Helicity Amplitudes for production of massive gravitino/goldstino

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Abstract. The spinor helicity formalism (SHF) has been efficiently applied to compute perturbative amplitudes in plenty of processes and reactions in gauge theories (including gravity), mostly in the massless case. Some work has been done in order to extend the SHF to the massive case. We have used these powerful tools to evaluate amplitudes in a local supersymmetric model where the gravitino is the lightest supersymmetric particle (LSP). Two decays have been evaluated in order to show the capability of the SHF in the massive extension, namely the two body neutralino decays $\tilde{\chi}_0 \rightarrow V \tilde{\Psi}$ with $V = \gamma, Z$. The comparisons of amplitudes with spin-3/2 gravitino and spin-1/2 goldstino are also presented.

1. Introduction

The traditional Feynman approach for perturbation theory has been developed for almost sixty years, and it has been undoubtedly a magnificent tool to evaluate amplitudes in gauge theories. Even with all the success of the Feynman rules, several computations for scattering processes especially in QCD suffer the lack of accuracy in order to be tested experimentally. Furthermore, processes with several external particles are extremely complicated to calculate with the traditional perturbative approach [1]. Recent developments for tree-level scattering amplitudes in the massless case using the helicity methods are now widely used [2, 3, 4, 5, 6, 7], allowing to efficiently evaluate amplitudes for processes with $n$ external legs in pure Yang-Mills theories. A huge theoretical progress has been done during the last two decades in this active and growing field, making manifest certain mathematical structures that underlie quantum field theories [8, 9, 10, 11, 12, 13].

We shall try to present a pragmatic approach in this work, starting still à la Feynman, building the scattering amplitudes from the Feynman diagrams but expressing the usual spinors (in any representation) appearing in the amplitudes as momentum twistors, this step shall hugely facilitate the calculations, as we shall see later. This is one of the goals of the SHF. Addressing the massive case requires the implementation of the so-called light cone decomposition (LCD) for the massive momenta [14, 15, 16, 17]. This technique allow us to express a massive momentum, namely $p^2 = -m^2$ as a lineal combination of two massless momentum $p_i = r_i + \alpha q_i$ where $r_i^2 = q_i^2 = 0$ and $\alpha = -\frac{m^2}{2q_i \cdot r_i}$. Several processes of the electroweak Standard Model have been evaluated using the massive SHF in Ref. [18], reproducing the well known results. More recently we have implemented this method for $\mathcal{N} = 1$ supergravity with LSP gravitinos [19], showing that the formalism is well suitable for amplitudes involving massive gravitinos in the final state.
Considering scenarios where the gravitino is the LSP and hence a good dark matter candidate, the nature of the next-to-lightest supersymmetric particle (NLSP) determines its phenomenology [20, 21, 22, 23, 24]. Potential candidates for NLSP include the lightest neutralino [25, 26], the chargino [27], and the lightest charged slepton [28]. The weakness of the gravitational interactions suggest that the NLSP could have a long lifetime leading to scenarios with a metastable charged sparticles that could have striking signatures at colliders [29, 30] and it could also affect the Big Bang nucleosynthesis [31, 32, 33]. Here we shall consider several decay modes for the neutralino as the NLSP. We also have considered the case when the gravitino can be approximated by the goldstino state, this is for the region where the gravitino mass is small compared to the neutralino mass \( \tilde{m} \ll m_{\tilde{\chi}_0} \). This proceeding contribution is structured as follows: Section 2 contains the solutions of the Rarita-Schwinger equation for the gravitino. The four states of the gravitino are expressed in terms of the momentum twistors after LCD is applied. We shall see that this is the key to obtain compact expression for the decay amplitudes. In Section 3 the helicity amplitudes (HAs) for the two decays \( \tilde{\chi}_0 \rightarrow \bar{V} \tilde{\Psi}_\mu \) with \( V = \gamma, Z \) are shown in terms of very compact mathematical expressions, in addition the HAs with the goldstino approximation are also shown. Finally in Section 4 we present some final remarks and comments.

