Not a prescription but an identity: Mandelstam-Leibbrandt.

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ABSTRACT

After some considerations and coincidences that appear when working in the light-cone gauge, in both the Mandelstam-Leibbrandt prescription and the covariantization method, we suspect that there must be some connection between them. This work shows that we were right and it is practically trivial to demonstrate that relationship. And since the covariantization method is not a prescription, this implies that the results of Mandelstam and Leibbrandt are not prescriptions too, but they are identities from light-cone gauge coordinates.

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Leading with true degrees of freedom was one of the relishes of Dirac as he showed us around in his famous paper “On Hamiltonian Forms” [1]. This should pointed out as the very beginning of light-cone gauge. Since then, there were periods where people payed special attention on it, as an example, the case of quantization of the (super)string [2]. But in quantum field theories some troubles appeared. The use of light-cone gauge in (super)Yang-Mills theories ought to use a prescription in order to avoid the pathological poles that arise on the propagator structure. The structure of this kind of poles brings the theory upon problems, just to mention three of them: Wick rotation is not allowed, naive power counting cannot be used and non physical divergences are obtained from computing Feynman integrals. The form of this poles, also known in the context of non covariant gauges as spurious poles, is

\[ \frac{1}{q \cdot n} = \frac{1}{q^+}, \]

where \( q \) is a vector of Minkowski spacetime of signature \((+, -, -, -)\) and \( n \) is a light-cone vector.

The use of the Principal Value (PV) prescription in Feynman integrals leads to not desire results, so another prescriptions were invented. The first that appear was proposed by S. Mandelstam in 1983 when he was working on \( N = 4 \) super Yang-Mills theories. One year later, G. Leibbrandt proposed another one, now in the context of Feynman integral calculations. These prescriptions were in fact one [3]. But more than solve the problems cited above, and many others, the Mandelstam-Leibbrandt (ML) prescription preserves causality, and this seems to be the mandatory property that all well behaved prescriptions [4] must have on the light-cone gauge. An interesting work on the studies of symmetries of poles in the axial gauge is found in reference [5]. One question may arise now, is a prescription really needed, mandatory, or
maybe there are other ways to obtain the same results but with prescriptionless methods?. Going in this way, A. Suzuki et al. [6] found another exit. But there is also another technique, coined by A. Suzuki [7] as covariantization, to treat light-cone integrals. This technique is causal and reproduces the results obtained through the use of the Mandelstam-Leibbrandt prescription [8]. This fact remains almost forgotten, even though the covariantization manifestly guarantees the absence of zero mode frequencies that spoil causality. Of course, here the technique uses parametric integrations and integration over components, and all the technology of complex analysis to treat the singularities in a proper way. The thrust of the technique lies in that it does not require a prescription for the light-cone pole; it “converts” this pole in a “covariant” pole whose treatment is well-grounded and established since the early days of quantum field theory. The burden of this technique and its most severe drawback is that it requires an additional parametric integration to be performed, a task which can be very demanding.

1 Light-cone coordinates.

Since our special purpose is to work with the light-cone gauge in quantum field theory, then it is necessary to start, in our case very briefly, with an introduction to light-cone coordinates. Using a signature (+, −, −), then a general contravariant four-vector is given by the coordinates:

\[ x^\mu = (x^+, x^-, x^i) , \ i = 1, 2, \]

where \( x^\pm \) are defined respect to usual Minkowski coordinates as:

\[ x^\pm = \frac{1}{\sqrt{2}}(x^0 \pm x^3), \]
\[ = \frac{1}{\sqrt{2}}(x_0 \mp x_3) = x_\mp, \]
Now, defining the dual base light-like four-vectors:

\[ n_\mu = \frac{1}{\sqrt{2}} (1, 0, 0, 1), \]
\[ m_\mu = \frac{1}{\sqrt{2}} (1, 0, 0, -1), \] \hspace{1cm} (1)

we observe that with the help of this base, the \( x^\pm \) coordinates can be expressed as

\[ x^+ = x^\mu n_\mu, \]
\[ x^- = x^\mu m_\mu. \]

Using this coordinates, the scalar product becomes:

\[ x^\mu y_\mu = x^+ y^- + x^- y^+ - \hat{x}\hat{y}, \]

i.e.,

\[ x^2 = 2x^+ x^- - \hat{x}^2, \] \hspace{1cm} (2)

we will use this last expression later. Also we have that the metric is now:

\[ g^{\mu\nu} = \begin{pmatrix} + & - & 1 & 2 \\ 0 & 1 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} \]

Now we are ready to go on.

