Dressed Gluon Exponentiation for Inclusive B–decay Spectra

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Plan of the talk

- Introduction: inclusive B decay spectra — the challenge
- Sudakov resummation with NNLL accuracy, divergence of perturbation theory
- Dressed Gluon Exponentiation: renormalon resummation in the Sudakov exponent
- The quark distribution function in a meson and in an on-shell heavy quark
- cancellation of the leading renormalon ambiguity
- Numerical results for inclusive B-decay spectra by DGE; comparison to data

Dressed Gluon Exponentiation for Inclusive B–decay Spectra

References

- Inclusive spectra in charmless semileptonic B decays by DGE,
  J.R. Andersen, E. Gardi, [hep-ph/0509360].

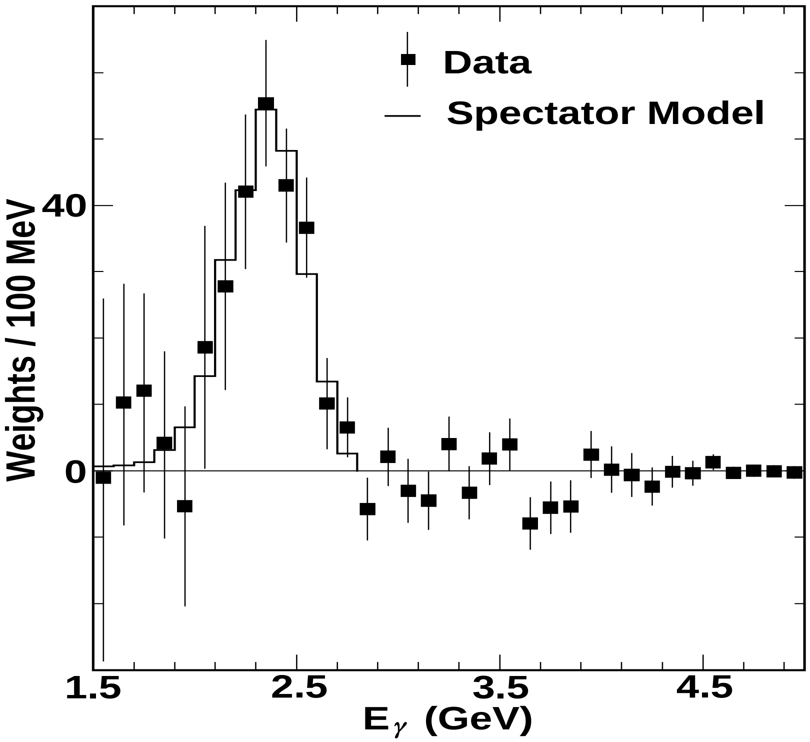
- Taming the $\bar{B} \longrightarrow X_s\gamma$ spectrum by Dressed Gluon Exponentiation,
  J.R. Andersen, E. Gardi, JHEP **0506** 030 (2005) [hep-ph/0502159].

- On the quark distribution in an on-shell heavy quark and its all-order relations with the perturbative fragmentation function,
  E. Gardi, JHEP **0502**, 053 (2005) [hep-ph/0501257].

- Radiative and semi-leptonic B-meson decay spectra: Sudakov resummation beyond logarithmic accuracy and the pole mass,
  E. Gardi, JHEP **0404**, 049 (2004) [hep-ph/0403249].

