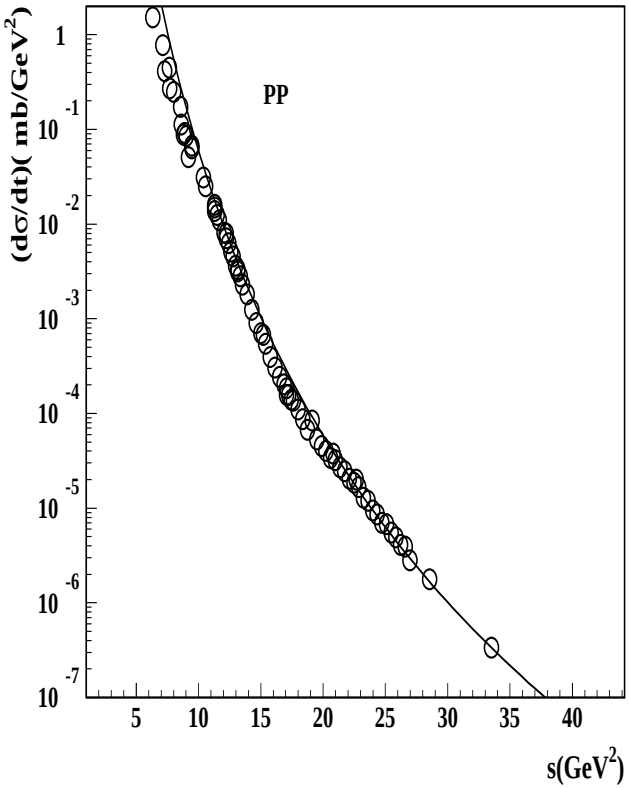
$NN \to NN$



$$\frac{d\sigma}{dt} \sim s^{-10}$$

Dimensional Counting Rules

pp elastic scattering



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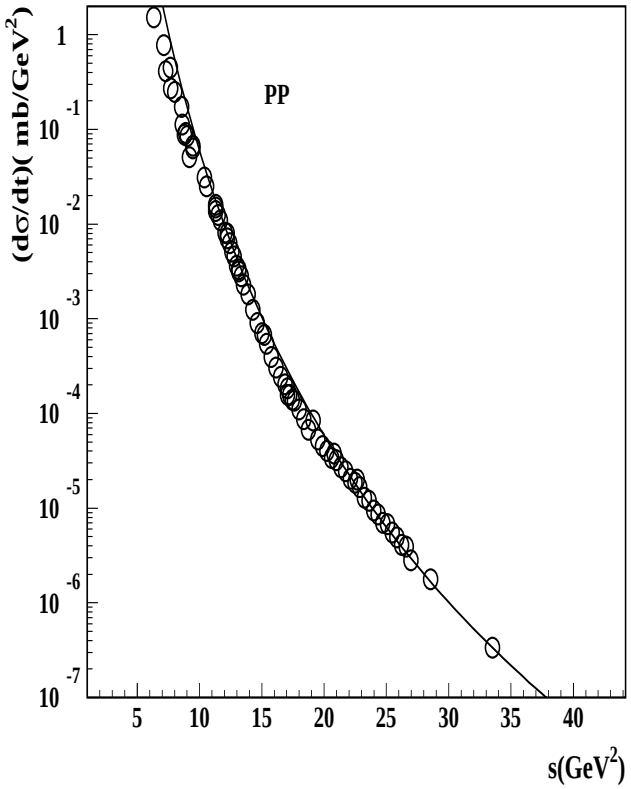
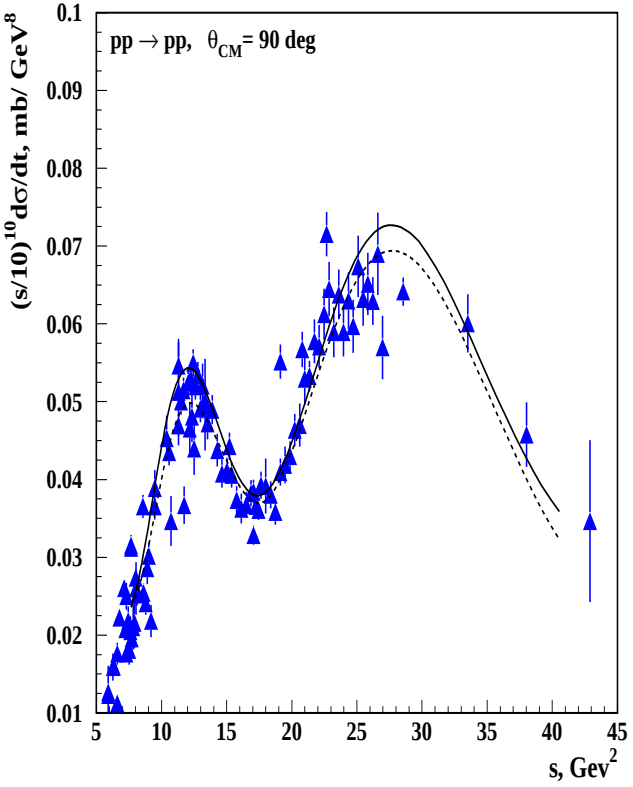
Dimensional Counting Rules

pp elastic scattering





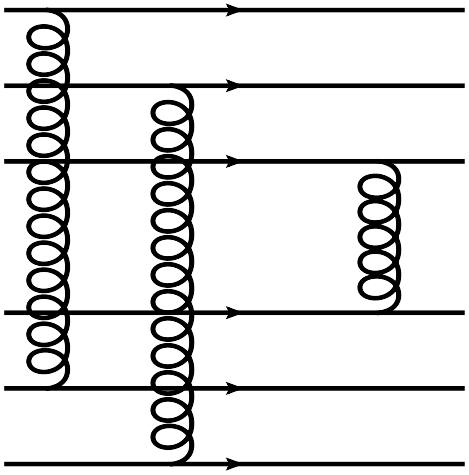
$$\frac{d\sigma}{dt} \sim s^{-10}$$

Oscillations point to a more complex picture of the NN interaction for the studied region.

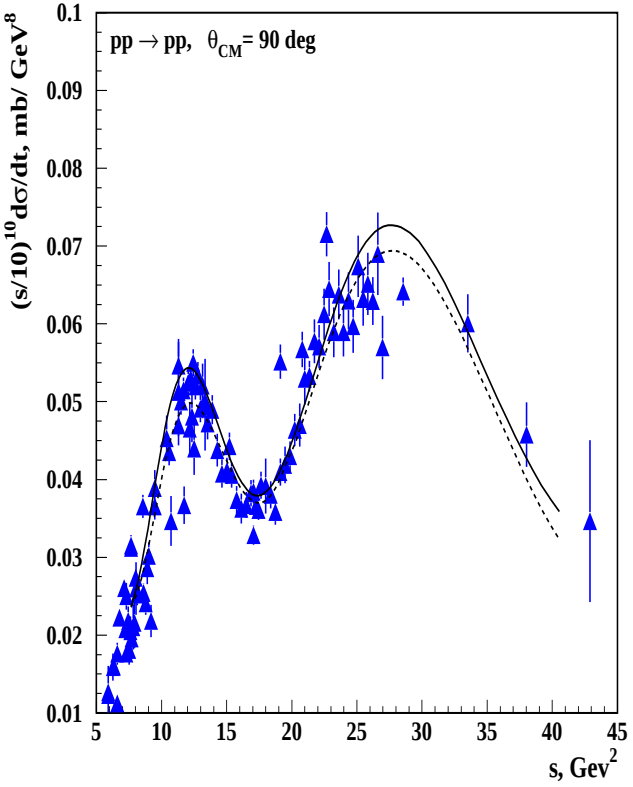
Dimensional Counting Rules

pp elastic scattering

- *Landshoff/Sudakof interference* (Pire,Ralston,1983)



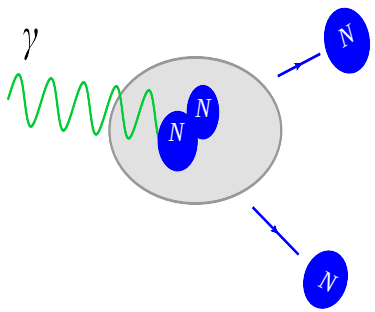
- *Heavy flavor resonances interfering with a pQCD background* (Brodsky,de Teramond,1988)



Oscillations point to a more complex picture of the NN interaction for the studied region.

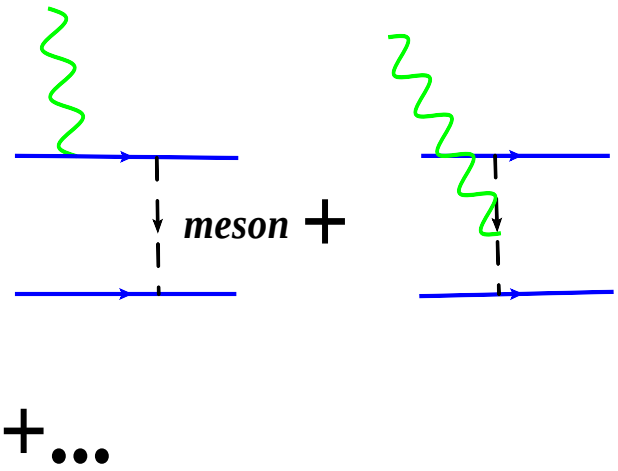
Hard Photodisintegration of a *NN* System in Nuclei I

Understanding of the strong force can be improved by studies of $\gamma + NN \to N + N$ *reactions*.

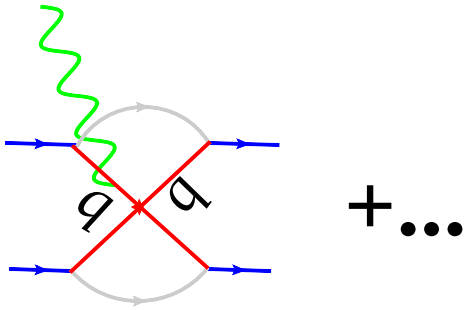


Hadronic and QCD descriptions of the NN force can be studied through such reactions.

Hadron picture



QCD picture



Kinematic Advantages of Hard $\gamma + NN \to N + N$

$N + N$ emerge at large transverse momentum from a short range NN interaction.

$$-t \sim_{s\to\infty} \frac{s}{2}(1 - cos(\theta_{c.m.}))$$

Large center of mass energy of the emerging nucleons $\sqrt{s}/2$ more efficiently reached in $\gamma + NN \to N + N$,

$$s_{\gamma NN} \approx 4m_N^2 + 2E_\gamma \cdot 2m_N,$$

compared to

$$s_{NN} = 2m_N^2 + 2E \cdot m_N$$

in $N_{beam} + N_{target} \longrightarrow N + N$ processes (with E being the energy of the nucleon beam).

Then, at moderate $E_\gamma \sim$GeV and at large $\theta_{c.m.}$, $\gamma + NN \to N + N$ reaches the hard kinematic regime.

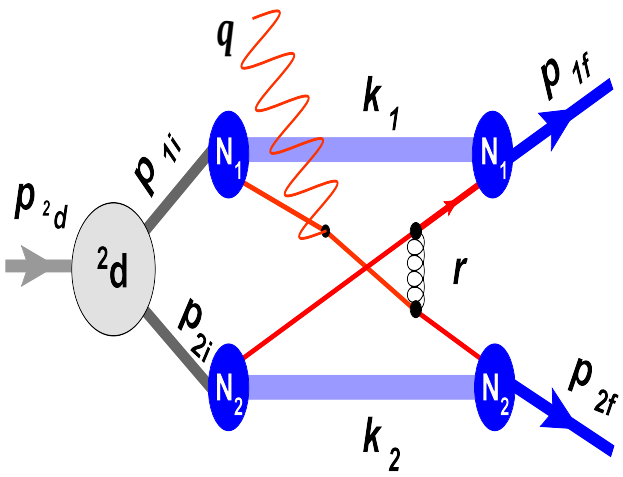
Producing High p_T $N + N$ in $\gamma + NN \to N + N$

There are two basic approaches,

- *Photon probes a preexisting compact NN system (large p_T component in initial state) in nucleus . This is assumed for instance in approaches such as the reduced nuclear amplitude (RNA)(Brodsky,Hiller,1983).*
  Or,
- *the energetic photon is absorbed by a low p_T NN system. The energy transfer triggers a final state interaction from which the nucleons emerge at large relative transverse.*

The latter is the scenario for the hard rescattering model (HRM) (Frankfurt et al. 2000) for which we develop applications in what follows.