\section{The ML prescription.}

In this section we will resume the ML prescription, showing how the Wick rotation is modified for the \((q.n)^{-1}\) poles when using this causal prescription.
First, the prescription reads
\[ \frac{1}{q^+} = \frac{1}{q.r} \rightarrow \frac{q^-}{q^+ q^- - \epsilon}. \]

(3)

As said above, this prescription was first proposed by S. Mandelstam and one year after, in 1983 by G. Leibbrandt. One of the important issues this prescription takes on is the fact that Wick rotation can be done.

## 2.1 Wick rotation.

If one decides to use the path integral formulation of quantum field theory, one of the problems we have to work out is the fact that the path integral is no well defined in Minkowski spacetime. This is a delicate matter that belongs to a field called axiomatic field theory and it is known as Wick rotation.

First, we will use this rotation in working out the covariant poles of type
\[ \frac{1}{q^2}, \]
this is shown in figure 1

When working with the light-cone gauge, a special care must be taken since poles of the type
\[ \frac{1}{q.r} \]
do not allow a Wick rotation if no causal prescription (e.g. ML prescription) is used. This is shown in figure 2 where both cases are considered: with and without the use of the ML prescription. Observe the location of poles in this two cases. This means that if we do not take care of use the ML prescription (or any other causal prescription) the result that we will obtain are axiomatically wrong since we are working as if we were in an Euclidean space when in fact we never leave the Minkowski spacetime. This explains
why all the pathologies are observed when the light-cone gauge is worked in
the same manner of the other gauges. So we have a key point here, the study
of Wick rotation in the light-cone gauge is crucial and should be considered
as mandatory.

2.2 Feynman integrals.

To see how this prescription works, we will take the following two integrals

$$A = \int d^2\omega q \cdot \frac{1}{(q - p)^2 q^+}$$

(4)

and

$$B = \int d^2\omega q \cdot \frac{1}{q^2 (q - p)^2 q^+}$$

(5)
as an examples. We are taking $D = 2\omega$ as the dimension of spacetime, i.e.;
the limit to four dimensions means $\omega \to 2$. Using the ML prescription given
in eq. (3) and replacing in the first integral (A) of eq. (4), we obtain

\[ A_{ML} = \int d^2 \omega \frac{q^-}{(q-p)^2(q^+q^- - \epsilon)} \]

the subscript \( ML \) remarks that we are using the ML prescription. At this point we have the chance of choosing a parametrization for the denominator, we will use the exponential parametrization (also known as Schwinger parametrization, as you can see in the appendix), then

\[
A_{ML} = -\int d^2 \omega \cdot q^- \int_0^\infty d\alpha d\beta \cdot e^{i(\alpha(q-p)^2 + \beta q^+ q^-)},
\]

\[
= -\int d^2 \omega \cdot q^- \int_0^\infty d\alpha d\beta \cdot e^{i\alpha p^2 \cdot e^{i(2q^+ q^- - \alpha q^2 - 2\alpha q^+ p^- - 2q^+ p^- + 2\alpha q^+ q^-)},}
\]

where we are working in light-cone coordinates. Interchanging the order of the integrals we obtain

