Collider signatures of Hylogenesis

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The concept of particle dark matter (DM) is the great mystery in physics [1]. In spite of the fact that DM constitutes a sizable fraction $\Omega_{DM}$ of the present Universe energy density and has amount almost five times larger than that of baryons $\Omega_B$ [2] particle physics experiments have so far failed to obtain any convincing signal from dark matter sector.

Existence of dark matter is one of the main phenomenological reasons to believe that Standard Model of particle physics should be replaced by a new theory. One of the explanations for DM is existence of a cosmologically stable weakly interacting massive particles (WIMPs). Being thermalized in the primordial plasma at the Hot stage of the Universe evolution WIMPs later become nonrelativistic and freeze-out through annihilation into SM particles. Then their relic abundance is determined by the corresponding annihilation cross section. Such particles can naturally appear in some SM extensions like supersymmetry or extra dimensions. However, in such theories in general the fact that $\Omega_{DM}$ is of the same order as $\Omega_B$ is accepted as a coincidence (“the WIMP miracle”).

However, the latter coincidence may be a hint at the common origin of baryonic and dark matter. Moreover, this may be related to another important problem of the Standard Model - asymmetry between number of particles and antiparticles in our visible Universe. This line of reasoning resulted in construction of asymmetric dark matter models where dark matter candidates have a net (anti)-baryonic charge (see, e.g. [3–9] and Ref. [10] for a review).

An attractive effective model of asymmetric dark matter—hylogenesis—has been recently proposed in [11] (a particular high-energy completion of the hylogenesis scenario in the framework of supersymmetric models is presented in [12]). This model exhibits several distinct phenomenological signatures: (i) the induced nucleon decay (IND) [11, 13, 14]; (ii) the direct production of dark matter particles in particle collisions, in particular, at LHC [13]. The idea of collider searches for dark matter [15–17] has been recently got much attention especially with on-going LHC experiments. This kind of searches is especially important for the class of models with asymmetric dark matter where other standard searches based on direct detection or indirect searches for annihilation products are not (so) effective.

The goal of this paper is to get bounds on parameters of the hylogenesis model from searches for processes with a monojet and missing energy in the 1st LHC run and to estimate the signal for future LHC runs. This is the common signature for all dark matter models including WIMPs, though in hylogenesis the jet is originated from the new physics interaction rather than incoming quark or gluon bremsstrahlung. Apart from this we discuss also a way of discriminating the hylogenesis mechanism from other dark matter models by looking for the process with four jets, three of which emerge from a heavy state decay.

The paper is organized as follows. In Sec. II we describe the main features of hylogenesis mechanism. In Sec. III we discuss bounds on its parameters obtained from results of searches for single jet and missing energy in the 1st LHC run. In Sec. IV we estimate sensitivity of the future LHC runs and suggest other possible collider signatures of the hylogenesis. The last Section contains discussion and conclusions.

II. THE MODEL

In the hylogenesis model [11] the asymmetric dark matter consists of two components: fermion $Y$ and scalar $\Phi$. Each carries nonzero baryonic charge (see, e.g. [3–9] and Ref. [10] for a review).

Each carries nonzero baryonic charge such that $B_Y + B_\Phi = -1$. These fields couple to visible matter via “neutron”-like portal as

$$\mathcal{L} = - \sum_{a=1,2} \frac{\lambda_{ijk}}{M^2} \bar{X}_a P_R d^i \bar{u}^j C P_R d^k + \frac{\zeta_a}{M} X_a \bar{Y}^C \Phi^* + h.c., \quad (1)$$

where $X_a, a = 1, 2$ are heavy fermionic mediators, indices $i,j,k$ label generations and superscript $C$ denotes charge conjugation. In the first term of (1) we implicitly assume that the color indices are convoluted to form a color singlet. Masses of $X_{1,2}$ are supposed to be at TeV scale and obey $m_{X_2} > m_{X_1}$. Nonrenormalizable interaction in Eq. (1) is suppressed by a high energy scale $M$ where the model must be UV completed to become renormalizable.