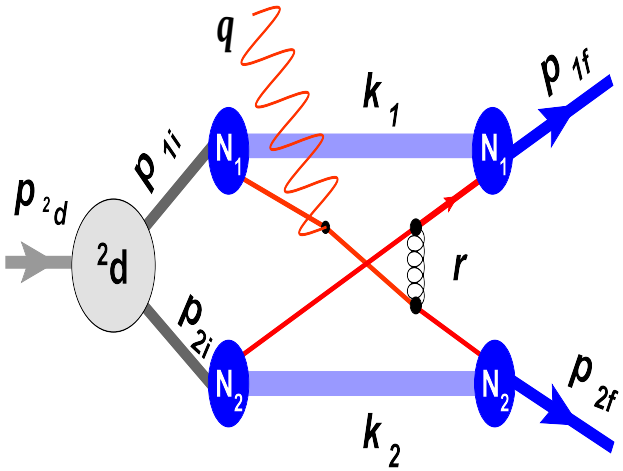
Hard rescattering model, HRM



L. Frankfurt et al.,Phys. Rev. Lett. 84, 3045 (2000)

▶

Hard rescattering model, HRM

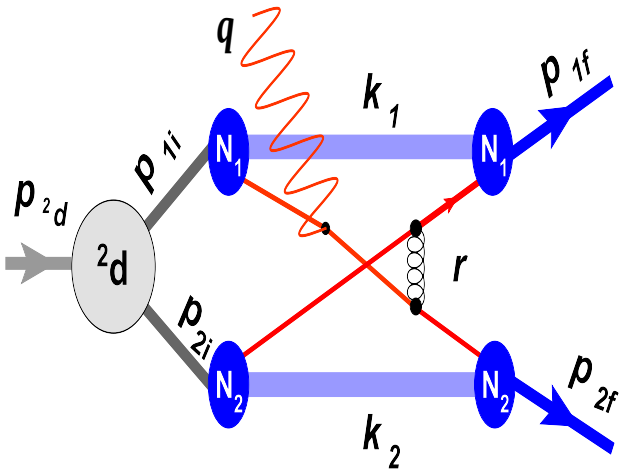


L. Frankfurt et al.,Phys. Rev. Lett. 84, 3045 (2000)

- $\gamma + d \to p + n$ *HRM amplitude results in a convolution of a hard process amplitude and a nuclear wave function,*

$$\langle \lambda_{1f}, \lambda_{2f} \mid \mathcal{M} \mid \lambda_\gamma, \lambda_d \rangle = -\frac{i[\lambda_\gamma] e Q_f \sqrt{2}(2\pi)^3}{\sqrt{2S'_{NN}}} \times \sum_{\lambda_{2i}} \int \langle \lambda_{2f}; \lambda_{1f} \mid T^{QIM}_{NN}(s_{NN}, t_N) \mid \lambda_\gamma; \lambda_{2i} \rangle$$
$$\times \quad \Psi^{\lambda_d \lambda_\gamma; \lambda_{2i}}_{d,NR}(\vec{p}_1, \vec{p}_2) m_N \frac{d^2 p_\perp}{(2\pi)^2},$$

Hard rescattering model, HRM



L. Frankfurt et al.,Phys. Rev. Lett. 84, 3045 (2000)

- *A parameter free cross section calculable through the input of experimental data,*

$$\frac{d\sigma^{\gamma d\to pn}}{dt}(s,\theta_{c.m.}) = \frac{8\alpha}{9}\pi^4\frac{1}{s'}C(\frac{\tilde{t}}{s}) \times \frac{\mathbf{d\sigma^{pn\to pn}}}{\mathbf{dt}}(\mathbf{s},\theta^{\mathbf{N}}_{\mathbf{c.m.}})$$
$$\times \left|\int \Psi_d^{NR}(p_z=0,p_t)\sqrt{m_n}\frac{d^2p_t}{(2\pi)^2}\right|^2$$

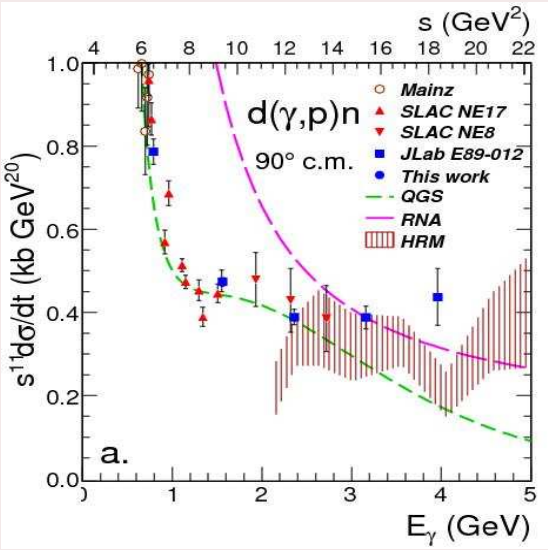
- *Energy and Angular distributions in pn elastic scattering data should be reflected by* $\gamma + d \longrightarrow p + n$ *data.*

Energy Distribution

$$\frac{d\sigma^{\gamma d\to pn}}{dt}(s,\theta_{c.m.}) = \frac{8\alpha}{9}\pi^4\frac{1}{s'}C(\frac{\tilde{t}}{s}) \times \frac{d\sigma^{pn\to pn}}{dt}(s,\theta^{N}_{c.m.})$$

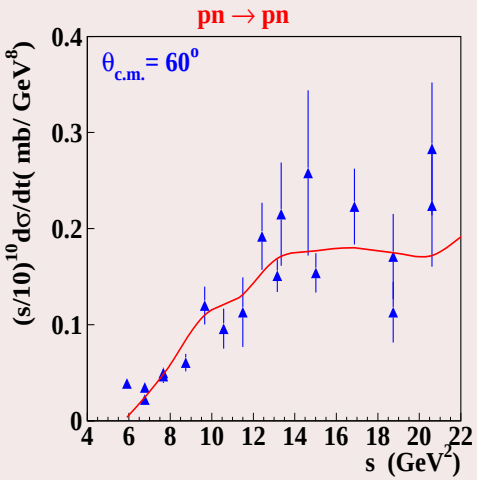
$$\times \left|\int \Psi_d^{NR}(p_z=0,p_t)\sqrt{m_n}\frac{d^2p_t}{(2\pi)^2}\right|^2$$

$\gamma + d \to p + n$



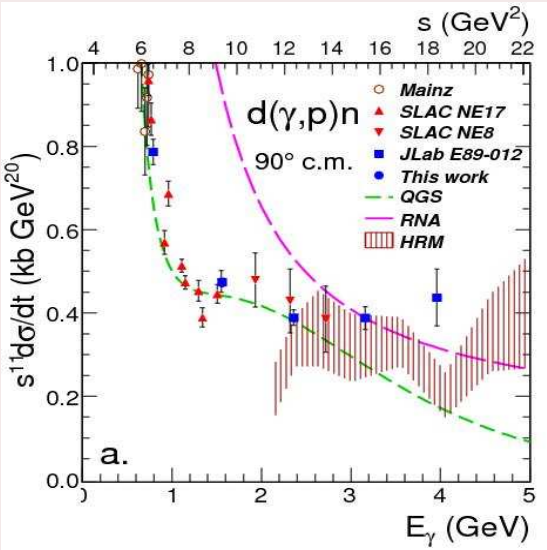
Bochna et al. 1998

$p + n \to p + n$



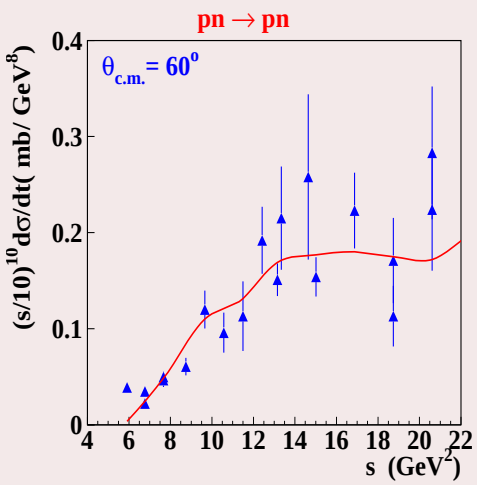
Energy Distribution

$\gamma + d \rightarrow p + n$



Bochna et al. 1998

$p + n \rightarrow p + n$

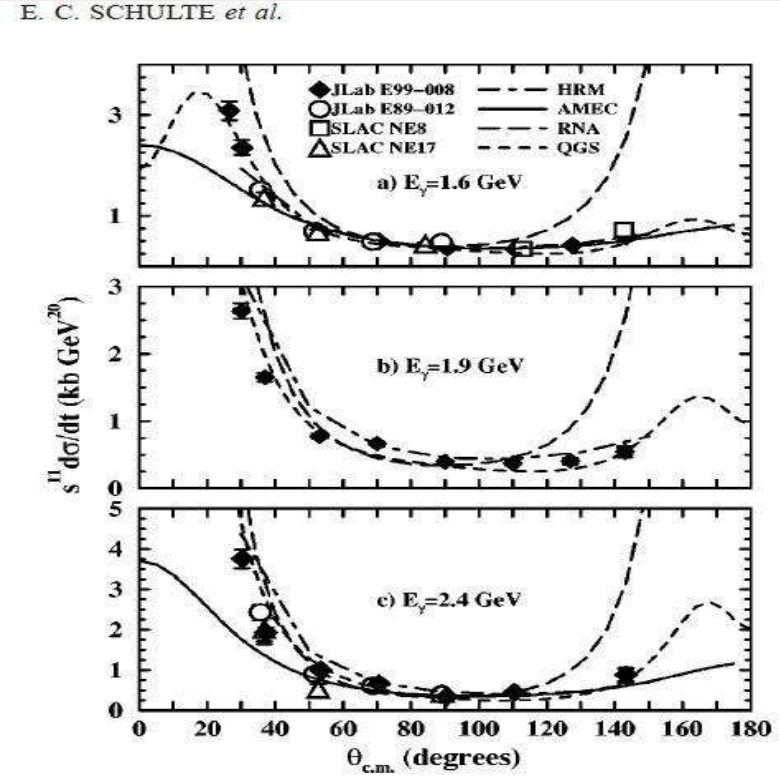


$\gamma + d \rightarrow p + n \sim s^{-11}$ ***for*** $E_\gamma > 2GeV$, $s > 10\boldsymbol{GeV}^2$

However, pn elastic scattering data is not good enough to clearly confirm HRM prediction

Angular Distributions

$\gamma + d \rightarrow p + n$



Schulte et al., Phys.Rev.C66:042201,2002.

$p + n \rightarrow p + n$

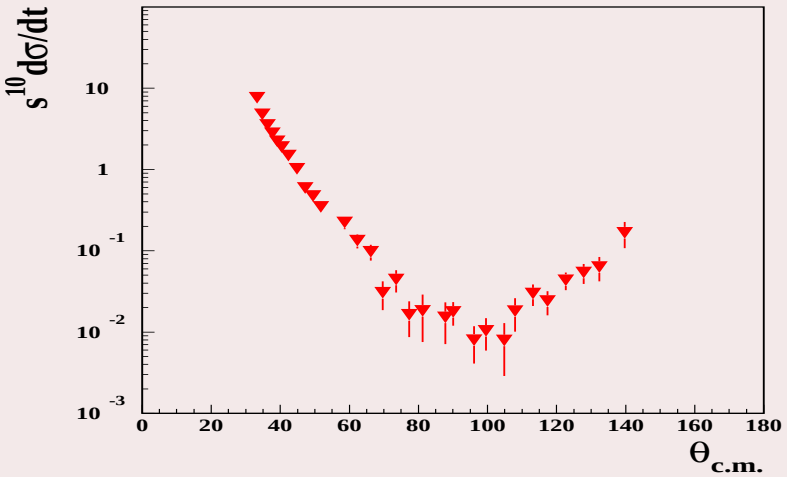


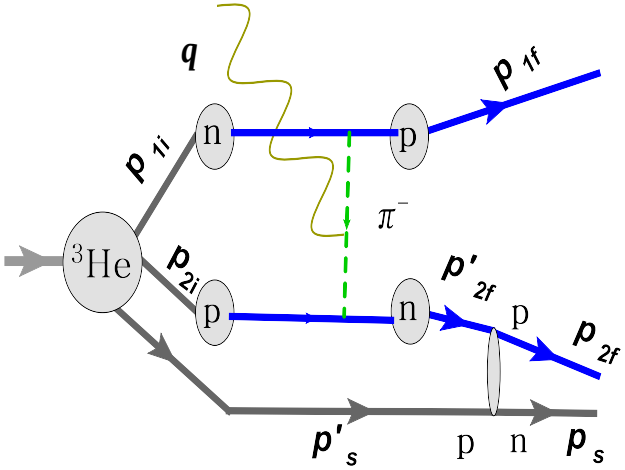
Figure: pn elastic scattering for p_{Lab}=8GeV/c