\[
A_{ML} = -\int_0^\infty d\alpha d\beta \cdot e^{i\alpha p^2} \cdot \int d^2 \omega \cdot q^- \int d^2 \omega \cdot q^- e^{i(2q^+ q^- - \alpha q^2 - 2\alpha q^+ p^- - 2q^+ p^- + 2\alpha q^+ q^-)},
\]
\[
\begin{align*}
&= -\int_0^\infty d\alpha d\beta \cdot e^{i\alpha p^2}.
\end{align*}
\]
\[
\int dq^+ \cdot e^{-i(\alpha q^2 - 2\alpha \hat{q} \hat{p})} \int_q^- dq^- \cdot q^- e^{-2i\alpha q^- p^+} \int_q^+ dq^+ \cdot e^{i(2\alpha q^- - 2\alpha p^- + \beta q^-)q^+}
\]

now we are ready to work out the momentum integral. Starting with the \(q^+\) component, it is nothing but the integral representation of Dirac distribution, this makes the \(q^-\) integral trivial, and the \(\hat{q}\) integral is of gaussian type. After some work, the result obtained is

\[
A_{ML} = -i\frac{\alpha^{2-\omega}}{(\alpha + \frac{\omega}{2})^2} e^{2i\phi_{\omega} p^+ p^-},
\]
working the parameter integral: ¹

\[
A_{ML} = i(-\pi)^{\omega} p^- (2p^+ p^-)^{\omega-1} \frac{\Gamma(2 - \omega) \Gamma(\omega - 1)}{\Gamma(\omega)}. \tag{6}
\]

This is the final result for the integral \(A_{ML}\).

Let’s compute now the integral \(B\), this will follow the same steps as above, then

\[
B_{ML} = \int d^2q : \left( \frac{q^-}{q^2(q - p)^2(q^+ q^- - \epsilon)} \right).
\]

that means in this case,

\[
B_{ML} = i \int d^2q \cdot q^- \int_0^\infty d\alpha d\beta d\gamma \cdot e^{i(\alpha(q-p)^2 + \beta q^+ q^- + \gamma q^2)},
\]

\[
= i \int d^2q \cdot q^- \int_0^\infty d\alpha d\beta \cdot e^{i[2(\alpha + \gamma)q^+ q^- - (\alpha + \gamma)q^2 - 2\alpha q^+ p^- - 2\alpha q^- p^+ + 2\alpha \hat{q} \hat{p} + \beta q^+ q^-]},
\]

¹observe that here we have made a rescaling \(\beta: \beta \rightarrow 2\beta\).
now we have 3 parameters

\[ B_{ML} = i \int_0^\infty d\alpha d\beta d\gamma \cdot e^{i\alpha p^2} \cdot \int d^2q \cdot q^- e^{[2(\alpha + \gamma)q^+ q^- - (\alpha + \gamma)q^2 - 2\alpha q^- p^- - 2\alpha q^+ p^+ + 2\alpha q^+ p^+ + 2\alpha q^- p^- + 2\alpha q^+ p^+]}, \]

\[ = i \int_0^\infty d\alpha d\beta d\gamma \cdot e^{i\alpha p^2} \cdot \int dq^+ \cdot e^{-i[(\alpha + \gamma)q^2 - 2\alpha q^+ p^+] \int_0^\infty dq^- \cdot q^- e^{-2i\alpha q^- p^+}, \]

\[ = i \int d\alpha d\beta d\gamma \cdot e^{i\alpha p^2}, \]

\[ = i \int d\alpha d\beta d\gamma \cdot e^{i\alpha p^2}. \]

Finally

\[ B_{ML} = i^{\omega+1} p^\omega \int_0^\infty d\alpha d\beta d\gamma \cdot \frac{\alpha}{(\alpha + \gamma/2)(\alpha + \beta + \gamma/2)^2} \cdot e^{i\alpha p^2 + i(\alpha + \gamma/2)p^2 - 2i\alpha q^- p^+}, \]

\[ = i(-\pi)^\omega p^\omega (p^2)^{\omega-2} \sum_{n=0}^\infty \frac{\Gamma(3 - \omega + n)\Gamma(\omega - 1 + n)}{\Gamma(\omega + n)\Gamma(n + 2)} \Gamma(\omega + n + 2) (1 - \eta^{n+1}) \]

where we defined

\[ \eta \equiv -\frac{\not{p}^2}{p^2}. \]

These results of integrals (6) and (7) are well known in the literature of non covariant gauges of axial type [3, 9]. For practically twenty years they are used as a base for other more complicated results.

