Impact of Dynamical Fermions on the Vacuum of Quantum Chromodynamics

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Appendix A

Papers by the author

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[8] K. Maltman, D. Leinweber, P. Moran and A. Sternbeck, *The determination of $\alpha_s(M_z)$ from perturbative analyses of short-distance-sensitive lattice QCD observables revisited*, PoS **LATTICE2008**, 214 (2008) [arXiv:0812.2484 [hep-lat]].

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Appendix B

Dirac γ matrix properties

A collection of useful properties for working with Dirac γ matrices.

B.1 Sakurai representation

B.1.1 γ matrices

- $\{\gamma_\mu, \gamma_\nu\} = 2\,\delta_{\mu\nu}$, where $\mu, \nu = 1, 2, 3, 4$
- $\gamma_\mu^\dagger = \gamma_\mu$
- $\gamma_\mu^2 = 1$
- $\gamma_5 = \gamma_1\gamma_2\gamma_3\gamma_4 = \frac{1}{4!}\epsilon_{\mu\nu\lambda\sigma}\gamma_\mu\gamma_\nu\gamma_\lambda\gamma_\sigma$
- $\{\gamma_5, \gamma_\mu\} = 0$, where $\mu, \nu = 1, 2, 3, 4$
- $\gamma_5^\dagger = \gamma_5$
- $\gamma_5^2 = 1$
- $\gamma_k = \begin{pmatrix} 0 & -i\sigma_k \\ i\sigma_k & 0 \end{pmatrix}$, where $k = 1, 2, 3$
- $\gamma_4 = \begin{pmatrix} I & 0 \\ 0 & -I \end{pmatrix}$
- $\gamma_5 = -\begin{pmatrix} 0 & I \\ I & 0 \end{pmatrix}$
- $\gamma_2^* = \gamma_2$, $\gamma_4^* = \gamma_4$, $\gamma_5^* = \gamma_5$
- $\sigma_{\mu\nu} = \frac{[\gamma_\mu, \gamma_\nu]}{2i} = -i\gamma_\mu\,\gamma_\nu$

B.1.2 Charge conjugation matrix C

- $C\,\gamma_\mu C^{-1} = -\gamma_\mu^T$, or $C^\dagger\,\gamma_\mu C = -\gamma_\mu^T$
- $C \equiv \gamma_4\gamma_2$
- $-C = C^T = C^{-1} = C^\dagger$
- $(C\,\gamma_\mu)^T = C\,\gamma_\mu$
- $(\gamma_\mu\,C)^T = \gamma_\mu\,C$

B.2 Bjorken and Drell representation

B.2.1 γ matrices

- $\{\gamma^\mu, \gamma^\nu\} = 2\,g^{\mu\nu}$, where $\mu, \nu = 0, 1, 2, 3$
- $g_{\mu\nu} = g^{\mu\nu}$, $g_{00} = 1$, $g_{kk} = -1$, $g_{\mu\nu} = 0$ if $\mu \neq \nu$
- $\gamma^{0\dagger} = \gamma^0$
- $\gamma^{k\dagger} = -\gamma^k$
- $a \cdot b = a^\mu b_\mu = a_0 b_0 - \vec{a} \cdot \vec{b}$, where $a^\mu = (a_0, \vec{a})$
- $a_\mu = g_{\mu\nu} a^\nu = (a_0, -\vec{a})$
- $\gamma_0 = \begin{pmatrix} I & 0 \\ 0 & -I \end{pmatrix}$
- $\gamma_k = \begin{pmatrix} 0 & \sigma_k \\ -\sigma_k & 0 \end{pmatrix}$, where $k = 1, 2, 3$
- $\gamma_5 = -\frac{i}{24}\,\epsilon_{\alpha\beta\gamma\delta}\,\gamma^\alpha\,\gamma^\beta\,\gamma^\gamma\,\gamma^\delta = i\,\gamma^0\,\gamma^1\,\gamma^2\,\gamma^3$
- $\gamma_5^2 = I$
- $\gamma_5^\dagger = \gamma_5$
- $\{\gamma_5, \gamma_\mu\} = 0$
- $\gamma_0\,\gamma_\mu^\dagger\,\gamma_0 = \gamma_\mu$

B.2.2 Charge conjugation matrix C

- $C\,\gamma_\mu C^{-1} = -\gamma_\mu^T$, or $C^\dagger\,\gamma_\mu C = -\gamma_\mu^T$
- $C \equiv -i\gamma_0\gamma_2$
- $-C = C^T = C^{-1} = C^\dagger$
- $C\gamma_5 C^\dagger = \gamma_5$
- $[C, \gamma_5] = \gamma_5$
- $(C\,\gamma_5)^T = -C\,\gamma_5$
- $(C\,\gamma_\mu)^T = C\,\gamma_\mu$
- $(\gamma_\mu\,C)^T = \gamma_\mu\,C$
- $\gamma_0\,C^\dagger\,\gamma_0 = C$
- $\gamma_0\,(C\,\gamma_\mu)^\dagger\,\gamma_0 = \gamma_\mu\,C$

Appendix C

Baryon interpolating fields

From simple considerations we derive the general form for a baryon interpolating field.