Baryon asymmetry is generated by CP-asymmetric decays of nonthermal population of $X_1$ and $\bar{X}_1$. The rate
of the dominant decay mode $X_1 \rightarrow \tilde{Y} \Phi^*$ is
$$
\Gamma(X_1 \rightarrow \tilde{Y} \Phi^*) = \frac{|\zeta_1|^2 m_{X_1}}{16 \pi}.
$$
(2)

These heavy fermions can decay also as $X_1 \rightarrow udd$ and $\tilde{X}_1 \rightarrow \bar{u}d\bar{d}$ (for coupling to the first quark generation) with tree level decay rate
$$
\Gamma(X_1 \rightarrow udd) = \frac{3|\lambda_1|^2 m_{X_1}^3}{1024 \pi^3 M^4}.
$$
(3)

The baryon asymmetry in the visible sector is generated at one-loop level by the latter process (see Ref. [11] for details), which yields for the microscopic asymmetry (asymmetry per one decay of $X_1, \tilde{X}_1$ pair)
$$
\delta \equiv \frac{1}{2 \Gamma_{X_1}} (\Gamma(X_1 \rightarrow udd) - \Gamma(\tilde{X}_1 \rightarrow \bar{u}d\bar{d}))
\approx \frac{m_{X_1}^5}{256 \pi^3 |\zeta_1|^2 M^4 m_{X_2}}
\times 10^{-4} \times \frac{\text{Im}[|\lambda_1|^2 \zeta_1 \zeta_2]}{|\zeta_1|^2} \left(\frac{m_{X_1}}{M}\right)^4 \left(\frac{m_{X_2}}{m_{X_1}}\right),
$$
(4)

assuming $m_{X_2} \gg m_{X_1}$.

Large enough asymmetry requires that the scale $M$ should not significantly exceed the masses of $X_{1,2}$. Indeed, if generated at hot stage without any suppression and subsequent washing out the macroscopic asymmetry $\Delta_B \equiv (n_B - n_\bar{B})/s$ is estimated as $\Delta_B \sim \delta/g_s$, where $s$ and $g_s$ are the entropy density and number of relativistic degrees of freedom in plasma at that epoch. Macroscopic asymmetry remains constant later, while the Universe expands, so one has $\Delta_B \sim 10^{-10}$ [26], hence $\delta \sim 10^{-8}$ for anticipated $g_s \sim 10^2$. If the macroscopic asymmetry gets suppressed, than successful baryogenesis asks for larger $\delta$. Introducing small parameter $\epsilon$, so that $m_{X_1} \sim \epsilon m_{X_2}$ and $m_{X_2} \sim \epsilon M$ one observes from (4), that even $\epsilon < 0.3$ is unacceptable without a hierarchy in coupling constants, $\zeta_1 \ll \zeta_2$ smaller values, $\epsilon \ll 0.3$, become allowed.

Both dark matter particles $Y$ and $\Phi$ are supposed to be at GeV scale and their stability is a consequence of baryon number conservation and kinematical constraints on their masses: $|m_Y - m_\Phi| < m_p + m_e$. In this mechanism the net baryonic charge of the Universe is zero and it is supposed to be perturbatively conserved during the late-time evolution of the Universe. As a consequence relic number densities of baryons and dark matter particles are related as $n_Y = n_\Phi = n_B$ and using cosmological data this gives the following mass range for dark matter particles: $1.7 \text{ GeV} \lesssim m_Y, m_\Phi \lesssim 2.9 \text{ GeV}$.