### 3 The covariantization method.

Now we are going to show how the calculus of Feynman integrals that present light-cone poles can be done without a prescription. We will introduce the
“covariantization” technique which was proposed by A. Suzuki [7]. The main idea is quite simple. When working with light-cone coordinates as was shown in eq. (2), the square of a four-momentum is:

\[ q^2 = 2q^+ q^- - \hat{q}^2, \]

As long as \( q^- \neq 0 \), we can write \( q^+ \) as

\[ q^+ = \frac{q^2 + \hat{q}^2}{2q^-} \]

We note that this dispersion relation almost guarantees that real gauge fields for which \( q^2 = 0 \) (real photons or real gluons for example) are transverse; the residual gauge freedom, that is left to be dealt with so that fields be manifestly transverse comes from the presence of the \( q^- \) in the denominator of the expression above. This implies that in the light-cone gauge the characteristic pole becomes

\[ \frac{1}{q^+} = \frac{2q^-}{q^2 + \hat{q}^2}, \tag{8} \]

The important thing here is that the condition \( q^- \neq 0 \) warranties the causal structure of this technique since it eliminates the troublesome \( q^- = 0 \) modes. Elimination of these modes restores the physically acceptable results as can manifestly be seen in the causal prescription [4] for the light-cone gauge.

### 3.1 Wick rotation.

As pointed out in the above section, the importance of Wick rotation has physical implications in the light-cone gauge, and this will be an important property that the covariantization method should have. From light-cone coordinates, we obtained

\[ \frac{1}{q^+} = \frac{2q^-}{q^2 + \hat{q}^2}. \]
This means that the denominator can be written as \((q^0)^2 - (q^3)^2\), making analytic continuation to the complex plane in the temporal component implies that the poles will be situated in

\[ q^0 = \pm |q^3| \mp i\epsilon \]

as is shown in figure 3.

![Figure 3: Poles that arise when using the covariantization.](image)

This shows that the covariantization method also preserves the Wick rotation, which, when violated, was the cornerstone of non-physical results. If it is true, the results of integrals \(A\) and \(B\) will coincide when using covariantization.
3.2 Feynman integrals.

Taking again the integrals (4) and (5), but now with the covariantization method, implies that

\[
A_{\text{cov}} = \int d^{2}\omega q \cdot \frac{2q^{-}}{(q - p)^{2}(q^{2} + \hat{q}^{2})}
\]

\[
A_{\text{cov}} = -2 \int d^{2}\omega q \cdot q^{-} \int_{0}^{\infty} d\alpha d\beta \cdot e^{i(\alpha(q-p)^{2} + 2\alpha\hat{q}p - q^{-})},
\]

here we have used the relation: \(q^{2} + \hat{q}^{2} = 2q^{+}q^{-}\). Doing this the appearance of \(A_{\text{cov}}\) is closely similar to that of \(A_{\text{ML}}\). This fact will help us when facing a possible relation between them.

\[
A_{\text{cov}} = -2 \int_{0}^{\infty} d\alpha d\beta \cdot e^{i\alpha p^{2}} \cdot \int d^{2}\omega q \cdot q^{-} \int_{0}^{\infty} d\alpha d\beta \cdot e^{i(2\alpha q^{+}q^{-} - \alpha\hat{q}^{2} - 2\alpha q^{+}p^{-} - 2\alpha q^{-}p^{+} + 2\alpha q^{-} + 2\beta q^{+}q^{-})},
\]

following the same steps as before

\[
A_{\text{cov}} = -2i^{\omega} \pi^{\omega} p^{-} \int_{0}^{\infty} d\alpha d\beta \cdot \frac{\alpha^{2-\omega}}{(\alpha + \beta)^{2}} e^{2i\omega p^{-} p^{+}}
\]