2. Gravitino wave functions
2.1. Rarita-Schwinger equations
The wave function for massive gravitino is the solution of the Rarita-Schwinger equation [34], this equation of motion (EOM) results from applying the Euler-Lagrange to the following lagrangian [35]

\[
\mathcal{L} = -\frac{1}{2} \epsilon^{\mu\nu\rho\sigma} \Psi_\mu C^I \gamma_5 \gamma_\nu \partial_\rho \Psi_\sigma + \frac{1}{4} \tilde{m} \Psi_\mu C^I [\gamma_\mu, \gamma_\nu] \Psi_\nu,
\]

the EOM for the gravitino are equivalent to the following equations [36]

\[
\gamma^\mu \Psi_\mu = 0, \\
\partial^\mu \Psi_\mu = 0, \\
(\gamma^\mu \partial_\mu - \tilde{m}) \Psi_\mu = 0,
\]

where \( \tilde{m} \) is the gravitino mass. Through all this work we shall use the following convention for the Minkowski metric \( \eta^{\mu\nu} = \text{diag}(-1,1,1,1) \), besides we have used the Dirac representation for the gamma matrices where they take the following form

\[
\gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{pmatrix},
\]

with \( \sigma^\mu = (1, \vec{\sigma}) \) and \( \bar{\sigma}^\mu = (1, -\vec{\sigma}) \). Returning to the EOM Eqs. (2)-(4), these admit the following solution [36]

\[
\tilde{\Psi}_\mu (\vec{p}, \lambda) = \sum_{s,m} \left\{ \begin{pmatrix} 1 \\ 2 \end{pmatrix} \begin{pmatrix} 1 \\ 2 \end{pmatrix} \right\} \begin{pmatrix} 3 \\ 2 \end{pmatrix} \begin{pmatrix} 1 \\ m \end{pmatrix} u(\vec{p}, s) \epsilon_\mu (\vec{p}, m),
\]
this is a Clebsch-Gordon expansion with coefficient \((\frac{1}{2}, \frac{3}{2}) (1, m) \left| (\frac{3}{2}, \lambda)\right.\), from Eq. (6), we know that there are four gravitino states, explicitly they are as follows

\[
\tilde{\Psi}^\mu_+(p) = \epsilon^\mu_+(p)u_+(p), \\
\tilde{\Psi}^\mu_-(p) = \epsilon^\mu_+(p)u_-(p), \\
\tilde{\Psi}^\mu_+(p) = \sqrt{\frac{2}{3}} \epsilon^\mu_0(p)u_+(p) + \frac{1}{\sqrt{3}} \epsilon^\mu_+(p)u_-(p), \\
\tilde{\Psi}^\mu_-(p) = \sqrt{\frac{2}{3}} \epsilon^\mu_0(p)u_+(p) + \frac{1}{\sqrt{3}} \epsilon^\mu_+(p)u_-(p).
\]

Replacing the momentum twistors in the four gravitino states Eqs. (7)-(10) is straightforward (See Ref. [19] for more details). Dealing with HAs involving gravitinos in the final state expressed now in these new variables shall avoid the large and messy expressions that appear when the trace technology is used (most of the time handled with \textsc{Mathematica}), we have to remember that the completeness relation for the gravitino takes the form

\[
D_{\mu\nu}(p) = \sum_{\lambda=1}^{3} \tilde{\Psi}_\mu(p, \tilde{\lambda}) \tilde{\Psi}_\nu(p, \tilde{\lambda}) = -\left(\not{p} + \tilde{m}\right) \times \left[ \left( g_{\mu\nu} - \frac{p_\mu p_\nu}{\tilde{m}^2}\right) \\
- \frac{1}{3} \left( g_{\mu\sigma} - \frac{p_\mu p_\sigma}{\tilde{m}^2}\right) \left( g_{\nu\lambda} - \frac{p_{\nu} p_\lambda}{\tilde{m}^2}\right) \gamma^\sigma \gamma^\lambda \right].
\]

Fortunately in the SHF the fundamental building block are the HAs (without square modulus) and we shall not need Eq. (11), neither another completeness relation in general.