Hard Photodisintegration of a NN System in ^3He I

NN breakup studies can be extended to $\gamma + pp \to p + p$

- *Because of much better pp elastic scattering data, the HRM approach applied to pp photodisintegration can be tested experimentally better than the pn case.*
- *The energy behavior of $\gamma^3He \to (p+p) + n_s$ potentially probes the transition from a Hadronic to a QCD description of the NN force.*

At low beam energies $\gamma + pp \to p + p$ is suppressed because of the neutral charge of the mesons mediating the force. $\gamma^3He \to (p+p) + n_s$ proceeds largely through a three body (two steps) interaction:

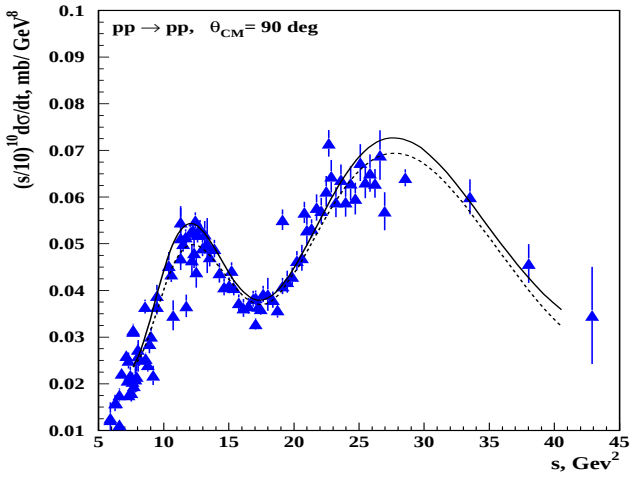


In hard pp breakup however, the photon couples to quark currents that now mediate the pp interaction making possible for

Hard Photodisintegration of a NN System in ^3He II

$\gamma^3He \to (p+p) + n_s$ *to proceed through* $\gamma + pp \to p + p$ *which becomes dominant.*

- *Features of* pp *elastic scattering such as oscillations in energy distributions are expected to appear as well in* $\gamma^3He \to (p+p) + n_s$.

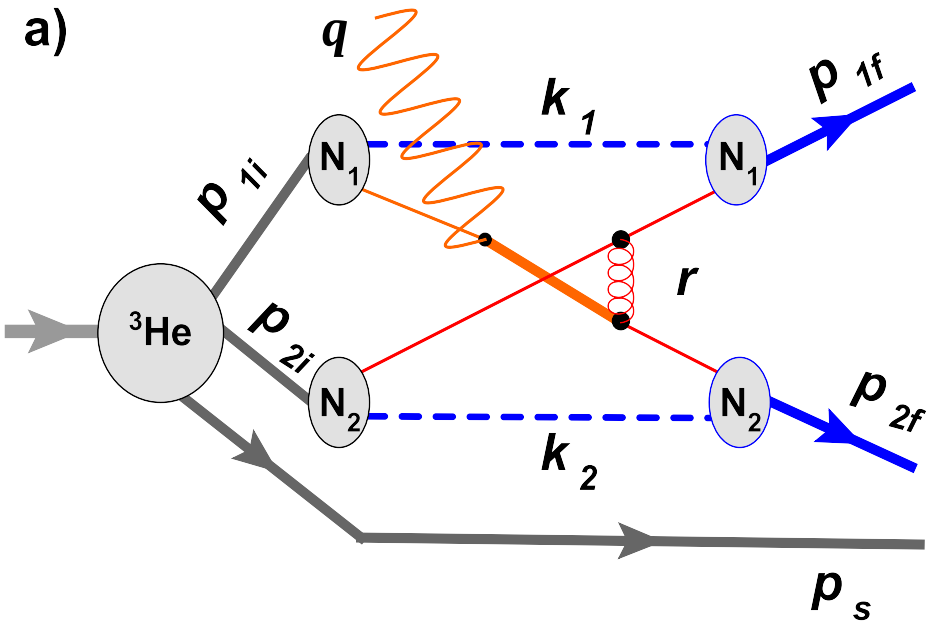


- *Predictions on spectator momentum distributions of* $\gamma^3He \to (p+p) + n_s$ *cross sections and on polarization observables provide yet more experimental checks on the validity of the HRM approach to* NN *breakup.*

^3He Photodisintegration within the HRM

Main assumptions

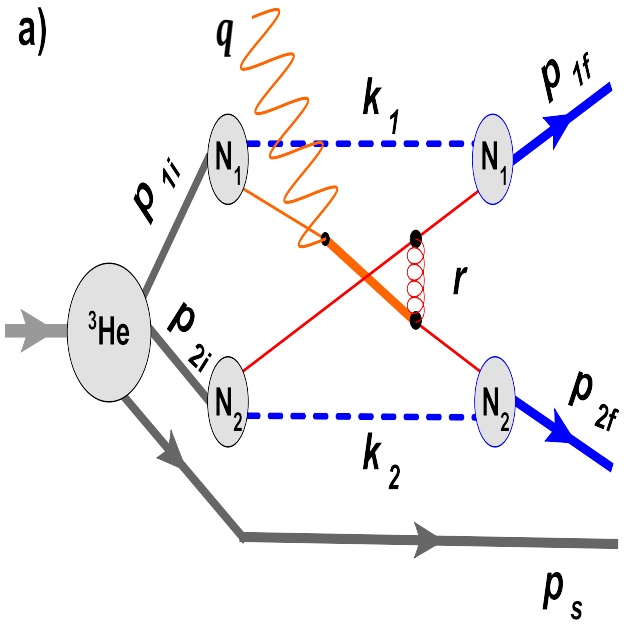
- *Photon interacts with a NN system with relative momentum* $p_\perp << m_N$
- *Third nucleon* (N_s) *is spectator to* γNN *reaction.*
- *The photon is absorbed by a valence quark of a nucleon from the NN system.*
- *Struck quark rescatters off valence quark of* 2nd *nucleon.*
- *This rescattering produces 2 nucleons emerging at large transverse momentum.*



S. J. Brodsky, et al., Phys. Lett. B578 (2004) 69.

M. M. Sargsian and C. Granados, Phys. Rev. C 80, 014612 (2009)

Calculation of Scattering Amplitude



$$\langle \lambda_{f1}, \lambda_{f2}, \lambda_s \mid \mathcal{M} \mid \lambda_\gamma, \lambda_A \rangle =$$

$$(N1): \int \frac{-i\Gamma^\dagger_{N_{1f}}\, i[\not{p}_{1f} - \not{k}_1 + m_q]}{(p_{1f} - k_1)^2 - m_q^2 + i\epsilon} \cdots [-igT_c^F \gamma_\mu] \cdots$$

$$iS(k_1) \frac{i[\not{p}_{1i} - \not{k}_1 + m_q](-i)\Gamma_{N_{1i}}}{(p_{1i} - k_1)^2 - m^2 + i\epsilon} \frac{d^4 k_1}{(2\pi)^4}$$

$$(\gamma q): \frac{i[\not{p}_{1i} - \not{k}_1 + \not{q} + m_q]}{(p_{1i} - k_1 + q)^2 - m_q^2 + i\epsilon} [-iQ_i e \epsilon^\perp \gamma^\perp]$$

$$(N2): \int \frac{-i\Gamma^\dagger_{N_{2f}}\, i[\not{p}_{2f} - \not{k}_2 + m_q]}{(p_{2f} - k_2)^2 - m_q^2 + i\epsilon} \cdots [-igT_c^F \gamma_\nu]$$

$$iS(k_2) \frac{i[\not{p}_{2i} - \not{k}_2 + m_q](-i)\Gamma_{N_{2i}}}{(p_{2i} - k_2)^2 - m^2 + \epsilon} \frac{d^4 k_2}{(2\pi)^4}$$

$$(^3He): \int \frac{\bar{u}_{\lambda_s}(p_s) i[\not{p}_{NN} - \not{p}_{2i} + m_N]}{(p_{NN} - p_{2i})^2 - m_N^2 + i\epsilon} \frac{i[\not{p}_{2i} + m_N]}{p_{2i}^2 - m_N^2 + i\epsilon} (-i)\Gamma_{^3\mathrm{He}} \frac{d^4 p_{2i}}{(2\pi)^4}$$

$$(g): \frac{i d^{\mu,\nu} \delta_{ab}}{[(p_{2i} - k_2) - (p_{1i} - k_1) - (q - l)]^2 + i\epsilon}, \tag{3}$$

Light cone variables

$$p = (p^+, p^-, p_\perp)$$

with

$$p^+ = p_0 + p_z; \quad p^- = p_0 - p_z; \quad p_\perp = (p_x, p_y).$$

Light cone momentum fractions

$$\alpha = \frac{p_{2i}^+}{p_{NN}^+},$$

and

$$x_l = \frac{k_l^+}{p_{li}^+}; \quad x_l' = \frac{k_l^+}{p_{lf}^+}.$$

Then,

$$d^4 p_{2i} = \frac{1}{2} p_{NN}^+ d\alpha dp_{2i}^- d^2 p_\perp,$$

and

$$d^4 k_l = \frac{1}{2} p_{li}^+ dx dk_l^- d^2 k_\perp$$

$\gamma - NN$ Kinematics

$\gamma - NN$ Reference Frame

We consider a reference frame where $q^+ = 0$. Having that $s_{NN} = (p_{NN} + q)^2$ and $s'_{NN} = s_{NN} - M^2_{NN}$, the 4-momenta of the photon and of the NN system are respectively:

$$q \equiv (q^+, q^-, q_\perp) = (0, \sqrt{s'_{NN}}, 0),$$

$$p \equiv (p^+_{NN}, p^-_{NN}, q_{NN\perp}) = (\sqrt{s'_{NN}}, \frac{M^2_{NN}}{\sqrt{s'_{NN}}}, 0).$$

For $p^+_{NN} >> M^2_{NN}$, the $\gamma - NN$ system approaches its center of mass reference frame.

HRM Scattering Amplitude for $\gamma +^3 He \to N + N + N_s$

After performing loop integrations on p_{21}^- and k^- variables, N_2 and spectator systems are put on mass shell. Using that,

$$\not{p} + m \approx \sum_s u_s(p)\bar{u}_s(p),$$

for internal lines off mass-shell in the $p_{NN}^+ >> M_{NN}^2$ approximation and by defining nuclear and nucleonic wave functions,

$$\Psi_{^3\mathrm{He}}^{\lambda_A,\lambda_1,\lambda_2,\lambda_s}(\alpha, p_\perp) = \frac{\bar{u}_{\lambda_1}(p_{NN} - p)\bar{u}_{\lambda_2}(p)\bar{u}_{\lambda_s}(p_s)\Gamma_{^3\mathrm{He}}^{\lambda_A}}{M_{NN}^2 - \frac{m_N^2 + p_\perp^2}{\alpha(1-\alpha)}} \tag{4}$$

$$\Psi_N^{\lambda,\eta}(p, x, k_\perp) = \frac{\bar{u}_\eta(p-k)\psi_s^\dagger(k)\Gamma_N u_N^\lambda(p)}{m_N^2 - \frac{m_s^2(1-x) + m_q^2 x + (k_\perp - x p_\perp)^2}{x(1-x)}} \tag{5}$$

HRM Scattering Amplitude for $\gamma +^3 He \to N + N + N_s$

After performing loop integrations on p_{21}^- and k^- variables, N_2 and spectator systems are put on mass shell. Using that,

$$\not{p} + m \approx \sum_s u_s(p)\bar{u}_s(p),$$

for internal lines off mass-shell in the $p_{NN}^+ >> M_{NN}^2$ approximation and by defining nuclear and nucleonic wave functions,
the invariant amplitude is simplified to:

$$\langle \lambda_{1f}, \lambda_{2f}, \lambda_s \mid \mathcal{M} \mid \lambda_\gamma, \lambda_A \rangle = \sum_{(\eta' s),(\lambda_i' s),(\zeta)} \int \left\{ \frac{\psi_N^{\dagger \lambda_{2f}, \eta_{2f}}(p_{2f}, x_2', k_{2\perp})}{1 - x_2'} \bar{u}_{\eta_{2f}}(p_{2f} - k_2)[-igT_c^F \gamma^\nu] \right.$$

$$\left. u_\zeta(p_{1i} - k_1 + q) \frac{\bar{u}_\zeta(p_{1i} - k_1 + q)[-iQ_i e \epsilon_\perp^{\lambda_\gamma} \gamma^\perp] u_{\eta_{1i}}(p_{1i} - k_1)}{(1 - x_1) s'(\alpha - (\alpha_c + i\epsilon))} \frac{\psi_N^{\lambda_{1i}, \eta_{1i}}(p_{1i}, x_1, k_{1\perp})}{(1 - x_1)} \right\}_1 \times$$

$$\left\{ \frac{\psi_N^{\dagger \lambda_{1f}, \eta_{1f}}(p_{1f}, x_1', k_{1\perp})}{1 - x_1'} \bar{u}_{\eta_{1f}}(p_{1f} - k_1)[-igT_c^F \gamma^\mu] u_{\eta_{2i}}(p_{2i} - k_2) \frac{\psi_N^{\lambda_{2i}, \eta_{2i}}(p_{2i}, x_2, k_{2\perp})}{(1 - x_2)} \right\}_2 \times$$

$$G^{\mu,\nu}(r) \frac{dx_1}{x_1} \frac{d^2 k_{1\perp}}{2(2\pi)^3} \frac{dx_2}{x_2} \frac{d^2 k_{2\perp}}{2(2\pi)^3} \frac{\Psi_{^3\mathrm{He}}^{\lambda_A, \lambda_{1i}, \lambda_{2i}, \lambda_s}(\alpha, p_\perp, p_s)}{(1 - \alpha)} \frac{d\alpha}{\alpha} \frac{d^2 p_\perp}{2(2\pi)^3} - (p_{1f} \longleftrightarrow p_{2f}), \quad (4)$$