finally

\[
A_{\text{cov}} = i(-\pi)^{\omega} p^{-} (2p^{+}p^{-})^{\omega-1} \frac{\Gamma(2 - \omega)\Gamma(\omega - 1)}{\Gamma(\omega)},
\]

which is the same result as obtained in eq. (6). This is in favor of the covariantization method, but to remove any doubt, the same computation
will be done for integral $B$ and then compare it with the result of integral $B_{ML}$ of eq. (7). Then, we have

$$B_{\text{cov}} = \int d^{2}\omega \cdot \frac{2q^{-}}{q^{2}(q-p)^{2}(q^{2}-\hat{q}^{2})}$$

again, following the same steps as before.

$$B_{\text{cov}} = 2i \int d^{2}\omega \cdot q^{-} \int_{0}^{\infty} d\alpha d\beta d\gamma \cdot e^{i(\alpha(q-p)^{2}+2\beta q^{+}q^{-}+\gamma q^{2})},$$

$$= 2i \int d^{2}\omega \cdot q^{-} \cdot 
\int_{0}^{\infty} d\alpha d\beta d\gamma \cdot e^{i2(\alpha+\gamma)q^{+}q^{-}-(\alpha+\gamma)q^{2}-2\alpha q^{+}p^{-}-2\alpha q^{-}p^{+}+2\alpha\hat{q}p+2\beta q^{+}q^{-}],}$$

this implies

$$B_{\text{cov}} = 2i \int_{0}^{\infty} d\alpha d\beta d\gamma \cdot e^{i\alpha p^{2}} \cdot 
\int d^{2}\omega \cdot q^{-} \cdot e^{i2(\alpha+\gamma)q^{+}q^{-}-(\alpha+\gamma)q^{2}-2\alpha q^{+}p^{-}-2\alpha q^{-}p^{+}+2\alpha\hat{q}p+2\beta q^{+}q^{-}],}$$

$$= 2i \int_{0}^{\infty} d\alpha d\beta d\gamma \cdot e^{i\alpha p^{2}} \cdot 
\int d\hat{q} \cdot e^{-i[(\alpha+\gamma)q^{2}-2\alpha\hat{q}p]} \int_{-\infty}^{\infty} dq^{-} \cdot q^{-} e^{-2i\alpha q^{-}p^{+}} \cdot 
\int_{-\infty}^{\infty} dq^{+} \cdot e^{i2(\alpha+\gamma)q^{+}q^{-}+2\beta q^{+}q^{-}]q^{+}}$$

doing the momentum integrals,

$$B_{\text{cov}} = 2i^{\omega+1} \pi^{\omega} \int_{0}^{\infty} d\alpha d\beta d\gamma \cdot \frac{\alpha}{(\alpha+\beta+\gamma)^{2}} \cdot e^{i\alpha p^{2}+i(\alpha+\gamma)\hat{p}^{2}-2i\frac{\beta^{2}}{(\alpha+\beta+\gamma)}p^{+}p^{-}},$$
doing the parameter integrals,
\[ B_{\text{cov}} = i(-\pi)^{\omega}p^{-}(p^2)^{-\omega-2} \sum_{n=0}^{\infty} \frac{\Gamma(3 - \omega + n)\Gamma(\omega - 1 + n)}{\Gamma(\omega + n)\Gamma(n + 2)} (1 - \eta^{n+1}) \] (10)
where again we defined
\[ \eta \equiv \frac{\hat{p}^2}{p^2} \]
this result agrees with that of \( B_{ML} \).
This is not an ill-fated result. Instead, this makes us to think about the possibility of a certain kind of relationship between the Mandelstam-Leibbrandt prescription and the covariantization method. And that will be the motif of the following section.