\subsection{2.2. High energy equivalence theorem}

We briefly discuss in this section the gravitino approximation to goldstino. When the gravitino mass \((\tilde{m})\) is small compared to the energy of the process, it is possible to apply the equivalence theorem which roughly speaking allow us to replace the longitudinal components of the gravitino by the derivative of the goldstino field [37, 38, 39, 40, 41, 42]. In practice the equivalence theorem tell us that the four gravitino states Eqs. (7)-(10), go as follows; \(\tilde{\Psi}^\mu_+(p) \approx 0, \tilde{\Psi}^\mu_-\approx 0, \tilde{\Psi}^\mu_-\approx 0, \tilde{\Psi}^\mu_+(p) \approx \sqrt{\frac{2}{3}} \left(\frac{p^\mu}{\tilde{m}}\right) u_-(p) \) and \(\tilde{\Psi}^\mu_+(p) \approx \sqrt{\frac{2}{3}} \left(\frac{p^\mu}{\tilde{m}}\right) u_+(p) \). At the end just two gravitino states shall survive in this approximation, one advantage is that they just depend on spinors and not anymore on the polarization vectors that makes calculations worse.

\section{3. Decay Amplitudes with LSP gravitino/goldstino}

In this section the idea is to present how the SHF works, with the amplitudes at hands the next step is to write the spinors and polarization vectors (for the gravitino) as momentum twistors, then the SHF makes its magic. We do not need to worry about addressing the massive case, because a massive momentum twistor is expressed in terms of two massless one [15, 19]. In this proceeding contribution two decay widths are worked out \((\tilde{\chi}_0 \rightarrow V\tilde{\psi}_0, V = \gamma, Z)\), both within the local supersymmetric extension of the Standard Model with gravitino LSP in the final state and with the neutralino as the NLSP. These calculations have been already evaluated in References [43, 44, 45] using the Feynman approach with the trace technology.

\subsection{3.1. Two-body Neutralino decay \(\tilde{\chi}_0 \rightarrow \tilde{\Psi}^\mu \gamma \)}

We build out the amplitudes for the two-body neutralino decays applying the Feynman rules of Ref. [35] for the spin-3/2 gravitino interactions. We shall start with the decay \(\tilde{\chi}_0 \rightarrow \tilde{\Psi}^\mu \gamma \)
corresponding to the diagram of Figure 1. The $\tilde{\chi}_0(p_1)$ denotes the neutralino with momentum $p_1$ and $\tilde{\Psi}^\mu(p_2)$ represents the spin-3/2 gravitino (in the goldstino approximation will be $\tilde{G}(p_2)$), and $\gamma(p_3)$ denotes the photon with momentum $p_3$. The amplitude reads as follows

$$\mathcal{M}_{\lambda_1, \lambda_2, \lambda_3} = \frac{C_{\chi\gamma}}{4M} \tilde{\Psi}_{\lambda_3} (p_2)(p_{3\nu}[\gamma^\nu, \gamma^\sigma]\gamma^\mu)\epsilon_{\sigma\lambda_3}(p_3)u_{\lambda_1}(p_1). \quad (12)$$

where $C_{\chi\gamma} = U_{i1} \cos \theta_W + U_{i2} \sin \theta_W$, $U_{ij}$ are the mixing matrices that diagonalize the neutralino factor, $M$ is the Plank mass, then $\lambda_1$, $\lambda_2$ and $\lambda_3$ are the helicity labels corresponding to the neutralino, gravitino and photon. After applying the SHF to the amplitude of Eq. (12), one notice that there are in principle 16 HAs, but 14 of them vanish. The SHF helps to identify the spurious quantities from the beginning, making the process to find physical observables expeditious. The nonzero HAs are shown in the Table 1.

| $\lambda_1$, $\lambda_2$, $\lambda_3$ | $\mathcal{M}_{\lambda_1, \lambda_2, \lambda_3}^{\pm/0}$ |
|----------------------------------|----------------------------------|
| $-, +++, +$                      | $\frac{C_{\chi\gamma}[r_{2q}]}{M(r_{1q})}m_{\tilde{\chi}_0}$ |
| $-, --, -$                       | $\frac{C_{\chi\gamma}[r_{2q}]}{\sqrt{Mm_{\tilde{\chi}_0}}}|r_{2q}|[r_{2q}]$ |

Table 1. Helicity Amplitudes for the two-body neutralino decay $\tilde{\chi}_0 \rightarrow \tilde{\Psi}^\mu \gamma$ with LSP gravitino in the final state. We are using $s_{ij} = -(p_i + p_j)^2$ for the Mandelstam variable. There are two more HAs ($+, --, +$ and $-, --, -$), but they are the complex conjugate of the two shown in this table, the same criterion shall apply to the next tables.

