Disentangling the Unparticles with polarized beams at $e^+e^-$ colliders

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Recently proposed idea of unparticles arising due to a scale invariant sector in the theory can give rise to effective operators with different Lorentz structures. We show that by using the different polarization options at the future linear $e^+e^-$ colliders, the nature of these effective operators can be easily understood. The unique feature of a complex phase in the propagator of the unparticle can also be understood uniquely for the different spins by exploiting the initial beam polarizations at the International Linear Collider (ILC).

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The recently proposed idea of the consequences of having a scale invariant sector in a theory [1, 2], which decouples from the low energy theory, has drawn enormous attention due to its rich phenomenology. Such a sector of scale invariant physics existing at a much higher scale can co-exist and still remain sterile to interactions at very low energies to hide its existence. Motivated by the Banks-Zaks theory [3], it was argued however, that a scale invariant sector with a nontrivial infrared fixed-point leads to peculiar low energy behavior in an effective field theory approach, with non-matter stuff coined as “unparticles” ($U$). The production of these unparticles, both on-shell and off-shell, can lead to interesting signatures at future accelerator experiments. Real production of the unparticles will lead to missing energy signatures. Another peculiar feature of the unparticles was identified as a complex “phase” originating in the propagators of the unparticles [2, 4] for time like momenta. These phases can be a crucial test in isolating signatures for unparticles at future experiments, and features related to the phases were shown to be quite interesting [2, 4, 5].

A lot of work already exists in the literature exploring the various possibilities for unparticles and their signatures beyond the Standard Model (SM) physics [6]. We show that this feature can be very easily captured by the use of transverse beam polarizations at the ILC. It has been shown that transverse beam polarizations can be a very useful tool at ILC to study CP-odd observables and effects of beyond the SM physics [8, 9, 10, 11, 12].

We choose the following notation for introducing the beam polarizations

$$
\sum_{s_1} \bar{u}(k_1, s_1) u(k_1, s_1) = \frac{1}{2} (1 + P_L \gamma_5 + \gamma_5 P_T \gamma_1) \gamma_1
$$

$$
\sum_{s_2} \bar{v}(k_2, s_2) v(k_2, s_2) = \frac{1}{2} (1 - P_L \gamma_5 + \gamma_5 P_T \gamma_2) \gamma_2
$$

where $t_{1,2}$ are the transverse beam polarization 4-vectors of the electron and positron beams, respectively. In the above equation, $P_L$ and $P_T$ represent the degree of the longitudinal and transverse beam polarizations. For our analysis, we chose $|P_L^t| = 0.8$ and $|P_T^t| = 0.6$. For the transverse beam polarization 4-vectors we assume $t_1^t = (0, 1, 0, 0) = -t_2^t$ (the results still hold for other choices). The scattering process we consider is mediated by the $\gamma$ and Z-boson propagators in the SM, while the “unparticle” effects will show up due to the s-channel contribution coming from the effective unparticle operators $\mathcal{O}_t$ respecting the SM gauge symmetry. We focus on three different operators $\gamma_1$ the scalar, vector and tensor operators given respectively by

$$
\lambda_0 \frac{1}{\Lambda_d} \bar{f} f \mathcal{O}_t
$$

$$
\lambda_1 \frac{1}{\Lambda_d} \bar{f} \mu (1 + \gamma_5) f \mathcal{O}_t^\mu
$$
The trivial scale invariant sector below the scale $\Lambda$ was found to have the following general form $\lambda_1$ and $\lambda_2$ are the dimensionless effective couplings obtained after the imposition of matching conditions on the nontrivial scale invariant sector below the scale $\Lambda_U$. The $d_{it}$ stands for the scaling dimension of the unparticle operator $O_{it}$. For simplicity, we have restricted ourselves to $\lambda_i = 1$ for all $i = 0, 1, 2$.

We now take up each of the different cases separately. One of the very distinct features that isolate the unparticles is the complex phase that appears for the s-channel propagator. The virtual propagator for the unparticles was found to have the following general form $\psi_{it}$ and $O_{it}$ represent the spin structures. The detailed expressions for them can be found in Ref. [2]. The interesting thing to note is that for $P^2 > 0$, the complex function $(- P^2)^{d_{it}-2}$ needs a branch cut. This results in a complex phase to be associated for an s-channel exchange of unparticles. This gives interesting effects on the cross-sections for physical process due to its interference effects. In this work we take up this feature of the unparticle propagator and try to analyze the effects it can have on the azimuthal distribution of the final state particles.