C.1 Baryon interpolating fields

Let us denote the annihilation operator for a quark at the space-time point x as,

$$\psi_\alpha^{af}(x)\,, \tag{C.1}$$

where the indices a, b, and c will label colour; α, β, and γ will label Dirac spinor indices; and f, g, and h represent the flavour of the quark, up, down, etc. Baryons are composed of three quarks and therefore we require the baryon interpolating field $\chi(x)$ to contain three fermion operators,

$$\chi(x) = \ldots\ \psi_\alpha^{af}(x)\,\psi_\beta^{bg}(x)\,\psi_\gamma^{ch}(x)\,. \tag{C.2}$$

where the ... represent terms in either colour, Dirac, or flavour space (or some combination thereof). One could include a term X^{fgh} that specifies the flavour combination of quarks, however in practice it is easier to just specify the required flavours when constructing the required baryon, e.g.

$$u_\alpha^a(x) \equiv \psi_\alpha^{au}(x)\,. \tag{C.3}$$

We therefore suppress the flavour indices from this point on and illustrate the flavour in place of ψ.

C.1.1 Colour indices

We require the baryon to be a colour-singlet state, and this implies the use of a Levi-Cevita tensor, ϵ^{abc}. Including this factor will also ensure local gauge invariance of the interpolating field. To see this, consider a local gauge transformation (suppressing Dirac indices for the moment),

$$\psi^a(x) \to G^{ab}(x)\psi^b(x)\,, \tag{C.4}$$

where $G(x) \in SU(3)$ for QCD. Under such a transformation we have,

$$\chi(x) \to \ldots\ G^{aa'}(x)\,G^{bb'}(x)\,G^{cc'}(x)\,\psi^{a'}(x)\,\psi^{b'}(x)\,\psi^{c'}(x)\,. \tag{C.5}$$

Now,

$$\epsilon^{abc}\, G^{aa'}(x)\, G^{bb'}(x)\, G^{cc'}(x) = \epsilon^{a'b'c'} \det(G(x)) \\ = \epsilon^{a'b'c'}\,, \tag{C.6}$$

where $\det(G(x)) = 1$ since $G(x) \in SU(3)$. Therefore,

$$\chi(x) = \ldots\, \epsilon^{abc}\, \psi^a_\alpha(x)\, \psi^b_\beta(x)\, \psi^c_\gamma(x)\,, \tag{C.7}$$

is a colour singlet state and invariant under both local and global gauge transformations as required.

C.1.2 Dirac indices

Our interpolator will be used to create a spin-1/2 baryon described by a Dirac spinor. Therefore, two of the Dirac indices must be contracted. We can contract the Dirac indices using arbitrary combinations of Dirac γ matrices, which we denote $\Gamma_{\{1,2\}}$, giving

$$\chi_\gamma(x) = \epsilon^{abc}\, \psi^a_\alpha(x)\, (\Gamma_1)_{\alpha\beta}\, \psi^b_\beta(x)\, (\Gamma_2)_{\gamma\gamma'}\, \psi^c_{\gamma'}(x)\,. \tag{C.8}$$

Note that the presence of the three fermion fields implies that our baryon interpolating field must be a Dirac spinor, $\chi_\gamma(x)$. One could also consider an alternate contraction of the Dirac indices,

$$\chi_\alpha(x) = \epsilon^{abc}\, \psi^a_{\alpha'}(x)\, (\Gamma_1)_{\alpha\alpha'}\, \psi^b_\beta(x)\, (\Gamma_2)_{\beta\gamma}\, \psi^c_\gamma(x)\,, \tag{C.9}$$

however a cyclic permutation of the colour indices results in the same interpolating field as Eq. (C.8).

C.1.3 Lorentz covariance

We also require that our baryon interpolating field transforms as a Dirac spinor under Lorentz transformations,

$$\psi^a_\alpha(x) \to S(\Lambda)_{\alpha\alpha'} \psi^a_{\alpha'}(\Lambda\, x)\,, \tag{C.10}$$

where $S(\Lambda) = e^{-i\omega^{\mu\nu}\sigma_{\mu\nu}}$. Under such a transformation $\chi_\gamma(x)$ transforms as (suppressing the space-time coordinate for clarity),

$$\chi_\gamma \to \epsilon^{abc} \left[\psi^a_{\alpha'}\, S(\Lambda)_{\alpha\alpha'}\, (\Gamma_1)_{\alpha\beta}\, S(\Lambda)_{\beta\beta'}\, \psi^b_{\beta'}\right] (\Gamma_2)_{\gamma\gamma'}\, S(\Lambda)_{\gamma'\gamma''}\, \psi^c_{\gamma''}\,, \tag{C.11}$$

where the factor in square brackets is a Dirac scalar. We can therefore move $S(\Lambda)_{\gamma'\gamma''}$ to the front if it commutes with Γ_2. Since $\{\gamma_5, \gamma_\mu\} = 0$ (see App. B), it is not too hard to show that $[\gamma_5, S] = 0$. In practice, it is sufficient to work with $\Gamma_2 = \gamma_5$ or I, with the choice depending on the desired parity transformation of the interpolator. Other forms may then be derived via Fiertz transforms.