γ-quark hard scattering factor

$$\frac{\bar{u}_\zeta(p_{1i}-k_1+q)[-iQ_i e\epsilon_\perp^{\lambda_\gamma}\gamma^\perp]u_{\eta_{1i}}(p_{1i}-k_1)}{(1-x_1)s'(\alpha-(\alpha_c+i\epsilon))}$$

can be calculated in the approximation $p_{1i}^+ - k_1^+ >> k_\perp, m_q$, *in which the spinor of a quark of a given helicity* $\eta_q = \pm 1$ *and a given energy* E_q *can be written*

$$u_{\eta_q} \approx \sqrt{E_q}\begin{pmatrix} \chi_{\eta_q} \\ \eta_q\chi_{\eta_q} \end{pmatrix}.$$

in which χ *are Pauli spinors.*

Recall that $x_1 = \frac{k_1^+}{p_{1i}^+}$, $\alpha = \frac{p_{2i}^+}{p_{NN}^+}$ *and having that,*

$$\alpha_c = 1 + \frac{1}{s'_{NN}}\left[\tilde{m}_N^2 - \frac{m_s^2(1-x_1)+m_q^2x_1+(k_1-x_1p_1)_\perp^2}{x_1(1-x_1)}\right]$$

Then using also that $\frac{1}{\alpha-\alpha_c+i\epsilon} \approx -i\pi\delta(\alpha-\alpha_c) + P(\frac{1}{\alpha-\alpha_c})$ *and neglecting the principal value term we obtain,*

$$\frac{\bar{u}_\zeta[-iQ_i e\epsilon_\perp^{\lambda_\gamma}\gamma^\perp]u_{\eta_{1i}}}{(1-x_1)s'(\alpha-(\alpha_c+i\epsilon))} \approx i\pi\delta(\alpha-\alpha_c)Q_i e\lambda_\gamma\sqrt{\frac{2(1-(1-x_1)(1-\alpha))}{s'(1-x_1)}}\delta_{\zeta\lambda_\gamma}\delta_{\eta_{1i}\lambda_\gamma}$$

then γ *only interacts with a quark having its same helicity, i.e.,if* $\eta_{1i} = \lambda_\gamma$ *and consequently we have that* $\zeta = \eta_{1i}$

α selection in $\gamma +^3 He \to (N+N) + N_s$

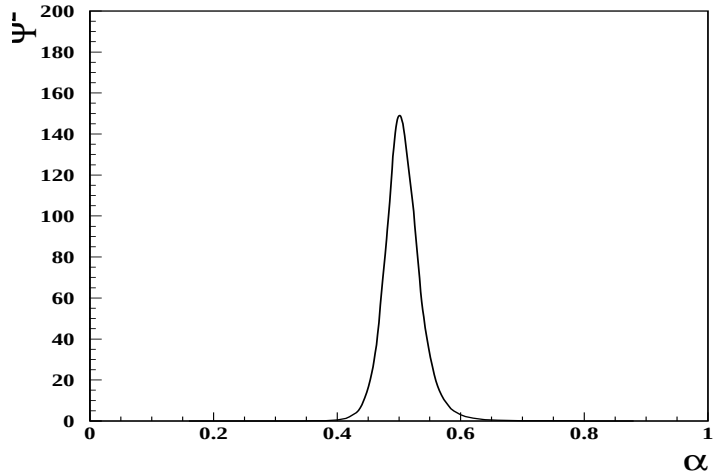
$$\frac{\bar{u}_\zeta[-iQ_i e\epsilon_\perp^{\lambda_\gamma}\gamma^\perp]u_{\eta_{1i}}}{(1-x_1)s'(\alpha-(\alpha_c+i\epsilon))} \approx i\pi\delta(\alpha-\alpha_c)Q_i e\lambda_\gamma\sqrt{\frac{2(1-(1-x_1)(1-\alpha))}{s'(1-x_1)}}\delta_{\zeta\lambda_\gamma}\delta_{\eta_{1i}\lambda_\gamma}$$

allows to perform the α integration in calculating $\mathcal{M}$.

α_c is then approximated to $\frac{1}{2}$ which corresponds to the situation in which both nucleons in NN contribute equally to p^+_{NN}. This is justified by the strong peak of $\Psi_{^3He}$'s magnitude at $\alpha = \frac{1}{2}$.

From

$$\alpha_c = 1 + \frac{1}{s'_{NN}}\left[\tilde{m}_N^2 - \frac{m_s^2(1-x_1)+m_q^2x_1+(k_1-x_1p_1)_\perp^2}{x_1(1-x_1)}\right].$$



in the large s' limit, $\alpha_c = 1/2$ leads to $x_1 \sim k_{1\perp}^2/s'$. Because $k_{1\perp}^2 << s'$, this corresponds to having the valence quark contributing to most of the momentum of its parent nucleon $\Rightarrow$ $p_q \approx p_N$ and $\eta_q \approx \lambda_N$.

α selection in $\gamma +^3 He \to (N+N) + N_s$

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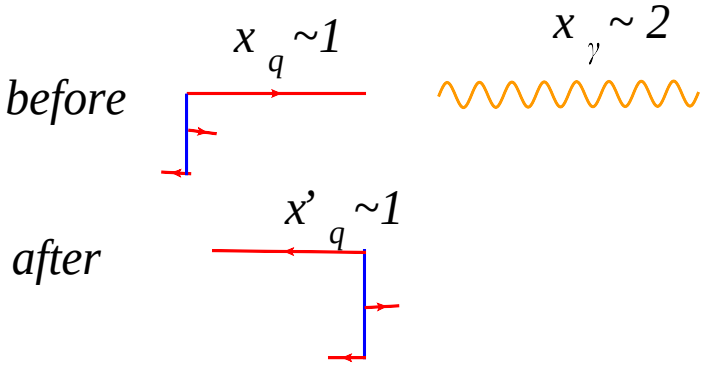
$$\alpha_c = 1 + \frac{1}{s'_{NN}}\left[\tilde{m}_N^2 - \frac{m_s^2(1-x_1)+m_q^2 x_1+(k_1-x_1p_1)_\perp^2}{x_1(1-x_1)}\right].$$

in the large s' limit, $\alpha_c = 1/2$ leads to $x_1 \sim k_{1\perp}^2/s'$. Because $k_{1\perp}^2 << s'$, this corresponds to having the valence quark contributing to most of the momentum of its parent nucleon $\Rightarrow$ $p_q \approx p_N$ and $\eta_q \approx \lambda_N$.



HRM scattering amplitude of $\gamma + He \to N + N + N_s$

After the considered approximations are applied to $\mathcal{M}$, the scattering amplitude for $\gamma + He \to N + N + N_s$ takes the form,

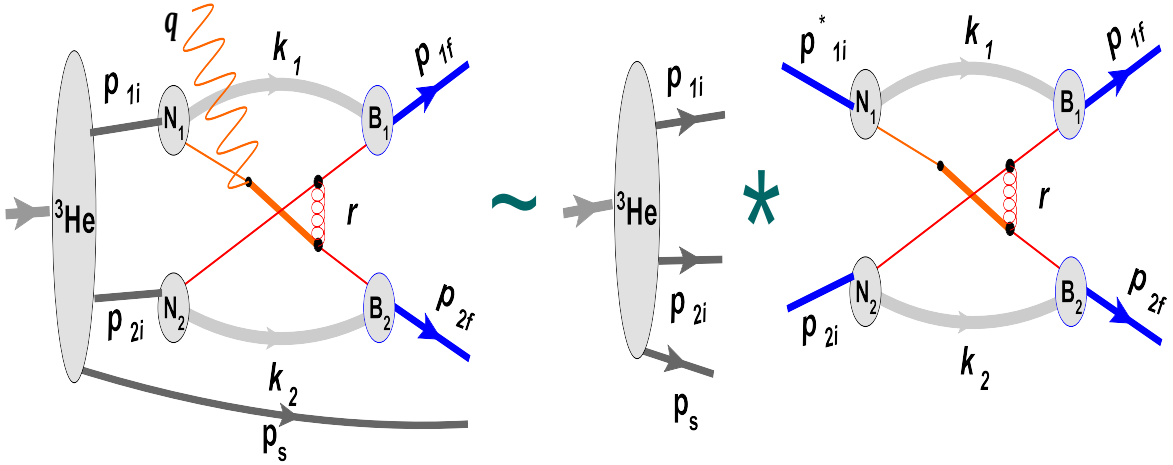
$$\langle \lambda_{1f}, \lambda_{2f} \lambda_s \mid \mathcal{M}_i \mid \lambda_\gamma, \lambda_A \rangle = i[\lambda_\gamma] e \sum_{(\eta_{1f},\eta_{2f}),(\eta_{2i}),\lambda_{1i},\lambda_{2i}} \int \frac{Q_i}{\sqrt{2s'}}$$

$$\left[\left\{ \frac{\psi_N^{\dagger \lambda_{2f},\eta_{2f}}}{1-x_2'} \bar{u}_{\eta_{2f}} [-igT_c^F \gamma^\nu] \, u_{\eta_{1i}} \frac{\psi_N^{\lambda_{1i},\eta_{1i}}}{(1-x_1)} \right\} \times \right.$$

$$\left\{ \frac{\psi_N^{\dagger \lambda_{1f},\eta_{1f}}}{1-x_1'} \bar{u}_{\eta_{1f}} [-igT_c^F \gamma^\mu] u_{\eta_{2i}} \frac{\psi_N^{\lambda_{2i},\eta_{2i}}}{(1-x_2)} \right\}$$

$$\left. G^{\mu,\nu} \frac{dx_1}{x_1} \frac{d^2 k_{1\perp}}{2(2\pi)^3} \frac{dx_2}{x_2} \frac{d^2 k_{2\perp}}{2(2\pi)^3} \right] \Psi_{^3He}^{\lambda_A,\lambda_{1i},\lambda_{2i}} (\alpha = \frac{1}{2}, p_{2\perp}) \frac{d^2 p_{2\perp}}{(2\pi)^2}.$$

which is written as a convolution of a NN elastic scattering amplitude with a component of a nuclear wave function.