The general expression for the differential cross-section of scattering process $e^+ e^- \rightarrow f \bar{f}$ can be written down as

$$
\frac{d\sigma}{d\Omega} = \frac{1}{16\pi^2} \frac{F}{2s^3} |M_{SM+U}|^2
$$

where $F = \lambda_{1/2}^i s, m_f, m_f$ and $d\Omega = \sin\theta d\theta d\phi$. The spin-averaged amplitudes for the SM and the different unparticles contributing in the virtual exchange is given by $|M_{SM+U}|^2$ and is listed individually in the appendix. We give the formulae for both longitudinal and transverse beam polarizations. Throughout our analysis, our choice for the free parameters of the unparticle sector are as follows

$$
\Lambda_U = 10 \text{ TeV}, \quad d_{it} = 1.5 \text{ (spin – 0 and spin – 1)}
\Lambda_U = 1 \text{ TeV}, \quad d_{it} = 1.3 \text{ (spin – 2)}
$$

![FIG. 1: The normalized differential cross-section plotted against the scattering angle $\theta$ for the final state $t$ quark in the scattering process $e^+ e^- \rightarrow t \bar{t}$](image1)

![FIG. 2: The normalized differential cross-section plotted against the scattering angle $\theta$ for the final state $b$ quark in the scattering process $e^+ e^- \rightarrow b \bar{b}$](image2)

![FIG. 3: The normalized differential cross-section plotted against the scattering angle $\theta$ for the final state $t$ quark in the scattering process $e^+ e^- \rightarrow t \bar{t}$](image3)
The above choice gives a significant number of new physics events for the different spin unparticles for an integrated luminosity in the range of $L = 10 - 100 \, fb^{-1}$.

In Figs.1 and 2, we plot the normalized differential cross-sections against the scattering angle for the final state fermions. One can use longitudinally polarized beams which help in suppressing the SM background in the analysis for unparticles with different spin options. Since our results are quite general, provided we have significant cross-section for new physics, we just consider one case where the initial beam polarization is $P_L = 0.8$ and $P_T = 0.6$. Thus we have preferred to show the normalized distributions only. The notation we follow for the cross-sections plotted for different spin cases of the unparticles is simply the excess over the SM cross-section. One can use longitudinally polarized beams which help in suppressing the SM background in the analysis for unparticles with different spin options.

Once integrated over the azimuthal angle, the transverse polarization does not affect the total cross-section for the SM or with the spin-1 and spin-2 exchange of the unparticles. The initial beams are transversely polarized with $P_T = -0.8$ and $P_T^e = -0.6$.

FIG. 4: The normalized differential cross-section plotted against the azimuthal angle $\phi$ for the final state $\mu$ in the scattering process $e^+e^- \rightarrow \mu^\pm\mu^\mp$. The initial beams are transversely polarized with $P_T^e = -0.8$ and $P_T^e = -0.6$.

FIG. 5: The normalized differential cross-section plotted against the azimuthal angle $\phi$ for the final state $b$ quark in the scattering process $e^+e^- \rightarrow bb$. The initial beams are transversely polarized with $P_T^e = -0.8$ and $P_T^e = -0.6$.

In Figs.3 and 4, we plot the normalized differential cross-section against the azimuthal angle for the transversely polarized initial beams, for the three different final states, particularly for the top quarks in the final state. However one needs to keep in mind that the heavy top quark will decay and the efficacy of this distinction would clearly be on how well the parent particle gets reconstructed from its decay products. Note here that although the complex phase in the s-channel has a nontrivial effect on the production cross-section as shown in earlier works [2, 4, 5], its effects on the scattering angle distribution are minimal.