Light dark matter particles can rescatter into quarks, $Y \Phi \rightarrow qgq$, and wash out the generated in the decays of $X$ asymmetry. This process is ineffective if reheating temperature is sufficiently low [11]
$$
\frac{T_{rh}}{2 \text{ GeV}} \lesssim \left( \sum_{a,b} \lambda_a \lambda_b \xi_a \xi_b \frac{\text{TeV}^6}{M^4 m_{X_a} m_{X_b}} \right)^{-1/5},
$$
(5)
yet it must exceed few MeV required by the standard Big Band Nucleosynthesis [26]. Note that mild hierarchy $\epsilon \sim 0.3$ is allowed by (5). Moreover, the new particles may be sufficiently below TeV scale. Say, with $\lambda_1 \sim \lambda_2 \sim \zeta_2 \sim 1$ and $\zeta_1 \sim 10^{-5}$ and $\epsilon \sim 0.1$ the process (2) still dominates over (3), while asymmetry (4) is at acceptable level, $\delta \sim 10^{-8}$ and reheating temperature (5) is above MeV scale for $M = 3 \text{ TeV}$, while the mass spectrum is $m_{X_1} \sim 30 \text{ GeV}$, $m_{X_2} \sim 300 \text{ GeV}$.

Hylogenesis mechanism can work for different combinations of flavors in the interaction (1), each is characterized by coupling $\lambda_a^{YY}$. Hereafter we analyze the simplest cases where only one of these constants is non-zero and contributes to both the baryon asymmetry generation and the production of dark matter particles in proton-proton collisions. We consider four cases (models) with different types of interaction
$$
\mathcal{O}^{dud} = \frac{\lambda_a^{dud}}{M^2} (\bar{X}_a P_R d)(\bar{u}^C P_R d),
$$
(6)
$$
\mathcal{O}^{ dus} = \frac{\lambda_a^{ dus}}{M^2} (\bar{X}_a P_R d)(\bar{u}^C P_R s),
$$
(7)
$$
\mathcal{O}^{dub} = \frac{\lambda_a^{dub}}{M^2} (\bar{X}_a P_R d)(\bar{u}^C P_R b),
$$
(8)
$$
\mathcal{O}^{dtd} = -\frac{\lambda_a^{dtd}}{M^2} (\bar{X}_a P_R d)(\bar{t}^C P_R d),
$$
(9)

where again convolution of color indices for quarks with $SU(3)$ antisymmetric tensor is implicitly assumed. Sure enough one can write down also operators with different permutations of quarks fields. Collider experiments are sensitive to all types of interactions presented above but the operators in (6), (7) also provide with IND signatures discussed in [13, 14] and can be probed in this way. As to operators (8), (9) collider searches is the only way to probe this scenario, since contribution from (8), (9) to proton decay are suppressed due to absence of the sea heavy quarks in proton.

TeV scale for $X_{1,2}$ was adopted in [11] because for smaller masses the proton lifetime starts to exceed the typical upper limits from Super-Kamiokande (see e.g. [27, 28]). Though the kinematics of IND (due to dark matter inelastic scattering) is different, so that Super-Kamiokande bounds are not directly applicable, light $X_{1,2}$ are expected to be disfavored from these searches. If operators (6), (7) are suppressed (absent), while (8), (9) are at work, the proton limits are significantly weaker and lower masses are acceptable even without strong hierarchy in coupling constants provided $\delta \gtrsim 10^{-8}$. With a hierarchy in coupling constants (e.g. like one suggested above) the masses may be significantly lower.

III. LHC Bounds

Here we discuss bounds on the model parameters from results of LHC searches for monojet and missing energy signature [18–20]. Sensitivity of the collider (Tevatron
and LHC) searches to hylogenesis model has been discussed in [13]. Below we consider the direct production of $X_{1,2}$ in association with a single jet. Since $X_{1,2}$ decay dominantly into dark matter particles $Y$ and $\Phi$ the corresponding events exhibit missing energy signature. In our calculations we limit ourselves to the case when narrow width approximation is valid. Corresponding tree level differential cross sections for different operators in (6)-(9) look as follows:

- Operator $\mathcal{O}^{dud}$. There are two main subprocesses contributing to reaction $pp \rightarrow jet + E_T^{miss}$, subprocess $d(p_1)d(p_2) \rightarrow \bar{u}(k_2)X_a(k_2)$ with differential cross section

$$\frac{d\sigma}{dt} = \frac{|\lambda^{dud}_a|^2}{96\pi s^2M^4} \left[4t^2 + 4(s - m^2_{X_a}) + s(s - m^2_{X_a})\right]$$

and subprocess $d(p_1)u(p_2) \rightarrow d(k_1)X_a(k_2)$ with

$$\frac{d\sigma}{dt} = \frac{|\lambda^{dud}_a|^2}{96\pi s^2M^4} \left[4s^2 + 4s(t - m^2_{X_a}) + t(t - m^2_{X_a})\right]$$

hereafter $s \equiv (p_1 + p_2)^2$, $t \equiv (p_1 - k_1)^2$.