For the factor in square brackets we have the product,

$$S(\Lambda)_{\alpha\alpha'}\, (\Gamma_1)_{\alpha\beta}\, S(\Lambda)_{\beta\beta'} = (S^T \Gamma_1 S)_{\alpha'\beta'}\,. \tag{C.12}$$

We already know that $S^{-1}\,\gamma_5\,S = \gamma_5$, and we can also make use of the identity, $S^{-1}\,\gamma_\mu\,S = \Lambda_\mu^{\ \nu}\,\gamma_\nu$. However, it is S^T that appears in Eq. (C.12), not S^{-1}. Fortunately we also have $S^T\,C = C\,S^{-1}$ where C is the charge conjugation matrix. To see this, consider

$$\begin{aligned} S^T\,C &= e^{-i\omega^{\mu\nu}\sigma^T_{\mu\nu}}\,C \\ &= C\,C^T\left(1 - i\omega^{\mu\nu}\sigma^T_{\mu\nu} + (-i\omega^{\mu\nu}\sigma^T_{\mu\nu})^2 + \ldots\right)C \\ &= C\,e^{-i\omega^{\mu\nu}C\sigma^T_{\mu\nu}C^T} \\ &= C\,e^{i\omega^{\mu\nu}\sigma_{\mu\nu}} \\ S^T\,C &= C\,S^{-1}\,, \end{aligned} \tag{C.13}$$

where we have used two useful properties of γ matrices,

$$C\,C^T = I \quad \text{and} \quad C\,\sigma^T_{\mu\nu}\,C^T = -\sigma_{\mu\nu}\,. \tag{C.14}$$

Therefore, in order to ensure Lorentz covariance of our baryon interpolating field we require that $\Gamma_1 \to C\,\Gamma_1$, where Γ_1 can now be γ_5 or I, with the selection dependent on the baryon required. We note that for spin-3/2 baryons, the Rarita-Schwinger spin-vector requires a Lorentz index. In this case Γ_1 can be a combination of γ_μ and γ_5 or I. The final version for the interpolating field (suppressing Dirac and flavour indices) is therefore,

$$\chi_{(\mu)}(x) = \epsilon^{abc}\left(\psi^a(x)\,C\,\Gamma_{(\mu)}\,\psi^b(x)\right)\Gamma\,\psi^c\,, \tag{C.15}$$

where the Lorentz index μ appears only for decuplet baryons.

Appendix D

The $U + U^*$ trick

The action for Lattice QCD is invariant under the change $U \rightarrow U^*$, and the link variables $\{U\}$ and $\{U^*\}$ are therefore gauge field configurations of equal weight. Accounting for both sets of configurations when creating correlation functions will therefore reduce statistical noise. In theory, one could create an entirely new propagator from each $\{U^*\}$ gauge field, however this would take much compute time. Fortunately we have the $U + U^*$ trick [223].

D.1 Derivation

Assume we have some γ matrix $\widetilde{C}$ with the property,

$$\widetilde{C}\,\gamma_\mu\,\widetilde{C}^{-1} = \gamma_\mu^* \,. \tag{D.1}$$

For a general clover-like action it is not too difficult to derive the fermion-matrix property,

$$M(\{U^*\}) = (\widetilde{C}\,M(\{U\})\,\widetilde{C}^{-1})^* \,. \tag{D.2}$$

To see this, first consider the γ-matrix, $\sigma_{\mu\nu} = c\,i\,[\gamma_\mu, \gamma_\nu]$ where c is some real constant that depends on the γ-matrix representation used. Multiplying by $\widetilde{C}$ and $\widetilde{C}^{-1}$ gives,

$$\begin{aligned}\widetilde{C}\,\sigma_{\mu\nu}\,\widetilde{C}^{-1} &= c\,i\left(\widetilde{C}\,\gamma_\mu\,\gamma_\nu\,\widetilde{C}^{-1} - \widetilde{C}\,\gamma_\nu\,\gamma_\mu\,\widetilde{C}^{-1}\right) \\ &= c\,i\left(\gamma_\mu^*\,\gamma_\nu^* - \gamma_\nu^*\,\gamma_\mu^*\right)\,,\end{aligned} \tag{D.3}$$

where we have inserted the identity $\widetilde{C}\,\widetilde{C}^{-1} = I$ and used the property of Eq. (D.1). From Eq. (D.3) it is clear that $(\widetilde{C}\,\sigma_{\mu\nu}\,\widetilde{C}^{-1})^* = -\sigma_{\mu\nu}$. Recalling the general form for a clover-improved fermion matrix,

$$\begin{aligned}M^W_{xy}[U]\,a = \delta_{xy} \,-\, \kappa \sum_\mu &\Big[(r - \gamma_\mu)\,U_{x,\mu}\,\delta_{x,y-\mu} + (r + \gamma_\mu)\,U^\dagger_{x-\mu,\mu}\,\delta_{x,y+\mu}\Big] \\ &- \frac{i\,g\,a\,C_{SW}\,r}{4}\,\sigma_{\mu\nu}\,F_{\mu\nu}\,\delta_{xy}\,,\end{aligned} \tag{D.4}$$

it is then easy to see that Eq. (D.2) is true in general. Since the quark propagator is the inverse of the fermion matrix, it follows that,

$$S(x,0;\{U^*\}) = (\widetilde{C}\,S(x,0;\{U\})\,\widetilde{C}^{-1})^* \,. \tag{D.5}$$