At the next generation linear collider which would run for one or two fixed center-of-mass energies, it would be very difficult to see the effects of the complex phase on the cross-section. So we should try to identify a variable which is susceptible to the complex phase. The most simple variable which can have a nontrivial dependence is the azimuthal angle ($\phi$). The $\phi$-dependent terms in the spin-averaged matrix amplitude square for the scattering process $e^+e^- \rightarrow f\bar{f}$ can be accessible if the initial beams are transversely polarized [10]. The azimuthal angle is defined by the directions of the $e^\pm$ transverse polarization and the plane of the momenta of the outgoing fermions in the $e^+e^- \rightarrow f\bar{f}$ process. Writing the various matrix amplitudes for the SM and the unparticle exchanges in the s-channel with transversely polarized beams, one realizes that the complex phase gives a non vanishing imaginary component in the interference terms which are proportional to the azimuthal angle and contributes by interfering with the imaginary part of the $Z$-boson propagator.

Once integrated over the azimuthal angle, the transverse polarizations do not affect the total cross-section for the SM or with the spin-1 and spin-2 exchange of the unparticles [11]. But one finds an imprint of the transversely polarized beams in the azimuthal angle distribution of the final state fermions. We try to exploit this fact in exploring the effect of the complex phase due to the s-channel unparticle propagator in the simplest scattering process at the linear collider.

In Figs.4 and 5, we plot the normalized differential cross-section against the azimuthal angle for the transversely polarized initial beams, for the three different final states viz. $\mu^+\mu^-$, $bb$ and $tt$ respectively. One finds
that for the spin-0 unparticle exchange, the azimuthal angle dependence is proportional to the mass of the final state fermions. Hence in Fig. 4 we find that for the scalar exchange there is no azimuthal dependence, as the interference terms with the photon and the Z-boson vanish simultaneously for vanishing fermion mass. However, the other unparticle exchanges do show a non trivial dependence. The most interesting thing to notice is the case for the spin-1 exchange of the unparticle. It shows a clear phase difference of $\pi/4$ with the SM distribution. This is purely due to the complex phase, as the choice of $d_U$ is seen to control its distribution. The spin-2 exchange is seen to behave like a pure photon exchange distribution, with no $\gamma_5$ dependence. To explore the role of mass of the final state fermions in the distributions, we plot similar distributions for massive fermions in the final state. The choice for the input scale $\Lambda_U$ and the scaling dimension $d_U$ remain the same as before. In Fig. 4 we show the azimuthal dependence for the final state b-quarks. One marked difference is in the SM diagrams. To really highlight the effect of mass on the azimuthal distribution, we consider the top quark in the final state and show the corresponding distribution in Fig. 6. The most striking change is for the scalar unparticle exchange, which was expected since the interference terms were proportional to the mass of the final state fermions. But it is also worth noting the effect of the mass on the distribution for the spin-2 exchange of unparticles.

In order to show the effect of the complex phase associated with the s-channel unparticle propagator on the azimuthal distribution, we consider the process $e^+e^- \rightarrow t\bar{t}$ and set the complex phase to zero. We plot the corresponding distribution in Fig. 7. As expected the phase difference between the spin-1 unparticle exchange and SM vanishes, and even the substantially modified curve for the scalar unparticle exchange suggests the non-trivial dependence, which the complex phase has on its distribution. We however point out that when setting the complex phase to zero, the scaling dimension $d_U$ and the unparticle scale $\Lambda_U$ has been taken as

$$\Lambda_U = 1 \text{ TeV}, \quad d_U = 1.4 \ (spin - 0)$$

$$\Lambda_U = 2 \text{ TeV}, \quad d_U = 1.4 \ (spin - 1)$$

$$\Lambda_U = 1 \text{ TeV}, \quad d_U = 1.15 \ (spin - 2)$$

so as to have enough significant cross-sections for the different cases.

As pointed out earlier we find that for the spin-1 unparticle exchange, one can get a very strong estimate of the scaling dimension $d_U$ in an independent way. This can be done by just looking at the azimuthal distribution of the final state fermion and compare it with the pure SM distribution. To highlight this, we plot the normalised distributions for different values of $d_U$ for a spin-1 unparticle exchange.
FIG. 8: The normalized differential cross-section plotted against the azimuthal angle $\phi$ for the final state $\mu$ in the scattering process $e^+e^- \rightarrow \mu^+\mu^-$. We consider the spin-1 unparticle contribution in the s-channel and compare with the SM distribution. The initial beams are transversely polarized with $P_T^e = -0.8$ and $P_T^\mu = -0.6$.

exchange in Fig. 8. One finds that for $\Delta d_U = 0.25$, the corresponding phase difference with the SM is $\pi/8$. This gives a very clean measure of the scaling dimension $d_U$ by just studying the azimuthal distribution of the fermion in the final state.