- Operators $\mathcal{O}^{dus}$ and $\mathcal{O}^{dub}$. Here the main subprocess is $d(p_1)u(p_2) \rightarrow (s, b)(k_1)X_a(k_2)$ with differential cross section

$$\frac{d\sigma}{dt} = \frac{|\lambda^{dus(b)}_a|^2}{96\pi s^2M^4} (s + t - m^2_{X_a}) (s + t)$$

- Operator $\mathcal{O}^{dud}$. As in the previous case here we consider the dominant subprocess $d(p_1)d(p_2) \rightarrow t(k_1)X_a(k_2)$ whose differential cross section reads

$$\frac{d\sigma}{dt} = \frac{|\lambda^{dud}_a|^2}{96\pi s^2M^4} \left[4t^2 + 4(s - m^2_{X_a} - m^2_t) + s(s - m^2_{X_a} - m^2_t) + 4m^2_{X_a}m^2_t\right],$$

where we take into account the top quark mass.

For integration over the phase space with a particular set of cuts we use CompHEP package [21] with CTEQ6L1 [22] as a universal PDF set. The interactions of the type like the first term in Eq. (1) can not be introduced in the CompHEP directly (this soft can not deal with color singlet operators like $\epsilon_{\alpha\beta\gamma}q^\alpha q^\beta q^\gamma$). We generate some analog of considered process and replace the matrix element squared with those directly calculated from (6)-(9) (we have checked that the results (10)-(13) are reproduced by the code).

To obtain exclusion plots for model parameters we use the latest results of CMS searches for jet+$E_T^{miss}$ [20]. These searches at $\sqrt{s} = 8$ TeV with $19.7$ fb$^{-1}$ integrated luminosity adopted cut on jet rapidity $|\eta_{jet}| < 2.4$ and a cut on transverse missing energy $^{1}$.

$^{1}$ Moreover, the event selection procedure of [20] allows a second jet which is separated from the first by less than 2.5 azimuthal radius, $\Delta\phi(j_1, j_2) < 2.5$. Such angular requirement suppresses dijet QCD events.

The results of the analysis are presented as upper limits on number of events for a given cut on $E_T^{miss}$. We have found that the most stringent bounds on parameters of the model in question are obtained for the strongest cut $E_T^{miss} > 550$ GeV. Note that $E_T^{miss}$ in [20] is defined as the magnitude of vector sum of transverse momentum of all particles in the event. For our calculations this means that $E_T^{miss}$ coincides with transverse momentum of jet, $p_T^j$.

Using the procedure described above we calculate the number of signal events expected for some set of parameters and compare it with an observed 95% CL upper limit on number of events from new physics $N^{obs}_{obs} = 142$ for a given luminosity (see Table 3 in Ref. [20]). Note that we use tree level expressions for cross sections and expect that QCD corrections somewhat increase the values. So, the bounds obtained below are only conservative.

For a given operator of the type (1) the relevant model parameters are scale $M$, masses of heavy mediators $m_{X_1}, m_{X_2}$ and coupling constants $\lambda_1, \lambda_2$. To illustrate the obtained bounds we fix some of these parameters. We start with operator $\mathcal{O}^{dud}$. In Fig. 1 we show allowed regions in $(m_{X_1}, m_{X_2})$ plane for different values of $\lambda^{dud}(b)$ and $\lambda^{dud}(b) = 1$. Right panel: allowed regions in $(m_{X_1}, m_{X_2})$ plane for different values of $\lambda^{dub}(b)$ and $\lambda^{dub}(b) = 1$. We take $M = 3.5$ TeV and $m_{X_2} > m_{X_1}$.