The squared and averaged amplitude takes the form

$$\langle |\mathcal{M}|^2 \rangle = \frac{C_{\chi\gamma}^2}{2M^2} \left( |\mathcal{M}_{-,+++}|^2 + |\mathcal{M}_{-+-}|^2 + |\mathcal{M}_{+-,--}|^2 + |\mathcal{M}_{++,+}|^2 \right) \quad (13)$$

$$= \frac{C_{\chi\gamma}^2}{2M^2} \left( 2 \frac{s_{12}^2 m_{\tilde{\chi}_0}^2}{s_{12}} + 2 \frac{s_{12}^2 m_{\tilde{\chi}_0}^2}{s_{12}} s_{12} \right) \quad (14)$$

$$= \frac{C_{\chi\gamma}^2}{M^2} \left( \frac{(m_{\tilde{\chi}_0}^2 - \tilde{m}^2)^2}{3\tilde{m}^2} (3\tilde{m}^2 + m_{\tilde{\chi}_0}^2) \right) \quad (15)$$

$$= \frac{C_{\chi\gamma}^2 m_{\tilde{\chi}_0}^6}{M^2} \left( 1 - \tilde{m}^2 m_{\tilde{\chi}_0}^2 \right)^2 \left( \frac{1}{3} + \frac{\tilde{m}^2}{m_{\tilde{\chi}_0}^2} \right). \quad (16)$$

Figure 1. Feynman diagram for the 2-body neutralino decay.

The result of Eq. (16) conduces to the right decay width found in Ref. [44, 45], but our calculations result considerably simpler applying the SHF (massive).

We shall considerate now the approximation of the spin-3/2 gravitino to spin-1/2 goldstino, this
is due to the high energy equivalence theorem, physically the approximation is valid when the split mass between fermions and bosons is larger than the gravitino mass. Some work has been done in order to compare observables (lifetime $\tau$) with gravitino and goldstino in the final state [24], basically comparing the squared amplitudes. It shall be interesting to exploit the power of the SHF for comparing the HAs with gravitino and goldstino for a given process or reaction.

The amplitude for the decay $\tilde{\chi}_0 \rightarrow \widetilde{G} \gamma$ with goldstino in the final state is the following

$$M_{\lambda_1, \lambda_2, \lambda_3} = \frac{C_{\gamma} m_{\tilde{\chi}_0}}{2\sqrt{6} \tilde{m}} \bar{\psi}_{\lambda_2}(p_2)(p_{30} [\gamma^\nu, \gamma^\sigma]) \epsilon_{\sigma \lambda_3}(p_3) u_{\lambda_1}(p_1), \quad (17)$$

the kinematics of the process remains the same in the approximation to goldstino, but now the transversal degrees of freedom vanish. The nonzero HAs are shown in the Table 2

| $\lambda_1$, $\lambda_2$, $\lambda_3$ | $M^{1/2}_{\lambda_1, \lambda_2, \lambda_3}$ |
|-------------------------------|----------------------------------|
| $-, -, -$                     | $\frac{C_{\gamma} \gamma_{2q_2}}{\sqrt{3} \tilde{m} (r_2 q_2)} \langle r_2 q_2 | r_2 r_1 \rangle$ |

Table 2. Helicity Amplitudes for the two-body neutralino decay $\tilde{\chi}_0 \rightarrow \widetilde{G} \gamma$ with LSP goldstino in the final state.

In Table 1 we show two HAs, one of them includes the transversal d.o.f. ($-, +, +, +$) corresponding to the full gravitino contribution, but Table 2 shows just one HA, $M^{1/2}_{-, -, -}$, that is in fact identically to the HA $M^{3/2}_{-, -, -}$.