Thus we find that a simple analysis of the scattering angle and azimuthal dependence of the normalized differential cross-section for the simple minded scattering process $e^+e^- \rightarrow f\bar{f}$ at linear colliders with polarized beams can very effectively identify the effects of unparticle exchanges. By analyzing the above distributions at a 500 GeV linear electron-positron collider we find that one can easily exploit the azimuthal distribution to distinguish the spin of the exchanged unparticle by using transversely polarized initial beams, complemented by the scattering angle distributions. We must point out that one can have a non-trivial azimuthal dependence in supersymmetric theories which do not conserve R-parity, where one can consider scalar exchange in the propagator [13]. However, the contributions to the process $e^+e^- \rightarrow f\bar{f}$ would be through the exchange of squarks and sleptons. The angular distributions would be crucial and very effective in distinguishing the scalar exchanges of the supersymmetric particles from that of the spin-1 or spin-2 unparticle exchanges. To distinguish the effect of scalar unparticles from supersymmetric contributions, the azimuthal distribution will still be a useful tool as the distributions with the complex phase differs from the generic scalar exchanges, as shown in Fig. 7.

We also find an interesting fact that the spin-1 unparticle exchange shows a direct sensitivity to the choice the scaling dimension $d_U$ in the form of the phase difference with the SM distribution in $\phi$, and hence can provide for a direct measure of the scaling dimension, irrespective of the scale $\Lambda_U$. This is an extraordinary result, bearing in mind that the complex phase associated with the s-channel exchange is a unique property of the unparticle theory. One can also consider different asymmetries using the dependence of the azimuthal angle to explore the effect of the complex phase in the unparticle propagator with time-like momenta. Thus we see that with the different nature of polarizations at the ILC, the unparticle sector of nature, if it exists and is observable cannot escape our attention and will be found.

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Appendix:

The amplitude with longitudinally polarized beams ($P_L^e$, $P_L^\mu$) for the SM and unparticle exchange can be written down as

$$M_{\gamma\gamma}^2 = Q_e^2 Q_{\mu}^2 e^4 \frac{N_f}{s^2} \left(1 - P_L^e P_L^\mu\right) \left[4u^2 + 4su + 2s^2 + 4m_f^4 - 8um_f^2\right]$$

$$M_{ZZ}^2 = Q_e Q_{\mu} e^4 \frac{N_f}{16s^2 c_W^2} \left[\frac{(1 - P_L^e P_L^\mu)e_{A\mu} f_{\alpha A} + (P_L^e - P_L^\mu)e_{A\mu} f_{\alpha A}}{s - M_Z^2 + i\Gamma_Z M_Z}\right] \left[(4u^2 + 4su + 2s^2 + 4m_f^4 - 8um_f^2)\right]$$

$$- Q_e Q_{\mu} e^4 \frac{N_f}{16s^2 c_W^2} \left[\frac{(1 - P_L^e P_L^\mu)e_{A\mu} f_{\alpha A} + (P_L^e - P_L^\mu)e_{A\mu} f_{\alpha A}}{s - M_Z^2 + i\Gamma_Z M_Z}\right] \left[2s^2 + 4su - 4m_f^2\right]$$

$$M_{ZZ}^2 = \frac{e^4}{256s^2 c_W^4} \left[\frac{(1 - P_L^e P_L^\mu)(c_{\alpha A}^2 + e_{\alpha A}^2) + 2(P_L^e - P_L^\mu)e_{A\mu} f_{\alpha A}}{s - M_Z^2 + i\Gamma_Z M_Z^2}\right] \left[(4u^2 + 4su + 2s^2 + 4m_f^4 - 8um_f^2)\right]$$
\[-\frac{e^4}{256s^4c^4_w} N_c^f \left[ 4 \left( 1 - P_L^- P_L^+ \right) \right] \frac{4(1 - P_L^- P_L^+)}{s - M_Z^2 + (i \Gamma_Z M_Z)^2} [ e_V e_A f_V f_A (2s^2 + 4su - 4sm_\gamma^2) + 2sm_\gamma^2 f_A (e_V^2 + e_A^2)] \]