HRM scattering amplitude of $\gamma + He \to N + N + N_s$

Accounting for all possible quark interchange diagrams we obtain,

$$\langle \lambda_{1f}, \lambda_{2f}, \lambda_s \mid M \mid \lambda_\gamma, \lambda_A \rangle = \frac{i[\lambda_\gamma] e \sqrt{2} (2\pi)^3}{\sqrt{2S'_{NN}}} \times$$

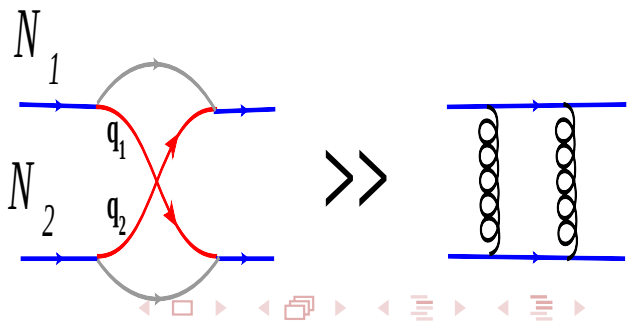
$$\left\{ Q_F^{N_1} \sum_{\lambda_{2i}} \int \langle \lambda_{2f}; \lambda_{1f} \mid T_{NN}^{\mathrm{QIM}}(s_{NN}, t_N) \mid \lambda_\gamma; \lambda_{2i} \rangle \Psi^{\lambda_A}_{{}^3\mathrm{He,NR}}(\vec{p}_1, \lambda_\gamma; \vec{p}_2, \lambda_{2i}; \vec{p}_s, \lambda_s) m_N \frac{d^2 p_\perp}{(2\pi)^2} + \right.$$

$$\left. Q_F^{N_2} \sum_{\lambda_{1i}} \int \langle \lambda_{2f}; \lambda_{1f} \mid T_{NN}^{\mathrm{QIM}}(s_{NN}, t_N) \mid \lambda_{1i}; \lambda_\lambda \rangle \Psi^{\lambda_A}_{{}^3\mathrm{He,NR}}(\vec{p}_1, \lambda_{1i}; \vec{p}_2, \lambda_\gamma; \vec{p}_s, \lambda_s) m_N \frac{d^2 p_\perp}{(2\pi)^2} \right\} \tag{5}$$

where charge factors Q_F are introduced such that,

$$\sum_{i \in N} Q_i^N \langle a' b' | T_{NN,i}^{\mathrm{QIM}} | ab \rangle = Q_F^N \cdot \langle a' b' | T_{NN}^{\mathrm{QIM}} | ab \rangle$$

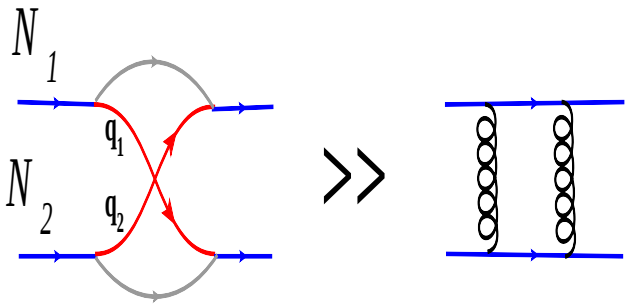
Dominance of the quark interchange (QI) mechanism (White et al. 1994)



$\gamma +^3 He \to N + N + N_s$ Scattering Amplitudes

NN scattering amplitudes I

Dominance of the quark interchange (QI) mechanism (*White et al. 1994*)



For the quark interchange NN amplitudes, we choose NN helicity amplitudes labeled as follows:

$$
\begin{aligned}
< +,+|T_{NN}^{\mathrm{QIM}}|+,+> &= \phi_1 \\
< +,+|T_{NN}^{\mathrm{QIM}}|+,-> &= \phi_5 \\
< +,+|T_{NN}^{\mathrm{QIM}}|-,-> &= \phi_2 \\
< +,-|T_{NN}^{\mathrm{QIM}}|+,-> &= \phi_3 \\
< +-|T_{NN}^{\mathrm{QIM}}|-+> &= -\phi_4. \qquad (6)
\end{aligned}
$$

All other helicity combinations can be related to the above amplitudes through the parity and time-reversal symmetry. The cross section of a NN scattering is defined as

$$
\frac{d\sigma^{NN\to NN}}{dt} = \frac{1}{16\pi}\frac{1}{s(s-4m_N^2)}\frac{1}{2}(|\phi_1|^2+|\phi_2|^2+|\phi_3|^2+|\phi_4|^2+4|\phi_5|^2). \qquad (7)
$$

$\gamma +^3 He \to N + N + N_s$ Scattering Amplitudes

NN scattering amplitudes II

Then for NN breakup in ^3He*, using the antisymmetry of the* ^3He *ground state wave function we have,*

$$\langle +,+,\lambda_s \mid \mathcal{M} \mid +,\lambda_A \rangle = B \int \left[Q_F \phi_5 \Psi^{\lambda_A}_{^3He}(+,-,\lambda_s) + Q_F \phi_1 \Psi^{\lambda_A}_{^3He}(+,+,\lambda_s) \right] m_N \frac{d^2 p_\perp}{(2\pi)^2}$$

$$\langle +,-,\lambda_s \mid \mathcal{M} \mid +,\lambda_A \rangle = B \int \left[(Q_F^{N_1} \phi_3 + Q_F^{N_2} \phi_4) \Psi^{\lambda_A}_{^3He}(+,-,\lambda_s) - Q_F \phi_5 \Psi^{\lambda_A}_{^3He}(+,+,\lambda_s) \right] \times m_N \frac{d^2 p_\perp}{(2\pi)^2}$$

$$\langle -,+,\lambda_s \mid \mathcal{M} \mid +,\lambda_A \rangle = B \int \left[-(Q_F^{N_1} \phi_4 + Q_F^{N_2} \phi_3) \Psi^{\lambda_A}_{^3He}(+,-,\lambda_s) + Q_F \phi_5 \Psi^{\lambda_A}_{^3He}(+,+,\lambda_s) \right] \times m_N \frac{d^2 p_\perp}{(2\pi)^2}$$

$$\langle -,-,\lambda_s \mid \mathcal{M} \mid +,\lambda_A \rangle = B \int \left[Q_F \phi_5 \Psi^{\lambda_A}_{^3He}(+,-,\lambda_s) + Q_F \phi_2 \Psi^{\lambda_A}_{^3He}(+,+,\lambda_s) \right] m_N \frac{d^2 p_\perp}{(2\pi)^2}$$

where $B = \frac{ie\sqrt{2}(2\pi)^3}{\sqrt{2s'_{NN}}}$.

NN Breakup cross sections

Then for the averaged squared amplitude one obtains,

$$|\bar{\mathcal{M}}|^2 = \frac{e^2 2(2\pi)^6}{2s'_{NN}} \frac{1}{2} \left\{ 2Q_F^2 |\phi_5|^2 S_0 + Q_F^2 (|\phi_1|^2 + |\phi_2|^2) S_{12} + \left[(Q_F^{N_1}\phi_3 + Q_F^{N_2}\phi_4)^2 + (Q_F^{N_1}\phi_4 + Q_F^{N_2}\phi_3)^2 \right] S_{34} \right\} \tag{8}$$

where $Q_F = Q_F^{N_1} + Q_F^{N_2}$ *and* S_{12}, S_{34}, *and* S_0 *are partially integrated nuclear spectral functions:*

$$S_{12}(t_1, t_2, \alpha, \vec{p}_s) = N_{NN} \sum_{\lambda_1=\lambda_2=-\frac{1}{2}}^{\frac{1}{2}} \sum_{\lambda_3=-\frac{1}{2}}^{\frac{1}{2}} \left| \int \Psi^{\frac{1}{2}}_{^3\mathrm{He,NR}}(\vec{p}_1, \lambda_1, t_1; \vec{p}_2, \lambda_2, t_2; \vec{p}_s, \lambda_3) m_N \frac{d^2 p_\perp}{(2\pi)^2} \right|^2,$$

$$S_{34}(t_1, t_2, \alpha, \vec{p}_s) = N_{NN} \sum_{\lambda_1=-\lambda_2=-\frac{1}{2}}^{\frac{1}{2}} \sum_{\lambda_3=-\frac{1}{2}}^{\frac{1}{2}} \left| \int \Psi^{\frac{1}{2}}_{^3\mathrm{He,NR}}(\vec{p}_1, \lambda_1, t_1; \vec{p}_2, \lambda_2, t_2; \vec{p}_s, \lambda_3) m_N \frac{d^2 p_\perp}{(2\pi)^2} \right|^2$$

and

$$S_0 = S_{12} + S_{34}. \tag{11}$$

NN Breakup cross sections

Assuming the massless approximation for the interchanging quarks, their corresponding helicities are conserved during the hard subprocess, leading to the vanishing of NN amplitudes that don't conserve helicity, i.e.,

$$\begin{aligned} \phi_2 &= 0 \\ \phi_5 &= 0 \end{aligned}$$

then for NN scattering,

$$\frac{d\sigma^{NN\to NN}}{dt} = \frac{1}{16\pi}\frac{1}{s(s-4m_N^2)}\frac{1}{2}(|\phi_1|^2+|\phi_3|^2+|\phi_4|^2), \tag{12}$$

while for NN breakup

$$\begin{aligned} |\bar{\mathcal{M}}|^2 = \quad & \frac{e^2 2(2\pi)^6}{2s'_{NN}}\frac{1}{2}\left\{Q_F^2(|\phi_1|^2)S_{12}+\right. \\ & \left.\left[(Q_F^{N_1}\phi_3+Q_F^{N_2}\phi_4)^2+(Q_F^{N_1}\phi_4+Q_F^{N_2}\phi_3)^2\right]S_{34}\right\} \end{aligned} \tag{13}$$

Hard pn Breakup

We now consider the $\gamma +^3 He \rightarrow p + n + p_s$ *reaction.*
Using that,

- *for large angle pn scattering* $\phi_3 \approx \phi_4$,
- *in the QI model it is found that* $Q_F^{pn} = \frac{1}{3}$,
- *and that for* ^3He *in its ground state,* $S_{12}^{pn} \approx S_{34}^{pn} \approx \frac{S_0^{pn}}{2}$,

we obtain that,

$$| \bar{M} |^2 = \frac{(eQ_{F,pn})^2 (2\pi)^6}{s'_{NN}} 16\pi s_{NN}(s_{NN} - 4m_N^2) \frac{d\sigma^{pn\rightarrow pn}(s_{NN}, t_N)}{dt_N} \frac{S_0^{pn}}{2},$$

and for the differential cross section,

$$\frac{d\sigma^{\gamma^3 He\rightarrow (pn)p}}{dt \frac{d^3 p_s}{E_s}} = \alpha Q_{F,pn}^2 16\pi^4 \frac{S_0^{pn}(\alpha = \frac{1}{2}, \vec{p}_s)}{2} \frac{s_{NN}(s_{NN} - 4m^2)}{(s_{NN} - p_{NN}^2)^2 (s - M_{^3He}^2)} \frac{d\sigma^{pn\rightarrow pn}(s_{NN}, t_N)}{dt_N},$$

Hard pp Breakup

To calculate $\gamma +^3 He \to p + p + n_s$ *we use that*

- *for pp scattering,* $\phi_4 \sim -\phi_3$ *for large* $\theta_{c.m}$*(Ramsey,Sivers,1992),*
- *from the exclusion principle and* S *state dominance of the nuclear wave function,* $S_{12}^{pp} << S_{34}^{pp}$*,*
- *and obtaining from QI diagrams that* $Q_F^{pp} = \frac{5}{3}$*.*

Then, for hard pp breakup in ^3He *photodisintegration,*

$$|\bar{\mathcal{M}}|^2 = \frac{(e^2 2(2\pi)^6}{2s'_{NN}} \frac{1}{2} \left\{ 2Q_F^{pp} (|\phi_3| - |\phi_4|)^2 S_{34} \right\},$$

and for the corresponding differential cross section,

$$\frac{d\sigma^{\gamma^3 He \to (pp)n}}{dt \frac{d^3 p_s}{E_s}} = \alpha Q_{F,pp}^2 16\pi^4 S_{34}^{pp} (\alpha = \frac{1}{2}, \vec{p}_s) \frac{2\beta^2}{1+2C^2} \frac{s_{NN}(s_{NN} - 4m_N^2)}{(s_{NN} - p_{NN}^2)^2 (s - M_{^3\mathrm{He}}^2)} \times \frac{d\sigma^{pp \to pp}(s_{NN}, t_N)}{dt},$$

where,

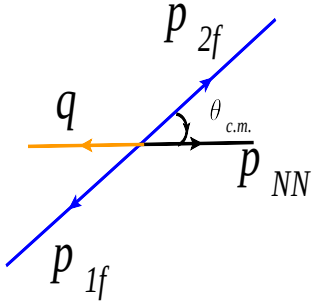
$$C^2 = \frac{\phi_3^2}{\phi_1^2} \approx \frac{\phi_4^2}{\phi_1^2}, \quad and, \quad \beta = \frac{|\phi_3| - |\phi_4|}{|\phi_1|},$$

γNN and $N + N$ Kinematics I

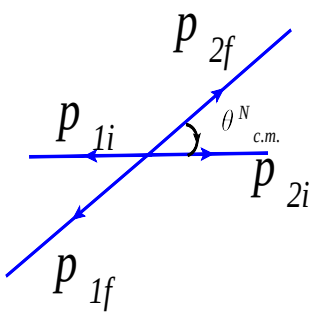
Calculating $\frac{d\sigma^{\gamma^3 He \to (pp)n}}{dt\,\frac{d^3 p_s}{E_s}}(s_{NN}, \theta_{c.m.})$ *within HRM, using* $\frac{d\sigma^{pp\to pp}(s_{NN}, \theta^N_{c.m.})}{dt}$ *requires using the right correspondence between* $\theta_{c.m}$ *and* $\theta^N_{c.m.}$.