The squared and averaged amplitude takes the following form

$$\langle |M|^2 \rangle = \frac{C_{\gamma}^2}{2M^2} (|M_{-, -, -}|^2 + |M_{+, +, +}|^2) \quad (18)$$

$$= \frac{C_{\gamma}^2 s_{2q_2}^2 r_{2q_2}}{3M^2 \tilde{m}^2} s_{r_2 r_1}^2 \quad (19)$$

$$= \frac{C_{\gamma}^2 (m_{\tilde{\chi}_0}^2 - \tilde{m}^2)^2 m_{\tilde{\chi}_0}^2}{3M^2 \tilde{m}^2}, \quad (20)$$

this result can be obtained directly from Eq. (16) in the limit that the gravitino mass ($\tilde{m}$) is small compared to the neutralino mass ($m_{\tilde{\chi}_0}$).

### 3.2. 2-body neutralino decay $\tilde{\chi}_0 \rightarrow \tilde{\Psi}^\mu Z$

In the case that the vector boson is massive (but neutral) the decay becomes $\tilde{\chi}_0(p_1) \rightarrow \tilde{\Psi}^\mu(p_2) Z(p_3)$. Some technical complications appear from having the $Z$ boson in the final state instead of the photon. The amplitude for this process with the spin-3/2 gravitino in the final state is as follows

$$M_{\lambda_1, \lambda_2, \lambda_3} = \frac{C_{\gamma} Z}{4M} \bar{\psi}_{\mu \lambda_2}(p_2)(p_{30}^{(u)}[\gamma_\rho, \gamma_\sigma] \gamma^\mu) \epsilon_{\rho \lambda_3}(p_3) u_{\lambda_1}(p_1), \quad (21)$$

where $C_{\chi Z} = -U_{i1} \sin \theta_W + U_{i2} \cos \theta_W$, and everything else remains equal to the decay with a photon in final state ($\tilde{\chi}_0 \rightarrow \tilde{\Psi}^\mu \gamma$), but now the helicity label corresponding to the $Z$ boson ($\lambda_3$) could be $+,-,0$.

From the 24 possible HAs resulting of Eq. (21), the nonzero are shown in the Table 3.
Table 3. Helicity Amplitudes for the two-body Neutralino decay $\tilde{\chi}_0 \rightarrow \tilde{\Psi}^\mu Z$ with LSP gravitino in the final state.

| $\lambda_1, \lambda_2, \lambda_3$ | $M_{\lambda_1, \lambda_2, \lambda_3}^{1/2}$ |
|-------------------------------|---------------------------------|
| $+, -, +$                     | $\frac{C_{\lambda_1, \lambda_2, \lambda_3}}{\sqrt{3M_{\bar{m}_{s_{2q_2}}}|Z_1| |r_1 r_2|}} (\tilde{m}^3 M^2 s_{2q_1 r_1} + s_{2q_2} m_{\tilde{\chi}_0})$ |
| $-, -, 0$                     | $\frac{C_{\lambda_1, \lambda_2, \lambda_3}}{\sqrt{3M_{\bar{m}_{s_{2q_2}}}|Z_1| |r_1 r_2|}} (\tilde{m} s_{2q_1 r_1} + s_{2q_2} m_{\tilde{\chi}_0})$ |
| $-, -, +$                     | $-\frac{C_{\lambda_1, \lambda_2, \lambda_3}}{M (r_1 |q_1 r_2|) |r_1 r_2|} (\tilde{m} M^2 m_{\tilde{\chi}_0} + s_{2q_2} s_{2q_1 r_1})$ |