\[-\frac{e^4}{256s^4c^4_w} N_c^f \frac{2(P_L^- - P_L^+)}{s - M_Z^2 + (i \Gamma_Z M_Z)^2} \left[ f_V f_A (e_V^2 + e_A^2)(2s^2 + 4su - 4sm_\gamma^2) \right] \]

\[M_{2 S S}^2 = \frac{\lambda_4^1 |\Delta_F(-s)|^2 N_c^f (1 + P_L^- P_L^+)}{\lambda_d^{4u-4}} (2s^2 - 5sm_\gamma^2), \quad M_{\gamma S}^2 = 0, \quad M_{Z S}^2 = 0 \]

\[M_{2 V V}^2 = \frac{\lambda_4^1 |\Delta_F(-s)|^2}{\lambda_d^{4u-4}} 16N_c^f (1 - P_L^- P_L^+ + P_L^- - P_L^+)(u^2 - 2um_\gamma^2 + m_f^2) \]

\[M_{2 V}^2 = \frac{Q e_f e^2 \lambda_4^2 \Delta_F(-s)}{16s^4 c_{24u-2}} 4N_c^f (1 - P_L^- P_L^+ + P_L^- - P_L^+)(u^2 - 2um_\gamma^2 + 2sm_\gamma^2 + m_f^2) \]

\[M_{2 V T}^2 = \frac{\lambda_4^1 |\Delta_F(-s)|^2}{\lambda_d^{4u-4}} N_c^f (1 - P_L^- P_L^+ + P_L^- - P_L^+)[256u^4 + 512su^3 + 336s^2u^2 + 80s^3u + 8u^4 - 32m_\gamma^2(32u^3 + 40su^23s^2u + s^3) + 16m_\gamma^2(96u^2 + 64su + 5s^3) - 256m_\gamma^2(4u + s - m_\gamma^2)] \]

\[M_{2 T}^2 = \frac{\lambda_4^1 |\Delta_F(-s)|^2}{\lambda_d^{4u-4}} N_c^f (1 - P_L^- P_L^+ + P_L^- - P_L^+)[32u^3 + 48su^2 + 24s^2u + 4s^3 - 8m_\gamma^2(12u^2 + 8su + s^2) + 16m_\gamma^2(6u + s - 2m_\gamma^2)] \]

\[M_{2 T T}^2 = \frac{-e^2 \lambda_4^2 \Delta_F(-s)}{256s^4 c^4_w} N_c^f \left[ (1 - P_L^- P_L^+)(4u^2 + 4su - 8um_\gamma^2 + 4m_\gamma^2) + N_c^f (2s^2 + P_L^- P_L^+)(2sm_\gamma^2 - 2s^2\sin^2 \theta \cos^2 \phi) \right] \]

Similarly the amplitude with transversely polarized beams \((P_T^- , P_T^+\)) for the SM and unparticle exchange can be written down as

\[M_{2 \gamma \gamma}^2 = \frac{Q e_f e^4}{s^2} N_c^f \left[ (1 - P_T^- P_T^+)(4u^2 + 4su - 8um_\gamma^2 + 4m_\gamma^2) + N_c^f (2s^2 + P_T^- P_T^+)(2sm_\gamma^2 - 2s^2\sin^2 \theta \cos^2 \phi) \right] \]

\[M_{2 Z}^2 = \frac{Q e_f e^4}{16s^4 c^4_w} N_c^f \left[ (1 - P_T^- P_T^+)(4u^2 + 4su + 4m_\gamma^2 - 8um_\gamma^2) - e_A f_A (2s^2 + 4su - 4sm_\gamma^2) \right] \]

\[+ 2e_V f_V s^2 + e_V f_V P_T^- P_T^+ (2sm_\gamma^2 - 2s^2\sin^2 \theta \cos^2 \phi) + i e_A f_V P_T^- P_T^+ (2sm_\gamma^2 - 2s^2\sin^2 \theta \cos^2 \phi) \]