Similarly, for an SST propagator,

$$\begin{aligned}\widehat{S}(x_2, 0; t_1, \vec{q}, \mu; \{U\}) = Q_f \sum_{x_1} e^{+i\vec{q}\cdot\vec{x}_1}\, \kappa \Big(\\ S(x_2, x_1 + \hat{\mu}; \{U\})\, (1 + \gamma^\mu)\, U^{\mu\dagger}(x_1)\, S(x_1, 0; \{U\}) \\ - S(x_2, x_1; \{U\})\, (1 - \gamma^\mu)\, U^\mu(x_1)\, S(x_1 + \hat{\mu}, 0; \{U\}) \Big) ,\end{aligned} \tag{D.6}$$

using Eq. (D.5) we see that,

$$\begin{aligned}\left(\widetilde{C}\, \widehat{S}(x_2, 0; t_1, -\vec{q}, \mu; \{U\})\, \widetilde{C}^{-1}\right)^* = Q_f \sum_{x_1} e^{+i\vec{q}\cdot\vec{x}_1}\, \kappa \Big(\\ S(x_2, x_1 + \hat{\mu}; \{U^*\})\, (1 + (\widetilde{C}\, \gamma^\mu\, \widetilde{C}^{-1})^*)\, (U^*)^{\mu\dagger}(x_1)\, S(x_1, 0; \{U^*\}) \\ - S(x_2, x_1; \{U^*\})\, (1 - (\widetilde{C}\, \gamma^\mu\, \widetilde{C}^{-1})^*)\, (U^*)^\mu(x_1)\, S(x_1 + \hat{\mu}, 0; \{U^*\}) \Big) ,\end{aligned} \tag{D.7}$$

where we can use Eq. (D.1) to get,

$$\widehat{S}(x_2, 0; t_1, \vec{q}, \mu; \{U^*\}) = \left(\widetilde{C}\, \widehat{S}(x_2, 0; t_1, -\vec{q}, \mu; \{U\})\, \widetilde{C}^{-1}\right)^* . \tag{D.8}$$

What we have shown is that, assuming we can find a suitable $\widetilde{C}$ satisfying Eq. (D.1), the original propagators from the $\{U\}$ gauge field configurations can be easily transformed into propagators from the $\{U^*\}$ configuration, saving much compute time. It is trivial to check that in the Sakurai γ matrix representation $\widetilde{C}$ has the form,

$$\widetilde{C} = C\, \gamma_5 , \tag{D.9}$$

and for the Bjorken-Drell representation one has,

$$\widetilde{C} = \gamma_0\, C\, \gamma_5 . \tag{D.10}$$

An example of how the U^* trick can be used to reduce statistical errors is given in App. E.

Appendix E

Form factor correlation function ratios

Here we investigate the choice of momenta in forming the ratios of two- and three-point correlation functions in Eqs (8.37) and (8.25). A sample calculation of the Δ^+ decuplet form factors is performed so that we can compare the two approaches.

E.1 Correlation function ratios

The ratio under consideration is,

$$\begin{aligned} R_\sigma{}^\mu{}_\tau(t_2,t_1;\vec{p}\,',\vec{p};\Gamma) &= \left(\frac{\left\langle\, G^{Bj^\mu B}_{\sigma\tau}(t_2,t_1;\vec{p}\,',\vec{p};\Gamma)\,\right\rangle \left\langle\, G^{Bj^\mu B}_{\sigma\tau}(t_2,t_1;\vec{p},\vec{p}\,';\Gamma)\,\right\rangle}{\left\langle\, G^{BB}_{\sigma\tau}(t_2;\vec{p}\,';\Gamma_4)\,\right\rangle \left\langle\, G^{BB}_{\sigma\tau}(t_2;\vec{p};\Gamma_4)\,\right\rangle} \right)^{1/2} \\ &\simeq \left(\frac{E_p+M}{2E_p}\right)^{1/2} \left(\frac{E_{p'}+M}{2E_{p'}}\right)^{1/2} \overline{R}_\sigma{}^\mu{}_\tau(\vec{p}\,',\vec{p};\Gamma)\,. \end{aligned} \tag{E.1}$$