The squared and averaged amplitude is as follows

$$\langle |M|^2 \rangle = \frac{C_{\lambda_1, \lambda_2, \lambda_3}^2}{2M^2} \left( |M_{+, -, +}|^2 + |M_{-, +, -}|^2 + |M_{-, -, 0}|^2 + |M_{+, +, 0}|^2 + |M_{-, -, +}|^2 + |M_{+, +, -}|^2 \right)$$

$$= \frac{C_{\lambda_1, \lambda_2, \lambda_3}^2}{M^2} \left[ \frac{1}{3\tilde{m}^2 s_{2q_2} s_{1q_2}} (\tilde{m}^3 M^2 s_{2q_1 r_1} + s_{2q_2} m_{\tilde{\chi}_0})^2 + \frac{2M^2}{3s_{2q_1 r_1}} (s_{2q_2} m_{\tilde{\chi}_0} + \tilde{m} s_{2q_1 r_1})^2 + \frac{1}{s_{2q_1 r_1} s_{2q_2}} (\tilde{m} m_{\tilde{\chi}_0} M^2 + s_{2q_1 r_1} s_{2q_2})^2 \right]$$

$$= \frac{C_{\lambda_1, \lambda_2, \lambda_3}^2}{M^2} \left[ \left( 1 - \frac{m_{\tilde{\chi}_0}}{m_{\tilde{\chi}_0}} \right)^2 \left( 1 + \frac{m_{\tilde{\chi}_0}}{m_{\tilde{\chi}_0}} \right) - \frac{m_{\tilde{\chi}_0}^2}{m_{\tilde{\chi}_0}^2} \left( 1 - \frac{m_{\tilde{\chi}_0}^2}{m_{\tilde{\chi}_0}^2} \right) - \frac{m_{\tilde{\chi}_0}^2}{3m_{\tilde{\chi}_0}^2} \left( 4 - \frac{m_{\tilde{\chi}_0}^2}{3m_{\tilde{\chi}_0}^2} \right) \right]$$

this result conduces to the known decay width found in [44, 45].

We repeat the process of the first example and approximate the gravitino to goldstino, then we compute the HAs.

$$M_{\lambda_1, \lambda_2, \lambda_3} = \frac{C_{\lambda_1, \lambda_2, \lambda_3} \bar{m}_{\tilde{\chi}_0} (p_2) (p_3[s_{\gamma\rho}, \gamma\sigma]) \epsilon_\lambda^\mu (p_3) u_\lambda (p_1),$$

the nonzero HA are shown in the Table 4

Table 4. Helicity Amplitudes for the two-body Neutralino decay $\chi_0 \rightarrow Z \tilde{G}$ with goldstino in the final state.

| $\lambda_1, \lambda_2, \lambda_3$ | $M_{\lambda_1, \lambda_2, \lambda_3}^{1/2}$ |
|-------------------------------|---------------------------------|
| $+, -, +$                     | $\frac{C_{\lambda_1, \lambda_2, \lambda_3}}{\sqrt{3M_{\bar{m}_{s_{2q_2}}}|Z_1| |r_1 r_2|}} (\tilde{m} m_{\tilde{\chi}_0} M^2 + s_{2q_2} s_{2q_1 r_1})$ |
| $-, -, 0$                     | $\frac{C_{\lambda_1, \lambda_2, \lambda_3}}{\sqrt{3M_{\bar{m}_{s_{2q_2}}}|Z_1| |r_1 r_2|}} (\tilde{m} s_{2q_1 r_1} + s_{2q_2} m_{\tilde{\chi}_0})$ |

We can notice in Table 3 that for the helicity labels $\lambda_1 = -, \lambda_2 = -$ and $\lambda_3 = 0$ the HA for the gravitino is up to a factor of $\sqrt{2}$ the same HA with goldstino, this is $M_{-, -, 0}^{3/2} = \sqrt{2} M_{-, -, 0}^{1/2}$. 