FIG. 1. Left panel: allowed regions in $(m_{X_1}, m_{X_2})$ plane for different values of $\lambda^{dud}(b)$ and $\lambda^{dud}(b) = 1$. Right panel: allowed regions in $(m_{X_1}, m_{X_2})$ plane for different values of $\lambda^{dub}(b)$ and $\lambda^{dub}(b) = 1$. We take $M = 3.5$ TeV and $m_{X_2} > m_{X_1}$.

FIG. 2. Left panel: allowed regions in $(m_{X_1}, m_{X_2})$ plane for different values of $\lambda^{dub}(b)$ and $\lambda^{dub}(b) = 1$. Right panel: allowed regions in $(m_{X_1}, m_{X_2})$ plane for different values of $\lambda^{dub}(b)$ and $\lambda^{dub}(b) = 1$. We take $M = 2.5$ TeV and $m_{X_2} > m_{X_1}$.
left plot we fix $\lambda_{dud}^d = 1$ and show the allowed regions varying $\lambda_2^d$ from 0.3 to 2. We see that the limit on the mass of $X_1$ varies from 1.1 to about 2 TeV depending on $\lambda_2^d$. The limit is sensitive to $\lambda_2^d$ in the region of light $X_2$. On the right plot we fix $\lambda_{dud}^d = 1$ and show the allowed regions for different values of $\lambda_2^d$ as in the previous case. The obtained limits on masses of $M_{X_1}$ and $m_{X_2}$ in this case are up to 1.8 TeV for $M_{X_1}$ and in the interval 1.1–1.8 TeV for $m_{X_2}$.

Events with single b- or s-jets are equivalent from view point of models (7) and (8). In Fig. 2 we show allowed regions in $(m_{X_1}, m_{X_2})$ plane for both $O_{dus}$ and $O_{dub}$ interactions (which are practically the same due to smallness of $m_b$) for $M = 2.5$ TeV. Note, that bounds of the same order are expected for operators (6)-(8) with different permutations of the quark fields.

Next, we can consider a limiting case when $m_{X_2} \gg m_{X_1}$ and neglect contribution to the cross section from the heaviest mediator. Equivalently, here we set $\lambda_{dus}^d = 0$. In this case we obtain an exclusion plot in $(m_{X_1}, M/\sqrt{|\lambda_1^d|})$ plane shown in Fig. 3. Here $a$ labels the type of interaction, $a = dus(d), dud$.

From this plot we can see how the limit on $m_{X_1}$ scales with combination $M/\sqrt{|\lambda_1^d|}$ in front of the interaction lagrangian (1). We note that for small values of $M/\sqrt{|\lambda_1^d|}$ (i.e. on the right part of this exclusion plot) formally the limits on mass of $X_1$ become stronger. In this case the effective theory (1) with the interaction term suppressed by $M$ may be invalid and instead a full theory should be considered. But in this case the actual bounds will depend on UV completion of the model and in particular the width of $X_2$ resonances, see e.g. [12].

It follows from (10), (11) and (12), that the rate of the process $pp \rightarrow X_1 s(b)$ is less than the rate of $pp \rightarrow X_1 t\bar{t}$, so that the lower bounds on parameter $M/\sqrt{|\lambda_1^d|}$ are below the limits on $M/\sqrt{|\lambda_1^{dus}|}$. The exclusion region for operator $O_{dus}$ is wider than that for $O^{dus}$.

\section{IV. FURTHER TESTS AT LHC}

Here we calculate tree level cross sections for processes

\begin{equation}
pp \rightarrow \bar{d}X \hspace{1cm} pp \rightarrow bX \hspace{1cm} pp \rightarrow t\bar{X}
\end{equation}

at $\sqrt{s} = 13$ TeV. Note that the two last processes involve operators which can not be probed by IND process. In Fig. 4 we plot the cross section for $pp \rightarrow \bar{d}X$ (right panel) and $pp \rightarrow bX$ (left panel) at $\sqrt{s} = 13$ TeV. Here we fix $\lambda_{dud}^d = 0.5$ and neglect contribution from $X_2$.