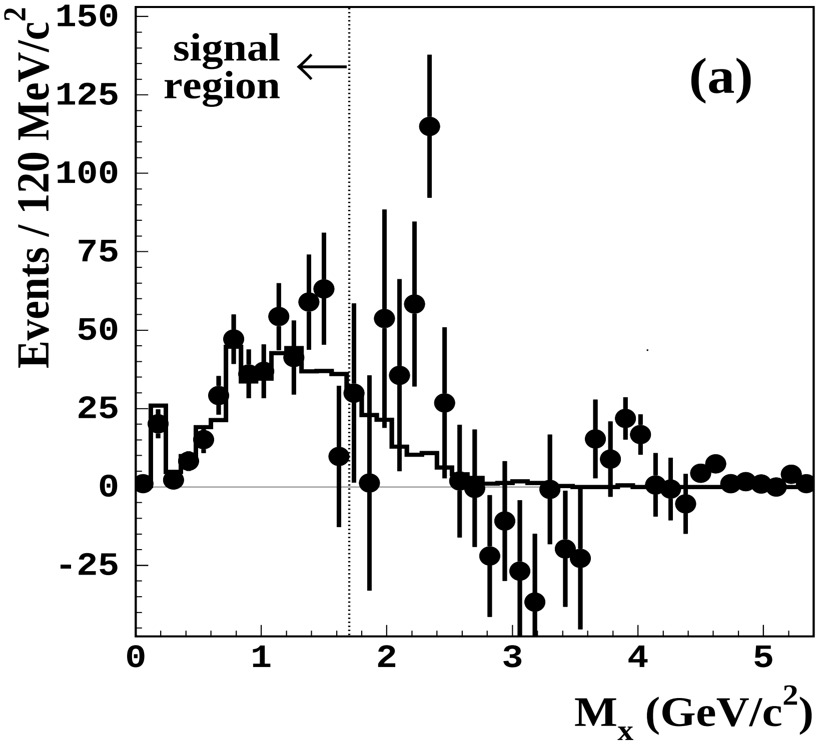
Inclusive B–decay Spectra

radiative decay: $\bar{B} \longrightarrow X_s\,\gamma$



CLEO

semi-leptonic decay: $\bar{B} \longrightarrow X_u\,l\bar{\nu}_l$



BELLE

The distribution peaks close to the endpoint $(E_\gamma \longrightarrow M_B/2; \quad \text{small } M_X)$

Example: extracting $|V_{ub}|$ from the semi-leptonic decay

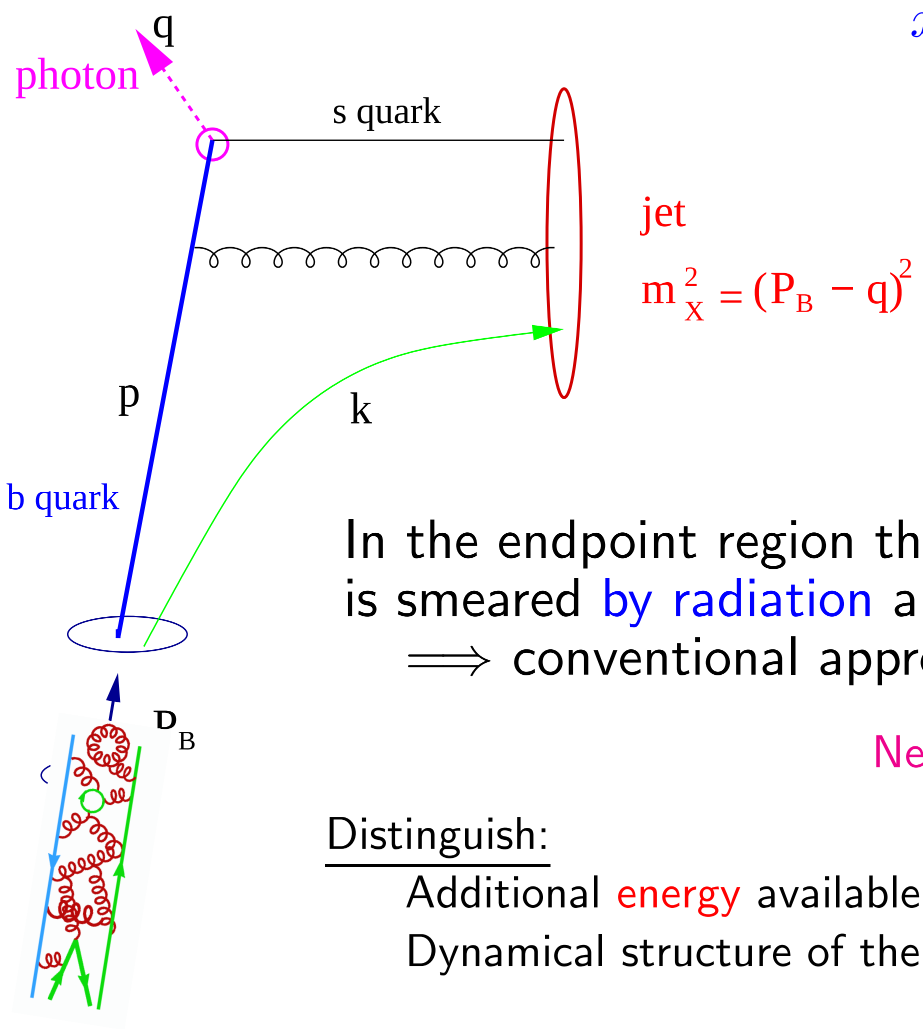
Precise measurements are restricted to the small M_X region (charm background)