\[M_{2 Z Z}^2 = \frac{e^4}{256s^4 c^4_w} N_c^f \left[ (1 - P_T^- P_T^+)(e_V^2 + e_A^2)(f_V^2 + f_A^2)(4u^2 + 4su + 4m_\gamma^2 - 8um_\gamma^2) \right] \]

\[+ \left\{ f_A^2(e_V^2 + e_A^2)(2s^2 - 8sm_\gamma^2) + 2s^2 f_A^2(e_V^2 + e_A^2) - 8e_V e_A f_V f_A (s^2 + 2s^2 - 2sm_\gamma^2) \right\} \]

\[M_{2 S S}^2 = \frac{\lambda_4^1 |\Delta_F(-s)|^2 N_c^f (1 - P_T^- P_T^+)}{2}(2s^2 - 5sm_\gamma^2) \]

\[M_{2 \gamma S}^2 = \frac{Q e f e^2 \lambda_4^2 \Delta_F(-s)}{2s^4 c_{24u-2}} \left[ -4i N_c^f (P_T^- - P_T^+) \sqrt{s - 4m_\gamma^2 sm_\gamma \sin \theta \cos \phi} \right] \]
\[ M_{2S}^2 = \frac{e^2 \lambda_1^2 \Delta_F(-s)}{16(s - M_Z^2 - i\Gamma_Z M_Z) \Lambda_{4d}^{2d_\ell - 2} s_W^2 c_W^2} \left[ 2N_F (P_T^e T - P_T^e T)^{f_V} \sqrt{s - 4m_f^2 \sin \theta (e_A \cos \phi - ie_V \sin \phi) } \right] \]

\[ M_{2V}^2 = \frac{\lambda_1^2 \Delta_F(-s)^2}{16 N_F^2 (u^2 - 2m_f^2 + m_f^2)} \]

\[ M_{2V}^2 = \frac{Q e^2 \lambda_1^2 \Delta_F(-s)}{s_W^2 c_W^2} N_F^2 \left[ (1 - P_T^e T T^e T^e T)^{(4u^2 - 8m_f^2 + 4m_f^4)} + 4s m_f^2 + 2s P_T^e T T^e T^e T (4u - (s - m_f^2) \sin^2 \theta \cos \phi e_i^f ) \right] \]

\[ M_{2T}^2 = \frac{\lambda_1^2 \Delta_F(-s)^2}{16 \Lambda_{4d}^{2d_\ell}} N_F^2 \left[ (1 - P_T^e T T^e T^e T)^{(2u^3 + 3su^2 - 6u^2 m_f^4 - 4su^2 m_f^4 + 6m_f^4 + 4m_f^4)} \right] \]

\[ M_{2T}^2 = \frac{-e^2 \lambda_1^2 \Delta_F(-s)}{16 \Lambda_{4d}^{2d_\ell}} N_F^2 \left[ 16(1 - P_T^e T T^e T^e T)^{(2u^3 + 3su^2 - 2u^2 m_f^4 - 4su^2 m_f^4 + 6m_f^4 + 4m_f^4)} \right] \]

In the above formulae for the matrix amplitude squares, \( m_f \) stands for the mass of the final state fermions, \( N_F^2 \) stands for the color factor which equals 1 for leptons and 3 for quarks. The \( s \) and \( u \) represent the usual Mandelstam variables while \( M_Z \) and \( \Gamma_Z \) stand for the mass and total width of the Z-boson. The variables \( \theta \) and \( \phi \) are the scattering angle and the azimuthal angles for the final state fermions respectively, while \( e_V = 2T^\ell - 4Q s_W^2 \), \( e_A = -2T^\ell \), and \( f_V = 2T^\ell - 4Q s_W^2 \). \( f_A = -2T^\ell \). We have absorbed the spin structures while calculating the amplitude squares and hence

\[ \Delta(-s) = Z_{d\ell} s_{d\ell}^{-2} e^{-id\ell \tau} \]

The spin averaged matrix amplitude square is given by the sum written in the form

\[ \sum_{i,j} M_{ij}^2 + 2Re \left( \sum_{i,j} M_{ij}^2 \right) \]

where the \( i, j = \gamma, Z, S(V, T) \) stand for the particles exchanged in the s-channel. The \( S, V \) and \( T \) stand for the scalar, vector and tensor unparticle respectively.

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