Squaring and averaging the HAs of Table 4, these take the following form

$$\langle |M|^2 \rangle = \left( |M_{+,+,+}|^2 + |M_{-,+,+}|^2 + |M_{-,+,0}|^2 + |M_{+,+,0}|^2 \right)$$

$$= \frac{C_P^2}{\lambda^2 m^2_{\chi_0}} \left( M_Z^2 (m_G s_{q_2} + s_{r_2 q_2} m_{\chi_0})^2 + 2 (m m_{\chi_0} M_Z^2 + s_{r_2 q_2} s_{q_2 r_1})^2 \right)$$

$$= \frac{2 C_P^2}{\lambda^2 m^2_{\chi_0}} \left[ \left( 1 - \frac{m^2}{m_{\chi_0}^2} \right)^2 - \frac{M_Z^2}{m_{\chi_0}^2} \left( \frac{1}{2} \left( 1 - \frac{m}{m_{\chi_0}} + \frac{m^2}{m_{\chi_0}^2} \right) + \frac{M_Z^2}{2 m_{\chi_0}^2} \right) \right]$$

In the traditional Feynman approach with trace technology, the calculations of the squared and averaged amplitudes with goldstinos are considerably simple, unlike amplitudes involving gravitinos. However, when the SHF is implemented the calculations with gravitino are still simple, and the helicity method is able to compute the decay widths without any approximation, as we have shown.

4. Conclusions

In this proceeding, we have presented how the spinor helicity formalism is well suitable to evaluate decay amplitudes, even in the massive case. Two examples have been recalculated with the SHF ($\chi_0 \to V \tilde{\Psi}^\mu$ with $V = \gamma, Z$). Furthermore, we briefly discussed the high energy equivalence theorem between gravitinos and goldstinos. It is possible in some cases to identify how the helicity amplitudes with goldstinos appear as a helicity amplitudes with gravitinos (for the longitudinal d.o.f.), this characteristic is difficult to identify in the traditional approach. We shall study and discuss more relationships between gravitino and goldstino amplitudes for another NLSP candidates in a future work to appear soon in a peer review journal.

5. Acknowledgements

The author wish to thank Lorenzo Diaz-Cruz for useful discussion, Roberto Garcia and Jose Zelaya for reading the manuscript. This work has been supported by CONACYT and was partly supported by DPC-SMF.