By constructing this ratio of correlation functions the decuplet form factors can be extracted using Eqs. (8.38) to (8.41). We prefer this choice over that used previously in Refs. [205, 213],

$$R_\sigma{}^\mu{}_\tau(t_2,t_1;\vec{p}\,',\vec{p};\Gamma) = \left(\frac{\left\langle\, G^{Bj^\mu B}_{\sigma\tau}(t_2,t_1;\vec{p}\,',\vec{p};\Gamma)\,\right\rangle \left\langle\, G^{Bj^\mu B}_{\sigma\tau}(t_2,t_1;-\vec{p},-\vec{p}\,';\Gamma)\,\right\rangle}{\left\langle\, G^{BB}_{\sigma\tau}(t_2;\vec{p}\,';\Gamma_4)\,\right\rangle \left\langle\, G^{BB}_{\sigma\tau}(t_2;-\vec{p};\Gamma_4)\,\right\rangle} \right)^{1/2} \tag{E.2}$$

as the symmetric combination in Eq. (E.1) results in a cancellation of the momentum dependent $Z_B(p)$ factors describing the overlap of the smeared interpolators with the baryon.

However, in contrast to the previous method, the symmetric choice (E.1) requires an extra SST propagator with momentum $\vec{p} = (-1, 0, 0)$. Fortunately, this propagator must be calculated anyway in order to apply the $U + U^*$ trick (see App. D), and this is not a problem.

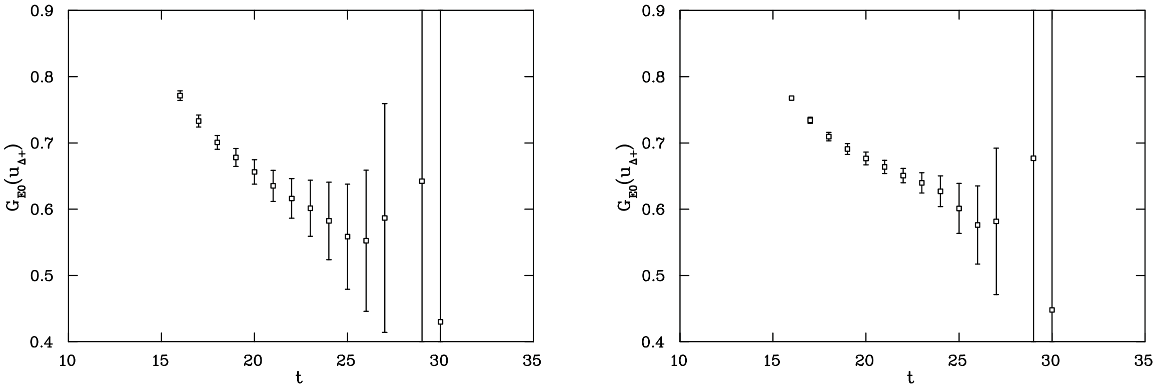

Figure E.1: The u quark contribution to the $E0$ form factor of the Δ^+ baryon. The figure on the left was calculated using the previous choice for construction the correlation function ratio (E.2). The new choice (E.1) was used to calculate the u quark contribution shown on the right. The U^* trick was not applied. We see a reduction in the statistical uncertainty when using the new approach.

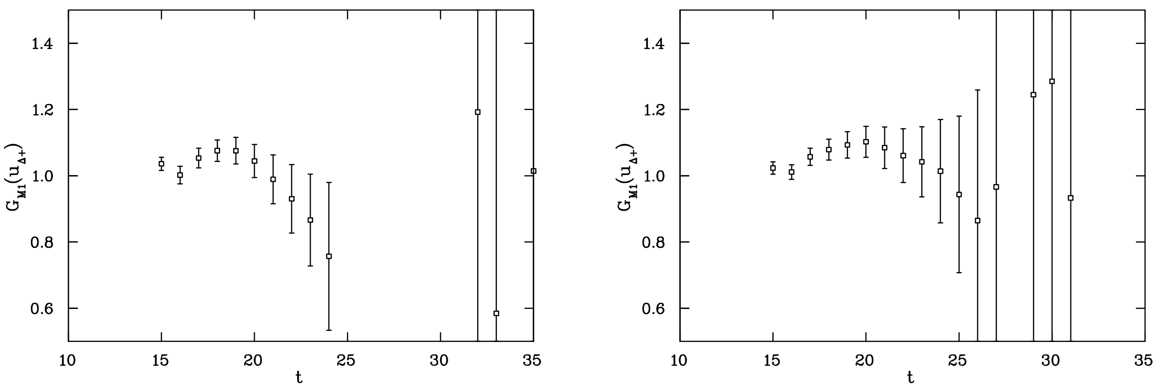

Figure E.2: Same as for Fig. E.1 but for the $M1$ form factor.

E.2 Comparison

We now compare the two approaches with a small calculation over a subset of 10 configurations from the Lüscher-Weisz quenched ensemble discussed in Chapter 8. Figure E.1 shows the u-quark sector contribution to the $E0$ form factor using both possible choices of correlation function ratios, without the U^* trick. We see a reduction in the statistical errors when using the new approach. This is also true for the $M1$ form factor as seen in Fig. E.2. The decrease in statistical uncertainty is due to the extra SST propagator required to form the new correlation function ratio. The extra momentum averaging this introduces leads to a reduction in statistical error via an improved estimator.