In the center of mass reference frame for,

$\gamma + NN \to N + N$



$N + N \to N + N$



$$t = (p_{NN} - p_{2f})^2$$

$$\begin{aligned} t^N &= (p_{2i} - p_{2f})^2 \\ &\approx (\frac{p_{NN}}{2} - p_{2f})^2 \end{aligned}$$

γNN and $N + N$ Kinematics II

which leads to,

$$t^N = \frac{t}{2} + \frac{m_N^2}{2} - \frac{M_{NN}^2}{4}$$

that together with

$$t^N = -\frac{s_{NN} - 4m_N^2}{2}(1 - cos(\theta_{c.m.}^N))$$

and

$$t = -\frac{(s_{NN} - M_{NN}^2)}{2}(1 - \sqrt{1 - 4\frac{m_N^2}{s_{NN}}}cos(\theta_{\text{c.m.}})) + m_N^2$$

is used to obtain that,

$$cos(\theta_{\text{c.m.}}^N) = 1 - \frac{(s_{NN} - M_{NN}^2)}{2(s_{NN} - 4m_N^2)}\frac{(\sqrt{s_{NN}} - \sqrt{s_{NN} - 4m_N^2}cos(\theta_{\text{c.m.}}))}{\sqrt{s_{NN}}} + \frac{4m_N^2 - M_{NN}^2}{2(s_{NN} - 4m_N^2)}.$$

Then for instance, to calculate a HRM ^3He *photodisintegration cross section for* $\theta_{c.m.} = 90^o$*, from the above equation we'll need to have a numerical input for the corresponding* NN *elastic scattering cross section at* $\theta_{c.m.}^N = 60^o$

pn and pp breakup at $\theta_{c.m} = 90^o$

- $$\frac{d\sigma^{\gamma\,^3He\to(pn)p}}{dt\,\frac{d^3p_s}{E_s}} = \alpha Q^2_{F,pn} 16\pi^4 \frac{S_0^{pn}(\alpha=\frac{1}{2},\vec{p}_s)}{2} \frac{s_{NN}(s_{NN}-4m^2)}{(s_{NN}-p^2_{NN})^2(s-M^2_{^3He})} \frac{d\sigma^{pn\to pn}(s_{NN},t_N)}{dt_N},$$

- $$\frac{d\sigma^{\gamma\,^3He\to(pp)n}}{dt\,\frac{d^3p_s}{E_s}} = \alpha Q^2_{F,pp} 16\pi^4 S_{34}^{pp}(\alpha=\frac{1}{2},\vec{p}_s)\frac{2\beta^2}{1+2C^2} \frac{s_{NN}(s_{NN}-4m_N^2)}{(s_{NN}-p^2_{NN})^2(s-M^2_{^3\mathrm{He}})} \times \frac{d\sigma^{pp\to pp}(s_{NN},t_N)}{dt},$$

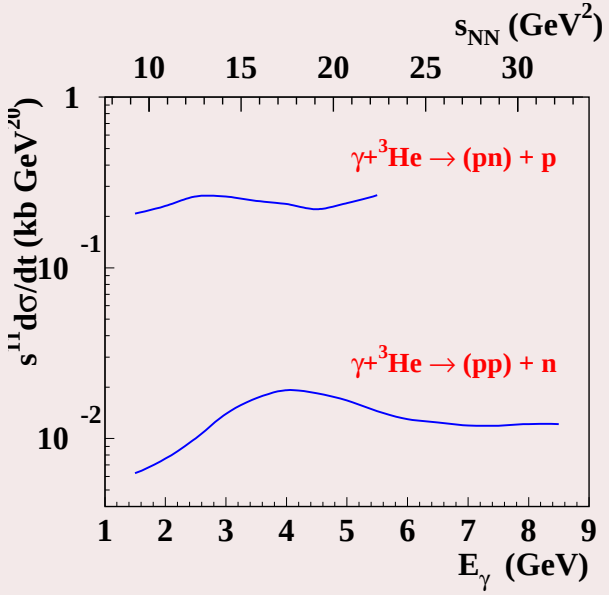
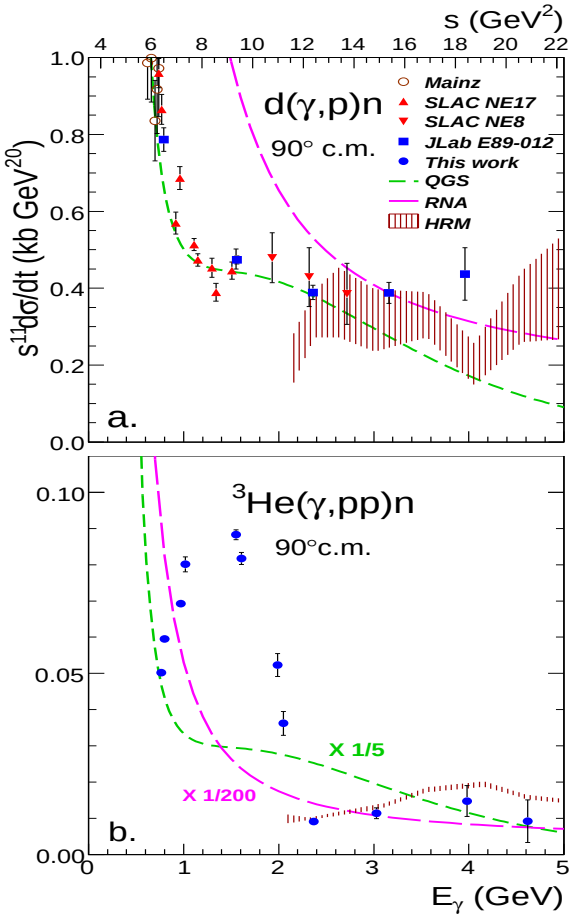
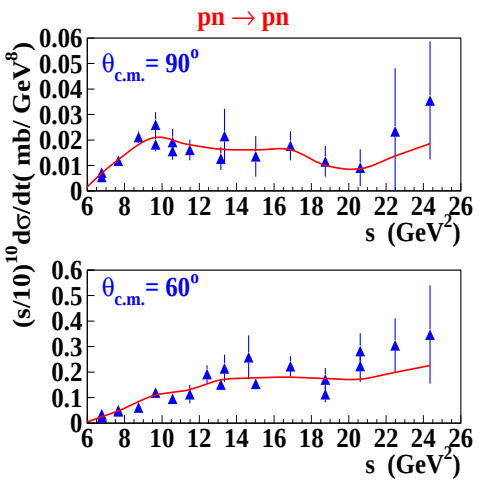


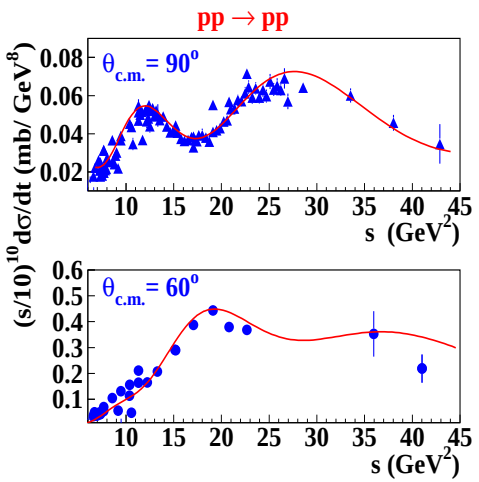
Figure: Energy dependence of s^{11} weighted differential cross sections at 90° c.m. angle scattering in "γ-NN" system. In these calculations one integrated over the spectator nucleon momenta in the range of 0-100 MeV/c.

M. M. Sargsian and C. Granados, Phys. Rev. C 80, 014612 (2009)

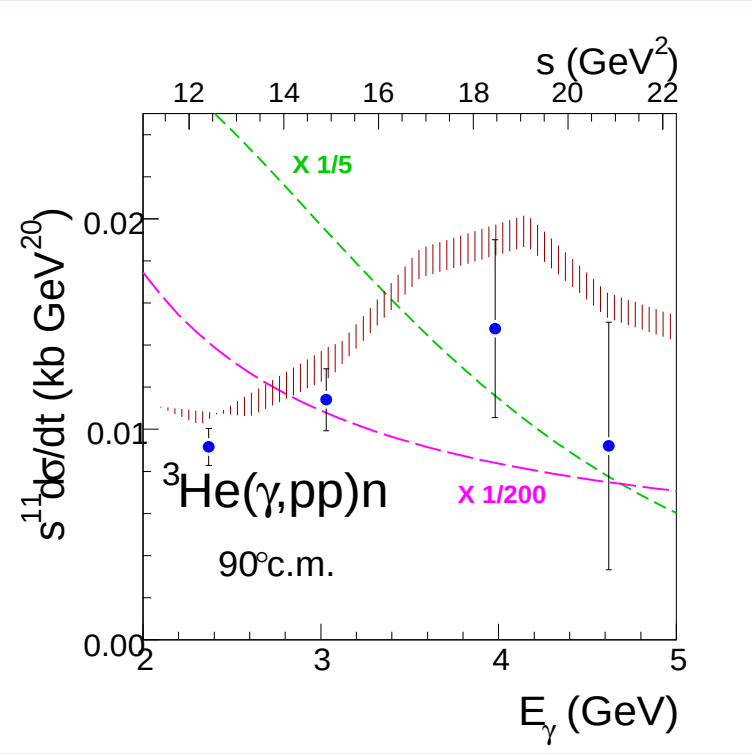
pn and *pp* breakup at $\theta_{c.m} = 90^o$



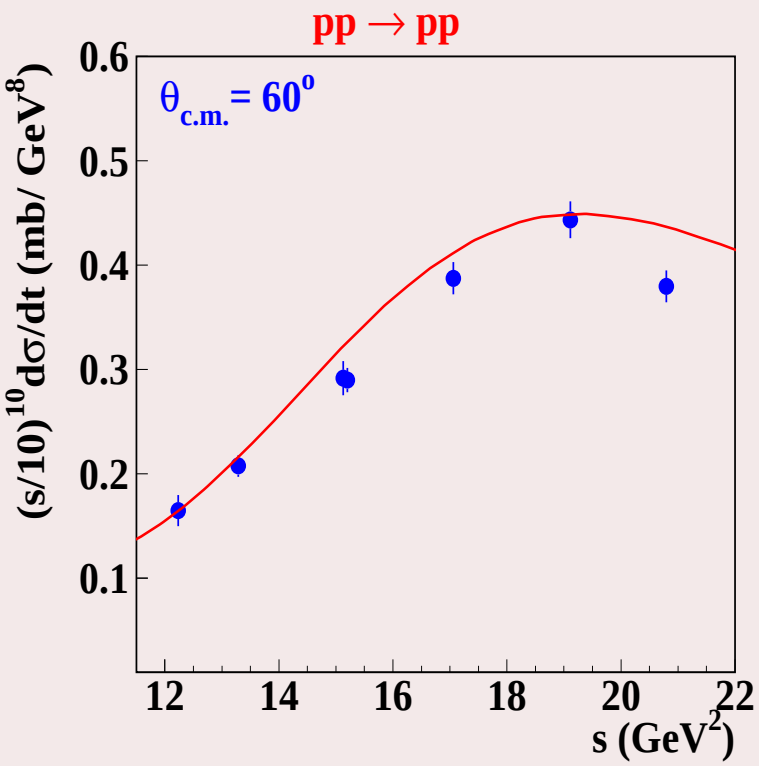




I. Pomerantz et al. [JLab Hall A Collaboration], Phys. Lett. B 684, 106 (2010) [arXiv:0908.2968 [nucl-ex]]



I. Pomerantz(2010)



- *HRM calculated* $^3He(\gamma pp)n$ *energy distribution in agreement with experimental data.*
- *Note*

$$s\sigma^{^3He(\gamma pp)n}(s) \sim \sigma^{pp}(s)$$

Polarization Transfer

For a circularly polarize photon C'_z measures the asymmetry of the hard breakup reaction with respect to the helicity of the outgoing proton.

$$C_{z'} = \frac{\sum_{\lambda_{2f},\lambda_s,\lambda_a} \left\{ |\langle +, \lambda_{2f}, \lambda_s \mid M \mid +, \lambda_A \rangle|^2 - |\langle -, \lambda_{2f}, \lambda_s \mid M \mid +, \lambda_A \rangle|^2 \right\}}{\sum_{\lambda_{1f}\lambda_{2f},\lambda_s,\lambda_a} |\langle \lambda_{1f}, \lambda_{2f}, \lambda_s \mid M \mid +, \lambda_A \rangle|^2}. \tag{14}$$

Then

$$C_{z'} = \frac{(|\phi_1|^2 - |\phi_2|^2)S^{++} + (|\phi_3|^2 - |\phi_4|^2)S^{+-}}{2|\phi_5|^2 S^{+} + (|\phi_1|^2 + |\phi_2|^2)S^{++} + (|\phi_3|^2 + |\phi_4|^2)S^{+-}}, \tag{15}$$

with,

$$S^{\pm,\pm}(t_1, t_2, \alpha, \vec{p}_s) = \sum_{\lambda_A=-\frac{1}{2}}^{\frac{1}{2}} \sum_{\lambda_3=-\frac{1}{2}}^{\frac{1}{2}} \left| \int \Psi^{\lambda_A}_{{}^3\mathrm{He,NR}}(\vec{p}_1, \lambda_1 = \pm\frac{1}{2}, t_1; \vec{p}_2, \lambda_2 = \pm\frac{1}{2}, t_2; \vec{p}_s, \lambda_3) m_N \frac{d^2 p_{2,\perp}}{(2\pi)^2} \right|^2 \tag{16}$$

and $S^+ = S^{++} + S^{+-}$.