References

[1] L. Dixon, arXiv:hep-ph/9601359.
[2] S.J. Parke and T. Taylor, Phys. Lett. 157 B:81 (1985); Z. Kunszt, Nucl. Phys. B 271:333 (1986); M. Mangano and S. Parke, Phys. Rep. 200:301 (1991).
[3] M. Srednicki, Quantum Field Theory, Cambridge University Press, 2007.
[4] M. Peskin, http://arxiv.org/pdf/1101.2414.pdf.
[5] M. D. Schwarz, Quantum Field Theory and the Standard Model, Cambridge University Press, 2014.
[6] J. M. Henn, J. C. Plefka, Scattering Amplitudes in Gauge Theories, in: Lecture Notes in Physics, vol. 883, Springer, 2014.
[7] H. Elvang and Y. Huang, Scattering Amplitudes in Gauge Theory and Gravity (Cambridge University Press, 2015), arXiv:1308.1697v2 [hep-th].
[8] E. Witten, Commun. Math. Phys. 252 189-258 (2004), arXiv:hep-th/0312171.
[9] R. Britto, F. Cachazo, B. Feng, E. Witten, Phys. Rev. Lett. 94 (2005) 181602. hep-th/0501052.
[10] L.J. Mason, D. Skinner, Phys. Lett. B 636 (2006) 60. hep-th/0510262.
[11] Z. Bern, J. J. M. Carrasco, H. Johansson, Phys. Rev. D 78 (2008) 085011. arXiv:0805.3993.
[12] N. Arkani-Hamed, J. Trnka, J. High Energy Phys. 10 (2014) 030. arXiv:1312.2007.
[13] F. Cachazo, S. He, E.Y. Yuan, Phys. Rev. D 90 (2014) 065001. arXiv:1406.6575.
[14] K. Ozeren and W. Stirling, Eur. Phys. J. C 48 159 (2006), arXiv:hep-ph/0603071.
[15] J. Kuczmarski, arXiv:1406.5612 [hep-ph].
[16] D. A. Kosower, Phys. Rev. D 71, 045007 (2005), [arXiv:hep-th/0406175].
[17] C. Schwinn and S. Weinzierl, JHEP 0704 (2007) 072, arXiv:hep-ph/0703021 [HEP-PH].
[18] J. Lorenzo Diaz-Cruz, Bryan O. Larios, O. Meza-Aldama, J. Phys. Conf. Ser. 761 (2016) no.1, 012012, arXiv:1608.04129 [hep-ph].
[19] J. Lorenzo Diaz-Cruz, Bryan O. Larios, arXiv:1612.04331 [hep-ph].
[20] J. L. Feng, A. Rajaraman and F. Takayama, Phys. Rev. Lett. 91 (2003) 011302, arXiv:hep-ph/0302215; Phys. Rev. D 68 (2003) 063504, arXiv:hep-ph/0306024.
[21] J. L. Feng, S. Su and F. Takayama, Phys. Rev. D 70 (2004) 075019, arXiv:hep-ph/0404231.
[22] J. R. Ellis, K. A. Olive, Y. Santoso and V. C. Spanos, Phys. Lett. B 588 (2004) 7, arXiv:hep-ph/0312262.
[23] J. L. Diaz-Cruz, John Ellis, Keith A. Olive, Yudi Santoso, JHEP 0705:003, 2007, arXiv:hep-ph/0701229v1.
[24] J. Lorenzo Diaz-Cruz, Bryan O. Larios, Eur. Phys. J. C76 (2016) no.3, 157, arXiv:1510.01447v2 [hep-ph].
[25] F. D. Steffen, arXiv:hep-ph/0711.1240.
[26] M. Johansen, J. Edsj, S. Hellman, J. Milstead , JHEP 1008 1-27 (2010).
[27] G. D. Kribs, A. Martin, and T. S. Roy, JHEP 0901 (2009) 023, arXiv:hep-ph/0807.4936.
[28] J. Heisig, J. Heising, JCAP 04 (2014) 023, arXiv:1310.6352.
[29] J. R. Ellis, A. R. Raklev and O. K. Oye, JHEP 0610, 061 (2006), arXiv:hep-ph/0607261.
[30] K. Hamaguchi, Y. Kuno, T. Nakaya and M. M. Nojiri, Phys. Rev. D 70 (2004) 115007, arXiv:hep-ph/0409248.
[31] R. H. Cyburt, J. R. Ellis, B. D. Fields and K. A. Olive, Phys. Rev. D 67 (2003) 103521, arXiv:astro-ph/0211258; J. R. Ellis, K. A. Olive and E. Vangioni, Phys. Lett. B 619 (2005) 30, arXiv:astro-ph/0503023.
[32] M. Kawasaki, K. Kohri and T. Moroi, Phys. Lett. B 625 (2005) 7, arXiv:astro-ph/0402490; Phys. Rev. D 71 (2005) 083502, arXiv:astro-ph/0408426.
[33] K. Kohri and Y. Santoso, Phys. Rev. D 79, 043514 (2009), arXiv:0811.1119 [hep-ph].
[34] W. Rarita and J. Schwinger, Phys. Rev. 60, 61 (1941).
[35] Takeo Moroi, arXiv:hep-ph/9503210v1.
[36] P.R. Auvil and J.J. Brehm, Phys. Rev. 145 (1966) 1152.
[37] R. Casalbuoni, S. De Curtis, D. Dominici, F. Feruglio, and R. Gatto, Physics Letters B 215, 313-316, 1988.
[38] P. Fayet, Phys.Lett. 70B (1977) 461.
[39] P. Fayet, Phys. Lett. B. 175 (1986) 471.
[40] P. Fayet, Phys. Lett. B. 84 (1979) 421.
[41] P. Fayet, Phys. Lett. B. 86 (1979) 272.
[42] P. Fayet, Conference Proc. LPTENS-81-9 (1981) 347.
[43] Laura Covi, Jasper Hasenkamp, Stefan Pokorski, Jonathan Roberts, JHEP 0911:003, 2009, arXiv:0908.3399v1 [hep-ph].
[44] John Ellis, Keith A. Olive, Yudi Santoso, Vassilis Spanos, Phys. Lett. B588:7-16, 2004, arXiv:hep-ph/0312262v4.
[45] Jonathan L. Feng, Shufang Su, Fumihiro Takayama, Phys. Rev. D70:075019, 2004, arXiv:hep-ph/0404231v2.