In Figs. E.3 and E.4 we repeat the calculation, this time including the U^* trick. By including the U^* trick the previous approach also receives the benefit from the extra momentum averaging introduced by the SST propagators. Therefore, the statistical errors for both choices are this time exactly equal and smaller than without the U^* trick.

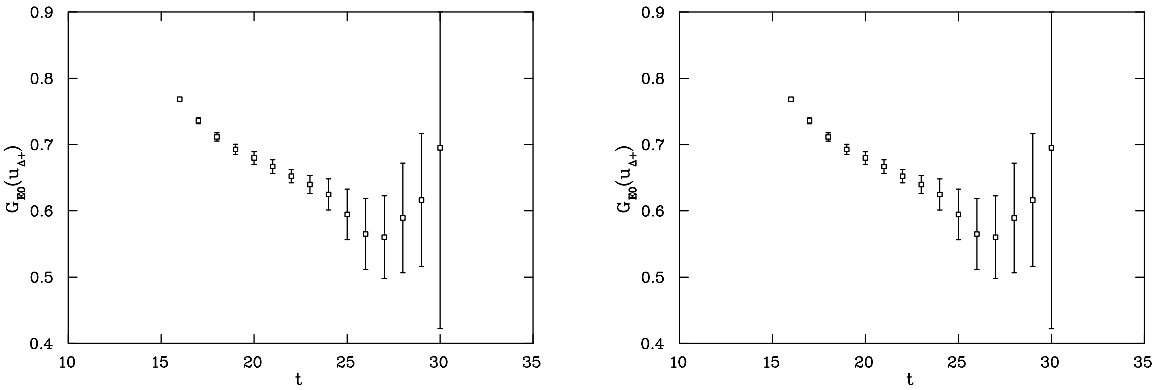

Figure E.3: Same as for Fig. E.1, however this time we include the U^* trick. This time both approaches are exactly equal.

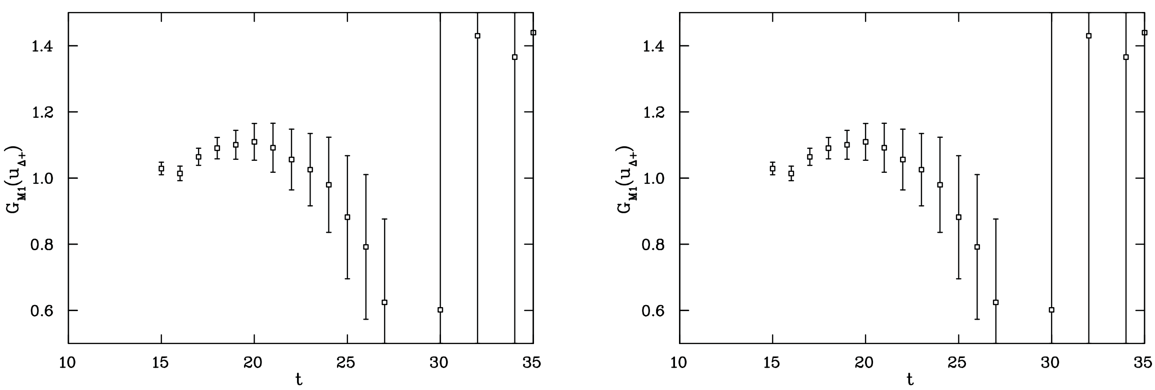

Figure E.4: Same as for Fig. E.3 but for the $M1$ form factor.

Appendix F

Reduce code for calculating decuplet form factors

Reduce scripts for calculating the required combinations of two- and three-point correlation functions used to extract the decuplet baryon factors.