Polarization Transfer

$$C_{z'} = \frac{(|\phi_1|^2 - |\phi_2|^2)S^{++} + (|\phi_3|^2 - |\phi_4|^2)S^{+-}}{2|\phi_5|^2 S^{+} + (|\phi_1|^2 + |\phi_2|^2)S^{++} + (|\phi_3|^2 + |\phi_4|^2)S^{+-}}, \tag{17}$$

Through previous assumptions for pp *and* pn *breakup,*

$$C_{z'}^{pp} \approx \frac{|\phi_3|^2 - |\phi_4|^2}{|\phi_3|^2 + |\phi_4|^2} \sim 0, \tag{18}$$

which results from the suppression of S^{++} *and from* $\phi_3 \sim \phi_4$*. And,*

$$C_{z'}^{pn} \approx \frac{|\phi_1|^2 + |\phi_3|^2 - |\phi_4|^2}{|\phi_1|^2 + |\phi_3|^2 + |\phi_4|^2} \sim \frac{2}{3}, \tag{19}$$

Summary I

Large angle high energy breakup of a NN system in ^3He has been studied in pp and pn breakup channels within the framework of the QCD hard rescattering model,HRM.
HRM predicted energy dependencies, in accordance with counting rules, agree with recent experimental observations which supports the explicit role of QCD degrees of freedom in the reaction.

- *Calculated HRM* $\gamma +^3 He \rightarrow p + p + n_s$ *differential cross section at* $\theta^{NN}_{c.m.} = 90^o$ *agrees well with experimental data without the introduction of adjustable parameters. (Note suppression from* $|\phi_3| - |\phi_4|$ *factor with respect to* pn *breakup).*
- *The* s^{-11} *behavior of the observed cross section for* $\gamma +^3 He \rightarrow p + p + n_s$ *indicates the dominance of the two body process in the reaction.*

Additional observables are calculated from the HRM applied to $\gamma +^3 He \rightarrow N + N + N_s$*, which predicts*

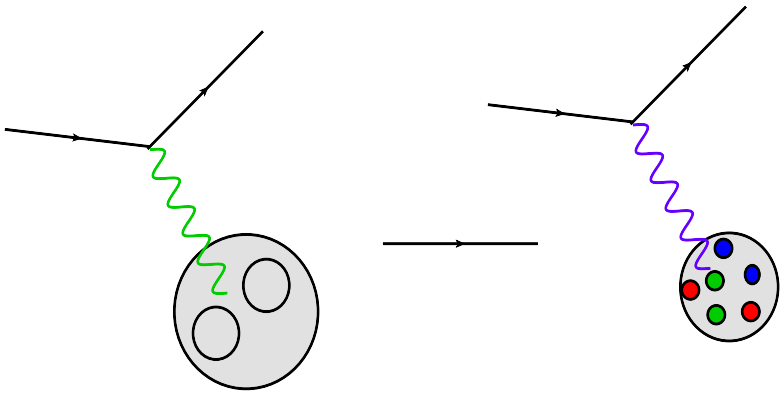
- *a broader spectator's momentum distribution for* pp *breakup in relations to* pn *breakup.*
- $C^{pp}_{z'}$ *suppression in relation to* $C^{pn}_{z'}$

Experimental confirmation of these predictions will further reinforce the HRM picture of breakup processes

The encouraging experimental results motivate the study of more processes that can be described through the hard rescattering model.

Δ-Isobar Production in Deuteron Breakup
Motivation I

The study of the deuteron photodisintegration into $\Delta\Delta$-isobars channels was proposed as a venue for investigating the evolution of a nucleon-nucleon system into a six quark system.



The onset of a six quark picture of the deuteron could then be marked by a large increase of the $\gamma d \to \Delta\Delta$ cross section. This prediction assumes that such cross section is small for a nucleon dominated deuteron wave function because of its suppressed $\Delta\Delta$ components.

In contrast, for a six quark deuteron, NN and $\Delta\Delta$ components contribute with comparable strength to the deuteron wave function (roughly 10% and 8% respectively) while more than 80% is contributed by CC (hiden color) components for which unlike N or Δ, C has a color charge (Brodsky,Ji,Lepage,1983).

$$\psi_d = \sqrt{\frac{1}{9}}\psi_{NN} + \sqrt{\frac{4}{45}}\psi_{\Delta\Delta} + \sqrt{\frac{4}{5}}\psi_{cc}$$

Δ-Isobar Production in Deuteron Breakup
Motivation II

High energy $\gamma d \to \Delta\Delta$ *with* $\Delta\Delta$ *emerging at large transverse momentum is thought to probe the onset of hidden color components in the deuteron. Assuming that the high* p_T $\Delta\Delta$ *system was created in the initial state of the interaction, in the asymptotic limit we have that*

$$\frac{d\sigma^{\gamma d \to \Delta\Delta}}{dt} \sim \frac{d\sigma^{\gamma d \to pn}}{dt}.$$

Under this same assumption we also have that,

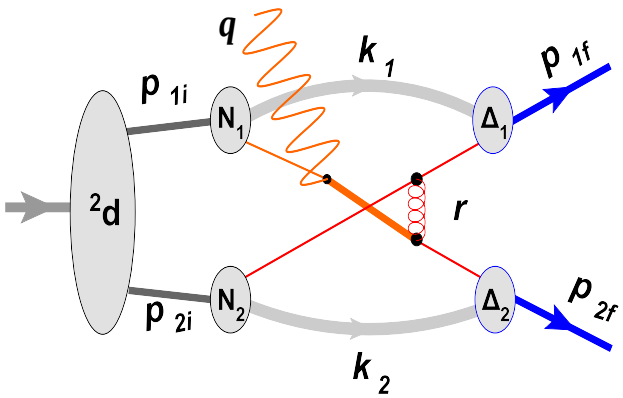
$$\frac{d\sigma^{\gamma d \to \Delta^{++}\Delta^{-}}}{dt} = \frac{d\sigma^{\gamma d \to \Delta^{+}\Delta^{0}}}{dt},$$

since both $\Delta\Delta$ *channels in general contribute with the same strength to the spin-isospin wave function of the deuteron. In the QCD hard rescattering model (HRM), the high* p_T $\Delta\Delta$ *system is created in a final state interaction mainly through a* pn *rescattering reaction. As it'll be shown in what follows under this scenario, the hard rescattering model predicts a dominance of* $\frac{d\sigma^{\gamma d \to \Delta^{++}\Delta^{-}}}{dt}$ *over* $\frac{d\sigma^{\gamma d \to \Delta^{+}\Delta^{0}}}{dt}$ *which contrasts the picture described above.*

Δ-Isobar Production in Deuteron Breakup
Motivation III

For comparison, we also estimate within the HRM the strength of $\Delta\Delta$ *channels in deuteron breakup relative to that of* $\gamma d \to pn$

Δ-Isobar Production in Deuteron Breakup
$\gamma d \to \Delta\Delta$ in the HRM



As it was the case for $\gamma d \to pn$, *through the HRM for* $\gamma d \to \Delta\Delta$ *we find that,*

$$\frac{d\sigma^{\gamma d \to \Delta\Delta}(s, \theta_{c.m.})}{dt} = \frac{\alpha Q^2_{F,\Delta\Delta} 8\pi^4}{s'} \frac{d\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt} \bar{S}_{0,NR},$$

But $\frac{\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt}$ *has not been well measured in the required kinematic regime.*

C. G. Granados and M. M. Sargsian,[arXiv: hep-ph 1007.4705] (2010)

Δ-Isobar Production in Deuteron Breakup $\gamma d \to \Delta\Delta$ in the HRM

As it was the case for $\gamma d \to pn$, through the HRM for $\gamma d \to \Delta\Delta$ we find that,

$$\frac{d\sigma^{\gamma d \to \Delta\Delta}(s, \theta_{c.m.})}{dt} = \frac{\alpha Q^2_{F,\Delta\Delta} 8\pi^4}{s'} \frac{d\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt} \bar{S}_{0,NR},$$

But $\frac{\sigma^{pn \to \Delta\Delta}(s, \theta^N_{c.m.})}{dt}$ has not been well measured in the required kinematic regime. Then model dependent $pn \to \Delta\Delta$ amplitudes ϕ are input in

$$\begin{aligned}
|\bar{\mathcal{M}}|^2{}_{\gamma d \to \Delta\Delta} &= \frac{1}{2}\frac{1}{3}\frac{e^2}{2s'} \\
&\times \Big[S_{12} \Big\{ |(\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_1|^2 + |(\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_6|^2 + |(\hat{Q}^{N_1} + \hat{Q}^{N_2})\phi_7|^2 \Big\} \\
&+ S_{34} \Big\{ |\hat{Q}^{N_1}\phi_3 + \hat{Q}^{N_2}\phi_4|^2 + |\hat{Q}^{N_1}\phi_4 + \hat{Q}^{N_2}\phi_3|^2 \\
&+ |\hat{Q}^{N_1}\phi_8 + \hat{Q}^{N_2}\phi_9|^2 + |\hat{Q}^{N_1}\phi_9 + \hat{Q}^{N_2}\phi_8|^2 \Big\} \Big],
\end{aligned}$$

C. G. Granados and M. M. Sargsian,[arXiv: hep-ph 1007.4705] (2010)

$\gamma d \to \Delta\Delta$ in the HRM

$pn \to \Delta\Delta$ scattering amplitudes

Helicity conserving $pn \to \Delta\Delta$ *amplitudes*

$$
\begin{aligned}
\left\langle +\frac{1}{2}, +\frac{1}{2} \,|T|\, +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_1 \\
\left\langle +\frac{1}{2}, -\frac{1}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_3 \\
\left\langle -\frac{1}{2}, +\frac{1}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_4 \\
\left\langle +\frac{3}{2}, -\frac{1}{2} \,|T|\, +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_6 \\
\left\langle -\frac{1}{2}, +\frac{3}{2} \,|T|\, +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_7 \\
\left\langle +\frac{3}{2}, -\frac{3}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_8 \\
\left\langle -\frac{3}{2}, +\frac{3}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_9,
\end{aligned}
$$

$\gamma d \to \Delta\Delta$ in the HRM

$pn \to \Delta\Delta$ scattering amplitudes

Helicity conserving $pn \to \Delta\Delta$ *amplitudes* *Obtained in the quark exchange model*

$pn \to \Delta^+\Delta^0$

$$\begin{aligned}
\phi_1 &= \frac{2}{9} N_{\Delta\Delta}(2f(\theta^N_{c.m.}) - f(\pi - \theta^N_{c.m.})) \\
\phi_3 &= \frac{1}{9} N_{\Delta\Delta}(4f(\theta^N_{c.m.}) + f(\pi - \theta^N_{c.m.})) \\
\phi_4 &= \frac{2}{9} N_{\Delta\Delta}(f(\theta^N_{c.m.})) + f(\pi - \theta^N_{c.m.}) \\
\phi_6 &= \frac{N_{\Delta\Delta}}{3\sqrt{3}}(2f(\theta^N_{c.m.}) - f(\pi - \theta^N_{c.m.})) \\
\phi_7 &= \frac{N_{\Delta\Delta}}{3\sqrt{3}}(2f(\theta^N_{c.m.}) - f(\pi - \theta^N_{c.m.})) \\
\phi_8 &= \frac{2}{9} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_9 &= \frac{1}{3} N_{\Delta\Delta} f(\pi - \theta^N_{c.m.}),
\end{aligned}$$

$pn \to \Delta^{++}\Delta^-$

$$\begin{aligned}
\phi_1 &= -\frac{2}{3} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_3 &= -\frac{2}{3} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_4 &= -\frac{1}{3} N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_6 &= \frac{-N_{\Delta\Delta}}{\sqrt{3}} f(\theta^N_{c.m.}) \\
\phi_7 &= \frac{-N_{\Delta\Delta}}{\sqrt{3}} f(\theta^N_{c.m.}) \\
\phi_8 &= -N_{\Delta\Delta} f(\theta^N_{c.m.}) \\
\phi_9 &= 0,
\end{aligned}$$