Integration over the phase space is performed with the following cuts: $|\eta| < 2.4$ and $p_T > 550$ GeV for $pp \rightarrow dX$ and $p_T > 50$ GeV for $pp \rightarrow bX$. Obtained cross section is up to 200 fb for $dX$ final state and of order 10-50 fb for $bX$. In Fig. 5 we plot the cross section of process

\begin{equation}
pp \rightarrow t\bar{X}
\end{equation}

at $\sqrt{s} = 8$ TeV (left panel) and $\sqrt{s} = 13$ TeV (right panel) for comparison. The process has monotop signature (see e.g [23]). The current limits on monotop searches from results of ATLAS [25] and CMS [24] are of order $0.1 - 1$ pb and the limits on parameters of the hylogenesis model are weak.

We finish the Section with brief discussion of other rather specific signatures. The successful baryogenesis requires sufficient CP-asymmetry in decays $X_1 \rightarrow udd$ and $X_1 \rightarrow \bar{u}dd$. The expression for partial width of these decays is given by (3). Comparing (3) with width of the dominant decay (2) one can see that for $|\lambda_1| \sim |\lambda| \sim 1$ and for $m_{X_1} \lesssim M$ the branching of $X_1$ decay to quarks can reach values $5 \times 10^{-3}$. Even larger values are possible if $|\lambda_1| < |\lambda|$. So, the invisible decay of $X_1$ is not the only signature to search for hylogenesis at LHC. Another

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig3.png}
\caption{Exclusion region in $(M/\sqrt{\lambda_1^d}, m_{X_1})$ plane for $\lambda_2^d = 0$.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig4.png}
\caption{Left panel: cross section of process $pp \rightarrow \bar{d}X$ at $\sqrt{s} = 13$ TeV. Right panel: cross section of process $pp \rightarrow bX$ at $\sqrt{s} = 13$ TeV.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig5.png}
\caption{Left panel: cross section for the process $pp \rightarrow bX$ at $\sqrt{s} = 8$ TeV. Right panel: cross section for the process $pp \rightarrow t\bar{X}$ for $\sqrt{s} = 13$ TeV.}
\end{figure}
interesting possibility appears when produced $X_1$ decays into three quarks and we have four-jet signature. For $O_{ud}^d$ operator processes of interest are

$$ud \rightarrow dX_1 \rightarrow d\,udd, \quad dd \rightarrow uX_1 \rightarrow u\,udd.$$  \hspace{1cm} (15)

In case of $O_{dub}^d$ operator the relevant to searches for hylogenesis signals are 4-jets (involving 2b-jets) in the final state,

$$ud \rightarrow bX_1 \rightarrow b\,dub.$$  \hspace{1cm} (16)

For $O_{dd}^d$ events with 2jets$+\ell\ell$ arise from subprocess

$$dd \rightarrow tX_1 \rightarrow t\,dd.$$  \hspace{1cm} (17)

We leave an analysis of these signatures for future study.

V. CONCLUSIONS

In this paper we explore collider phenomenology of the hylogenesis model. We have found that the current searches for jet and missing transverse energy signature at ATLAS and CMS place bounds on mass of the lightest mediator $X_1$ of order $0.7 - 2$ TeV for interactions mediated by operators $O_{ud}^d$, $O_{dus}^d$ and $O_{dub}^d$ suppressed by scale $M = 2.5 - 3.5$ TeV. At the same time current monotop searches do not allow to put considerable bounds on the parameters of the model for $M$ in TeV range. For the next run of LHC experiments at $\sqrt{s} = 13$ TeV we predict cross section up to 200 fb for jet$+E_T^{miss}$ searches and up to tens fb for monotop process. Four jets signature (including 2jets$+b\bar{b}$ and 2jets$+t\bar{t}$) is suggested to probe hylogenesis at LHC.

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