F.1 definitions.red

Provides some useful definitions.

```
%----
%  output formatting
ON div;

%  separation of components
ON factor;

%  disable traces over the line, L
NOSPUR L;

%----
%  variable definitions:
%
%  unit vectors
VECTOR e0,e1,e2,e3,e4;

%  using Bjorken and Drell representation => g_{\mu\nu} metric
LET e0.e0 = 1; LET e0.e1 =  0; LET e0.e2 =  0; LET e0.e3 =  0;
               LET e1.e1 = -1; LET e1.e2 =  0; LET e1.e3 =  0;
                               LET e2.e2 = -1; LET e2.e3 =  0;
                                               LET e3.e3 = -1;
LET e4 = e0;
LET eps(e0,e1,e2,e3) = 1;

%  useful defintions
VECTOR mu, nu, sig, ome;
INDEX sig, ome, al, b;
ORDER a1, a2, c1, c2;

%----
%  \Gamma matrices,
%
%  \Gamma_0 = \Gamma_4 = 1/2 ((I,0),(0,0))
%
%  Recall: in Dirac/Bjorken and Drell representation,
%          \gamma_0 = ((1,0,0,0),(0,1,0,0),(0,0,-1,0),(0,0,0,-1))
```

```
%      =>  \Gamma_0 = 1/4 (\gamma_0 + 1)
%
%  similarly for \Gamma_i,
%
%  \Gamma_i = 1/4 (\gamma_5 \gamma_0 \gamma_i - \gamma_5 \gamma_i)
%
FOR ALL L,mu LET Gam(L,mu) = (1/4)*(G(L,a,e0,mu) - G(L,a,mu));
FOR ALL L    LET Gam(L,e0) = (1/4)*(G(L,e0) + 1);
FOR ALL L    LET Gam(L,e4) = (1/4)*(G(L,e4) + 1);

%----
%  Rarita-Schwinger spin sum,
%
%   \sum_s u_sig(p,s) ubar_ome(p,s) = -(gam.p + M_D)/(2M_D) (
%     g_sigome - 1/3 gam_sig gam_ome - (2 p_sig p_ome)/(3 M_D^2) )
%
FOR ALL L,sig,ome,p,m
LET Spin32(L,sig,ome,p,m) = -(G(L,p) + m)/(2*m)*(
        (sig.ome) - (1/3)*G(L,sig,ome) - 2*(p.sig)*(p.ome)/(3*m^2)
        + ((p.sig)*G(L,ome) - (p.ome)*G(L,sig))/(3*m) );

%  O^{\alpha\mu\beta}
%
FOR ALL L,alpha,mu,beta,pp,p,m
    LET O(L,alpha,mu,beta,pp,p,m) =
      -(alpha.beta)*(a1*G(L,mu) + (a2/(2*m))*(pp.mu + p.mu)) -
       ((pp.alpha - p.alpha)*(pp.beta - p.beta))/((2*m)^2)*
       (c1*G(L,mu) + (c2/(2*m))*(pp.mu + p.mu));

%----
%  Correlation functions,
%
FOR ALL p,m LET En(p,m) =
    sqrt((p.e1)^2 + (p.e2)^2 + (p.e3)^3 + m^2);

% Include factor of 4 because reduce returns 1/4 of trace
FOR ALL L,p,sig,ome,m LET
    g2(L,p,sig,ome,m) =
       4*ZB*ZB*(m/En(p,m))*
       Gam(L,e0)*Spin32(L,sig,ome,p,m);

% Include factor of 4 because reduce returns 1/4 of trace
FOR ALL L,sig,ome,mu,pp,p,igam,m LET
    g3(L,sig,mu,ome,pp,p,igam,m) =
       4*Gam(L,igam)*
       ZB*sqrt(MB/En(pp,m))*Spin32(L,sig,al,pp,MB)*
       sqrt(MB^2/(En(p,m)*En(pp,m)))*O(L,al,mu,b,pp,p,MB)*
       ZB*sqrt(MB/En(p,m))*Spin32(L,b,ome,p,MB);
```

F.2 twopoint.red

Calculates the two-point correlation functions in Eqs. (8.18) to (8.21).

```
in "definitions.red";

%----
%  Define momenta, pplus  = (E_p,p_x,0,0)
%                  pzero  = (MB,0,0,0)
%                  pminus = (E_p,-p_x,0,0)
%
VECTOR pplus, pzero, pminus;
MASS pplus=MB, pzero=MB, pminus=MB;
MSHELL pplus, pzero, pminus;

LET pplus.e0 = E_p, pzero.e0 = MB, pminus.e0 = E_p;
```

```
LET pplus.e1 = p_x, pzero.e1 = 0,  pminus.e1 = -p_x;
LET pplus.e2 = 0,   pzero.e2 = 0,  pminus.e2 = 0;
LET pplus.e3 = 0,   pzero.e3 = 0,  pminus.e3 = 0;

%LET sqrt(p_x^2 + MB^2) = E_p;
%LET p_x^2 + MB^2 = E_p;
LET abs(mb) = mb;
LET abs(E_p) = E_p;
%LET E_p = p_x^2 + MB^2
LET p_x^2 = E_p^2 - MB^2;

ON FACTOR;

LET ZB = 1;

g2(pplus,e0,e0,mb);
g2(pplus,e1,e1,mb);
g2(pplus,e2,e2,mb);
g2(pplus,e3,e3,mb);

END;
```