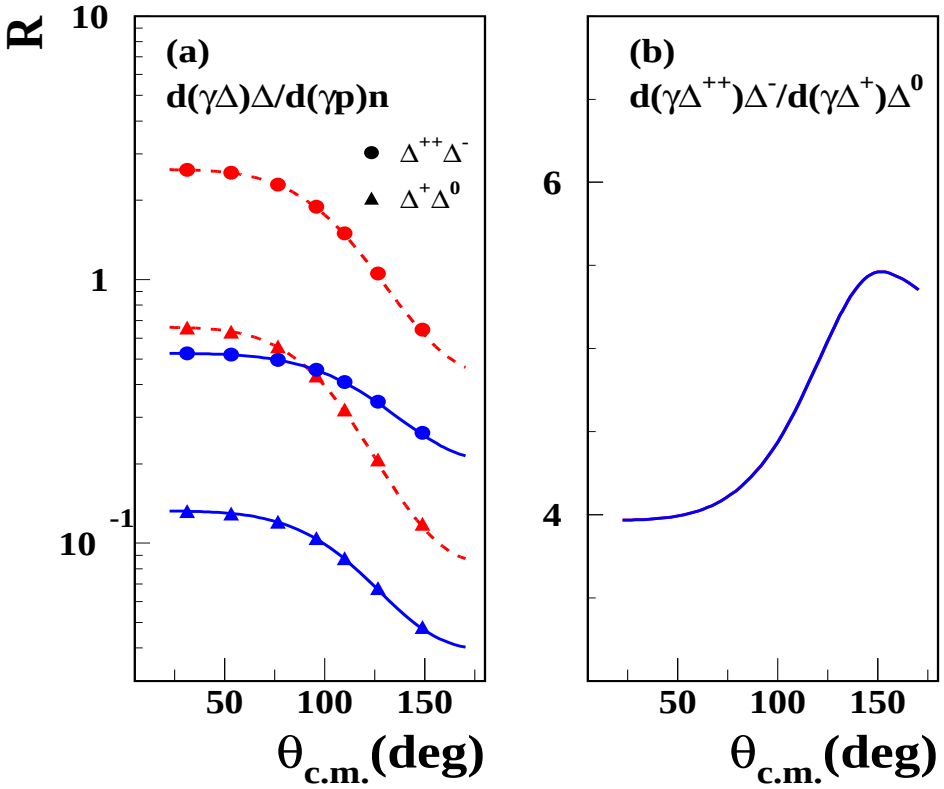
$\gamma d \to \Delta\Delta$ in the HRM

$$
\begin{aligned}
|\bar{\mathcal{M}}|^2_{\gamma d \longrightarrow \Delta^+\Delta^-} \quad = \quad & \frac{1}{6}\frac{e^2}{2s'} Q^2_{F,\Delta\Delta} \left\{ S_{12} \left[|\phi_1|^2 + |\phi_6|^2 + |\phi_7|^2 \right] \right. \\
& + S_{34} \left[\left(\frac{\phi_3^t + \phi_4^t}{2} + 2\phi_4^u - \phi_3^u \right)^2 + \left(\frac{\phi_4^t + \phi_3^t}{2} + 2\phi_3^u - \phi_4^u \right)^2 \right. \\
& \left. \left. + \left(\frac{\phi_8^t + \phi_9^t}{2} + 2\phi_9^u - \phi_8^u \right)^2 + \left(\frac{\phi_9^t + \phi_8^t}{2} + 2\phi_8^u - \phi_9^u \right)^2 \right] \right\},
\end{aligned}
$$

$$
\begin{aligned}
|\bar{\mathcal{M}}|^2_{\gamma d \longrightarrow \Delta^{++}\Delta^-} = & \frac{1}{6}\frac{e^2}{2s'} Q^2_{F,\Delta\Delta} \left\{ S_{12} \left(|\phi_1|^2 + |\phi_6|^2 + |\phi_7|^2 \right) \right. \\
& \left. + S_{34} \left[(2\phi_3 - \phi_4)^2 + (2\phi_4 - \phi_3)^2 + 5|\phi_8|^2 \right] \right\}.
\end{aligned}
$$

$\gamma d \to \Delta\Delta$ in the HRM

Angular Distributions



C. G. Granados and M. M. Sargsian, Phys. Rev. C 83, 054606 (2011).

R in Fig. (a)

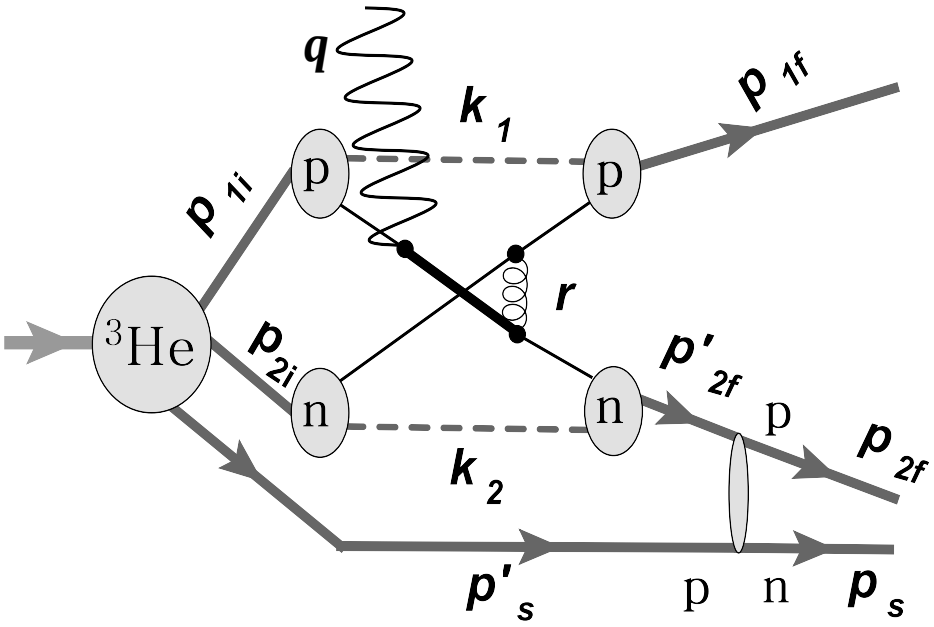
- *is sensitive to model of baryon wavefunction*
- *shows a strong dependence on $\theta_{c.m}$ contrasting R$\sim$ 1 expected from the onset of hidden color components.*

In Fig. (b)

- *R doesn't depend on the choice of wave function*
- *$\Delta^{++}\Delta^{-}$ channel is consistently larger than that of the $\Delta^{+}\Delta^{0}$ channel. If the isobars were produced from a $\Delta\Delta$ component of the initial deuteron wave function, both channels should have the same strength and R=1.*

Thank you

Three-body/two-step reaction



- *Dominant at low to intermediate beam energies ($E_\gamma \sim 200 MeV$).*
- *HRM amplitude does not interfere with two-body/one-step amplitude.*
- *HRM cross section scales like s^{-12} at large energies from second rescattering.*

Introducing the light-cone wave function of ^{3}*He [**?**, **?**, **?**]*

$$\Psi_{^3\mathrm{He}}^{\lambda_A,\lambda_1,\lambda_2,\lambda_s}(\alpha, p_\perp) = \frac{\Gamma_{^3\mathrm{He}}^{\lambda_A}\bar{u}_{\lambda_1}(p_{NN}-p)\bar{u}_{\lambda_2}(p)\bar{u}_{\lambda_s}(p_s)}{M_{NN}^2 - \frac{m_N^2+p_\perp^2}{\alpha(1-\alpha)}} \tag{20}$$

defining quark wave function of the nucleon as

$$\Psi_N^{\lambda,\eta}(p, x, k_\perp) = \frac{u_N^\lambda(p)\Gamma_N\bar{u}_\eta(p-k)\psi_s^\dagger(k)}{m_N^2 - \frac{m_s^2(1-x)+m_q^2x+(k_\perp-xp_\perp)^2}{x(1-x)}} \tag{21}$$

L.L. Frankfurt and M.I. Strikman, Phys. Rep. 76, 214 (1981)

$$\Psi_{^3\mathrm{He}}(\alpha, p_\perp, \alpha_s, p_{s,\perp}) = \sqrt{2}(2\pi)^3 m_N \Psi_{^3\mathrm{He,NR}}(\alpha, p_\perp, \alpha_s, p_{s,\perp}) \quad (22)$$

A. Nogga A. Kievsky, H. Kamada, W. Gloeckle, L. E. Marcucci, S. Rosati and M. Viviani, Phys. Rev. C 67, 034004 (2003).

SU(6)Helicity Amplitudes

For pp scattering:

Farrar et al.,(1979).

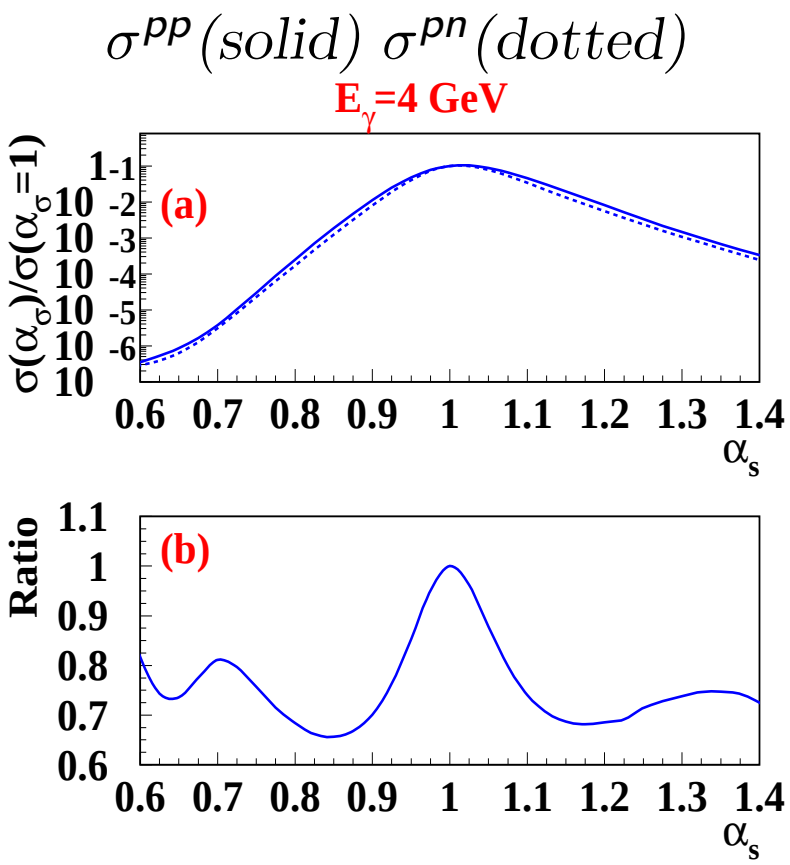
$$
\begin{aligned}
\phi_1(\theta_{CM}) &= 144f(\theta_{CM}) + 144f(\pi - \theta_{CM}) \qquad (23)\\
\phi_2(\theta_{CM}) &= 0\\
\phi_3(\theta_{CM}) &= 56f(\theta_{CM}) + 68f(\pi - \theta_{CM})\\
\phi_4(\theta_{CM}) &= -68f(\theta_{CM}) - 56f(\pi - \theta_{CM})\\
\phi_5(\theta_{CM}) &= 0
\end{aligned}
$$

Uncertainty in $\gamma^3\text{He} \rightarrow NN + N_s$

$$R(p_s) = \frac{4\Psi_{^3\text{He}}^{\lambda_A,\lambda_{1i},\lambda_{2i},\lambda_s}(\alpha_c = \frac{1}{2}, p_\perp, p_s)}{\int \frac{d\alpha_c}{\alpha_c} \frac{\Psi_{^3\text{He}}^{\lambda_A,\lambda_{1i},\lambda_{2i},\lambda_s}(\alpha_c, p_\perp, p_s)}{\sqrt{(1-\alpha_c)\alpha_c}}}. \tag{24}$$

This ratio depends on the kinematics of the spectator nucleon, and for the case of $p_s \leq 100$ *MeV/c*, $R(p_s) \approx 1.1$, *which corresponds to* $\sim$20% *of uncertainty in the cross section of the reaction calculated with the* $\alpha_c = \frac{1}{2}$ *approximation.*

Spectator Momentum Distributions



$$Ratio = \frac{\sigma^{pp}}{\sigma^{pn}}$$

$$\alpha_s \equiv \frac{E_s - p_{s,z}}{M_A/A} = \alpha_A - \alpha_{NN}$$

with

$$s_{NN} = M_{NN}^2 + E_\gamma m_n \alpha_{NN}.$$

- *Asymmetry around* $\alpha = 1$ *due to* s^{-11} *dependence.*
- pp *distribution broader than* pn.
- $\boldsymbol{R}$ *drops around* $\alpha = 1$ *from the suppression of same helicity two proton components of the nuclear wave function at small momenta.*

M. M. Sargsian and C. Granados, Phys. Rev. C 80, 014612 (2009)