Determination of $|V_{ub}|$ relies on calculation of the spectrum.

Kinematics in $\bar{B} \longrightarrow X_s \, \gamma$

In the B meson the b quark is close to its mass shell.
Therefore, perturbation theory (with an on-shell quark initial state) applies (up to power corrections...).

$$x \equiv \frac{2E_\gamma}{m_b}; \qquad \frac{1}{\Gamma_{\text{tot}}} \frac{d\Gamma}{dx}\bigg|_{\text{LO}} = \delta(1-x)$$



Perturbative endpoint: $x = 1$
Physical endpoint: $x = M_B/m_b > 1$

In the endpoint region the distribution is smeared by radiation and by the primordial motion of the quark
$\Longrightarrow$ conventional approach: leading power NP "shape function".

Neubert; Bigi, Shifman, Uraltsev & Vainshtein (93)

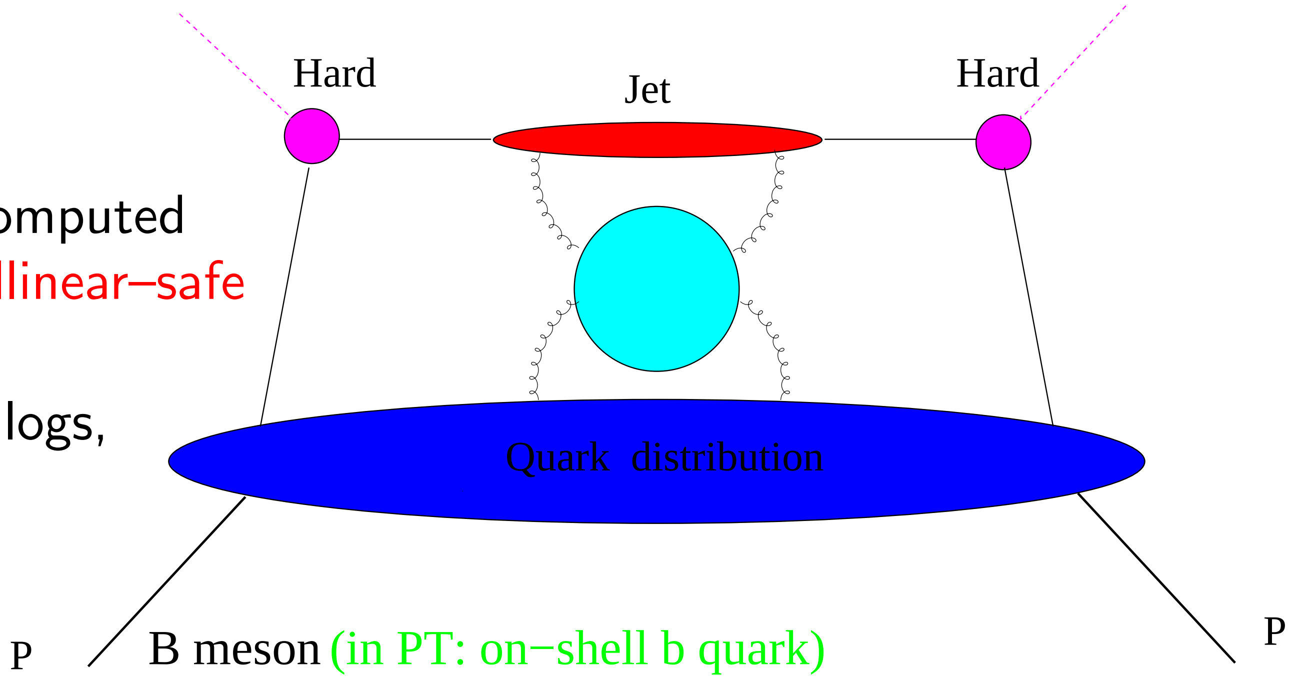
Distinguish:

Additional energy available in the meson $\bar{\Lambda} = M_B - m_b$
Dynamical structure of the meson

Large–x factorization in inclusive B decays

The spectrum can be computed in PT: infrared– and collinear–safe

Dominated by Sudakov logs, $\ln(1-x)$



scales:

Hard: m

Jet: $m_X^2 = (P_b - q)^2 \simeq m^2(1-x) \Longrightarrow m^2/N$

Soft: $m(1-x) \Longrightarrow m/N$

Korchemsky & Sterman (94)

Spectral moments:

$$\begin{aligned} \Gamma_N^{\rm PT} &\equiv \int_0^1 dx x^{N-1} \frac{1}{\Gamma_{\rm tot}^{\rm PT}} \frac{d\Gamma^{\rm PT}}{dx} \\ &= H(m) J(m^2/N;\mu) S_{\rm PT}(m/N;\mu) \; + \; \mathcal{O}(1/N) \\ &\equiv H(m)\,{\rm Sud}(N,m) \qquad + \; \mathcal{O}(1/N) \end{aligned}$$

Coefficients in the Sudakov exponent

$$\mathrm{Sud}(N,m) = \exp\left\{ -\sum_{n=1}^{\infty}\sum_{k=1}^{n+1} C_{n,k} \ln^k N \left(\frac{\alpha_s^{\overline{\mathrm{MS}}}(m^2)}{\pi}\right)^n \right\}$$

The coefficients $C_{n,k}$ are known **exactly** to **NNLL accuracy** [Gardi (2005)]
For $N_f = 4$ $C_{n,k}$ are:

$\boldsymbol{k} \quad \longrightarrow$

$\boldsymbol{n}$ $\downarrow$

-1.564	0.667	0	0	0	0	0	0
3.837	-0.078	1.389	0	0	0	0	0
$?$	20.579	6.339	3.376	0	0	0	0
$?$	$?$	116.464	33.024	9.042	0	0	0
$?$	$?$	$?$	597.221	138.600	25.955	0	0
$?$	$?$	$?$	$?$	2859.284	548.170	78.492	0
$?$	$?$	$?$	$?$	$?$	13141.289	2129.058	247.233
$?$	$?$	$?$	$?$	$?$	$?$	58941.217	8238.359
$?$	$?$	$?$	$?$	$?$	$?$	$?$	260391.559