F.3 threepoint.red

Calculates the required combinations of correlation function ratios in Eqs. (8.38) to (8.41).

```
in "definitions.red";

%----
%  Define momenta, pplus  = (E_p,p_x,0,0)
%                 pzero  = (MB,0,0,0)
%                 pminus = (E_p,-p_x,0,0)
%
% Note that we use E_p = MB*(2*tau + 1), which
% should be checked for values of p' and p used.
% [Relation is true for momentum values used in
%  Leinweber and Hedditch work]
%
VECTOR pplus, pzero, pminus;
MASS pplus=MB, pzero=MB, pminus=MB;
MSHELL pplus, pzero, pminus;

LET pplus.e0 = E_p, pzero.e0 = MB, pminus.e0 = E_p;
LET pplus.e1 = p_x, pzero.e1 = 0,  pminus.e1 = -p_x;
LET pplus.e2 = 0,   pzero.e2 = 0,  pminus.e2 = 0;
LET pplus.e3 = 0,   pzero.e3 = 0,  pminus.e3 = 0;

LET eps(pplus, e1,e2,e3) = pplus.e0  * eps(e0,e1,e2,e3);
LET eps(pzero, e1,e2,e3) = pzero.e0  * eps(e0,e1,e2,e3);
LET eps(pminus,e1,e2,e3) = pminus.e0 * eps(e0,e1,e2,e3);

LET eps(pplus, e0,e2,e3) = pplus.e1  * eps(e0,e1,e2,e3);
LET eps(pzero, e0,e2,e3) = pzero.e1  * eps(e0,e1,e2,e3);
LET eps(pminus,e0,e2,e3) = pminus.e1 * eps(e0,e1,e2,e3);

LET eps(pzero,pplus,e2,e3)  = - pzero.e0 * pplus.e1 * eps(e0,e1,e2,e3)
                              + pzero.e1 * pplus.e0 * eps(e0,e1,e2,e3);
LET eps(pzero,pminus,e2,e3) = - pzero.e0 * pminus.e1 * eps(e0,e1,e2,e3)
                              + pzero.e1 * pminus.e0 * eps(e0,e1,e2,e3);
LET eps(pplus,pminus,e2,e3) = - pplus.e0 * pminus.e1 * eps(e0,e1,e2,e3)
                              + pplus.e1 * pminus.e0 * eps(e0,e1,e2,e3);

LET pplus.pzero  = MB*E_p;
LET pminus.pzero = MB*E_p;
%LET pplus.pminus = E_p*MB - p_x^2;
```

```
%LET sqrt(p_x^2 + MB^2) = E_p;
%LET p_x^2 + MB^2 = E_p^2;
%LET abs(E_p) = E_p;

LET abs(MB) = MB;
%LET p_x = sqrt(E_p^2 - MB^2);

%----
% Factor out p_x and E_p, replace with MB and tau
% *** NOTE ***
% Here we assume one of p', p is zero and the
% other is p+ or p-, this affects the value
% of tau.
%
LET p_x = 2*MB*sqrt(tau*(tau + 1));
LET E_p = MB*(2*tau + 1);

OFF FACTOR;

% Don't care about ZB factors
LET ZB = 1;

% Check that G3(\vec{p},0) = G3(0,-\vec{p})
%
LET Rele11 = g3(L3,e1,e0,e1,pplus,pzero,e0,MB);
LET Rele22 = g3(L3,e2,e0,e2,pplus,pzero,e0,MB);
LET Rele33 = g3(L3,e3,e0,e3,pplus,pzero,e0,MB);

LET Sele11 = g3(L3,e1,e0,e1,pzero,pplus,e0,MB);
LET Sele22 = g3(L3,e2,e0,e2,pzero,pplus,e0,MB);
LET Sele33 = g3(L3,e3,e0,e3,pzero,pplus,e0,MB);

LET Rmag11 = g3(L3,e1,e3,e1,pplus,pzero,e2,MB);
LET Rmag22 = g3(L3,e2,e3,e2,pplus,pzero,e2,MB);
LET Rmag33 = g3(L3,e3,e3,e3,pplus,pzero,e2,MB);

LET Smag11 = g3(L3,e1,e3,e1,pzero,pplus,e2,MB);
LET Smag22 = g3(L3,e2,e3,e2,pzero,pplus,e2,MB);
LET Smag33 = g3(L3,e3,e3,e3,pzero,pplus,e2,MB);

% Set normalisations due to dividing by 2-pt correlation
% functions (calculate these by hand).
% *** note for NORM11 we assume one of the momenta
% *** will be zero
%
% NORM11 = MB/E_p, with E_p = MB*2(tau + 1)
LET NORM11 = (3/2)*(1/(2*tau + 1));
LET NORM22 = (3/2);
LET NORM33 = (3/2);

% NORM11 for both non-zero momenta
% NORM11 = (3/2)*(1/(2*tau + 1)^2);

% Plus include factor of (MB + Ep)/(2Ep), again this
% assumes one of the momenta are zero. Write in terms
% of tau.
%
LET NORM = (2*tau + 1)/(tau + 1);

ON FACTOR;

% Finally need norm for ele0 and ele2
%
LET ELE0 = (1/3);
LET ELE2 = 2*MB*(En(pplus,MB) + MB)/p_x^2;
LET MAG1 = (-3/5)*(En(pplus,MB) + MB)/p_x;
LET MAG3 = -4*MB*(En(pplus,MB) + MB)^2/p_x^3;

ELE0*NORM*(NORM11*Rele11 + NORM22*Rele22 + NORM33*Rele33);
```

```
ELE2*NORM*(NORM11*Rele11 + NORM22*Rele22 - 2*NORM33*Rele33);
MAG1*NORM*(NORM11*Rmag11 + NORM22*Rmag22 + NORM33*Rmag33);
MAG3*NORM*(NORM11*Rmag11 + NORM22*Rmag22 - (3/2)*NORM33*Rmag33);

END;
```

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