4 The “proof”.

This proof actually is not a real complicated proof, it is in fact almost trivial as we can see now.
Let’s take the form of a more general one loop Feynman integral in the light-cone gauge,
\[ I(p) = \int q.d^{2}\omega \frac{f(q,p)}{f_1(q,p) \cdots f_n(q,p)} \frac{1}{q^+}. \] (11)
Making use of the ML prescription, the integral (11) takes the form
\[ I_{L}(p) = \int q.d^{2}\omega \frac{f(q,p)}{f_1(q,p) \cdots f_n(q,p)} \frac{q^-}{q^+q^- + i\epsilon}, \]
after the Schwinger parametrization
\[ I_{L}(p) = (-i)^{n+1} \int q.d^{2}\omega 0 \int d\alpha_1 \cdots d\alpha_n d\beta \cdot q^- e^{i \left[ \sum_{i=1}^{n} \alpha_i f_i(q,p) + \beta q^+ q^- \right]} . \] (12)
But if instead of using the prescription, we use the covariantization method, we obtain
\[ I_{\text{cov}}(p) = \int q.d^{2}\omega \frac{f(q,p)}{f_1(q,p) \cdots f_n(q,p)} \frac{2q^-}{q^2 + \hat{q}^2}, \]
Schwinger parametrization implies

\[ I_{\text{cov}}(p) = (-i)^{n+1} \int q.d^{2}\omega \int_{0}^{\infty} d\alpha_{1} \cdots d\alpha_{n} d\beta \cdot 2q^{-} e^{i \left[ \sum_{1}^{n} \alpha_{i} f_{i}(q,p) + \beta 2q^{+} q^{-} \right]}, \]  

where the term that multiplies the \( \beta \) parameter was changed using again the identity of light-cone coordinates \( q^{2} + q^{2} \equiv 2q^{+}q^{-} \), in this way, the resemblance between the eqs. (12) and (13) is more evident. In fact, if the following change of variable is performed

\[ \beta \rightarrow 2\beta, \]
\[ d\beta \rightarrow 2d\beta, \]

in eq. (12), this will coincide with that of eq. (13). And also this is saying us that the covariantization method and the ML prescription are related throughout a rescaling of coordinates in the space of parameters.

5 Conclusions.

In this work we have shown that there exist a very close relationship between the wide used Mandelstam-Leibbrandt prescription and the covariantization method. Also, it is known that the covariantization method is related \[10\] to another technique, the negative dimension integration (NDIM) for light-cone gauge \[6\], which is considered as a prescriptionless technique too. Since the covariantization was just derived from a light-cone coordinates identity,

\[ \frac{1}{q^{+}} = \frac{2q^{-}}{q^{2} + q^{2} + q^{-} q^{-}}, \]

this seems that something similar can be do for the ML prescription, in this case,

\[ \frac{1}{q^{+}} = \frac{q^{-}}{q^{+}q^{-} - \epsilon} \rightarrow \frac{q^{-}}{q^{+}q^{-}}. \]
in both cases the same limit is reached, and in this limit the important issue
is the pole structure
\[
\frac{1}{q^+q^-},
\]
which guarantees the Wick rotation and all well behaved results of the theory.

Summing up this work in few words, the ML prescription in fact should be
called as the Mandelstam-Leibbrandt identity or the Mandelstam-Leibbrandt
relation.

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A Schwinger parametrization.

The Schwinger parametrization consists on the following, starting from the
identity
\[
\frac{1}{A} = -i \int_0^\infty d\alpha \cdot e^{i\alpha A},
\]
it is possible generalize this to the case where we have \( n \) denominators
\[
\frac{1}{A_1A_2\cdots A_n} = (-i)^n \int_0^\infty d\alpha_1 d\alpha_2 \cdots d\alpha_n e^{i \sum_{j=1}^n \alpha_j A_j}
\tag{14}
\]
in the case of working with the light-cone integrals, the form of the denomi-
nators are
\[
\frac{1}{(q - a_1)^2(q - a_2)^2\cdots(q - a_n)^2 [[q^+]}} = (-i)^n \int_0^\infty d\alpha_1 d\alpha_2 \cdots d\alpha_n d\beta e^{i[\alpha_1 q^2 - 2a_1, q] + i\alpha_2 a_2^2 + i\beta[q^+]} 
\]
where \( [q^+] \) stands for that we are using some causal prescription or the covariantization method.

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