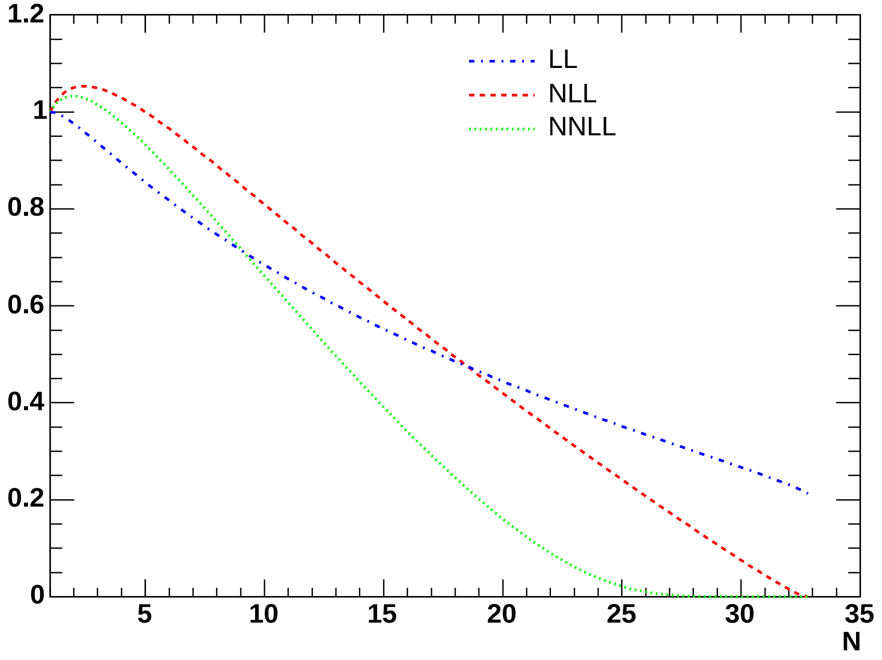
- At a given order in α_s the coefficients of **subleading logs** (lower k) get large...
- Is the fixed–logarithmic–accuracy approximation at LL / NLL / NNLL good?

Conventional Sudakov resummation with NNLL accuracy

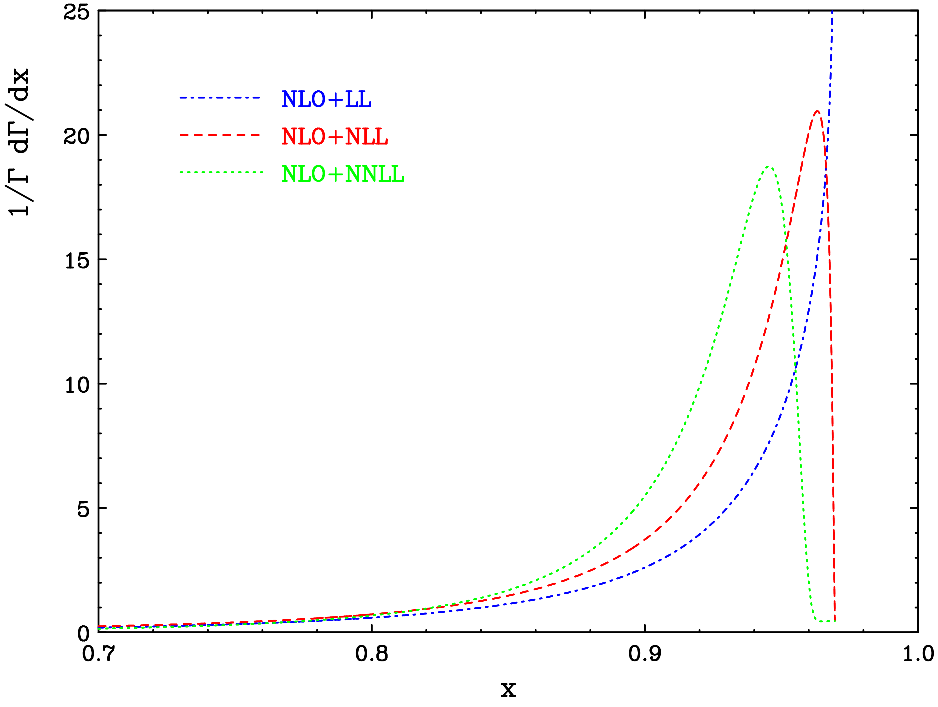
$$\mathrm{Sud}(N,m) \quad = \quad \exp\left\{\sum_{n=0}^{\infty} g_n(\lambda)\left(\frac{\alpha_s^{\overline{\mathrm{MS}}}(m^2)}{\pi}\right)^{n-1}\right\}; \qquad \lambda \equiv \frac{\alpha_s^{\overline{\mathrm{MS}}}(m^2)}{\pi}\beta_0 \ln N$$

$$g_0(\lambda) = \frac{C_F}{{\beta_0}^2}\left[(1-\lambda)\,\ln(1-\lambda) - \frac{1}{2}(1-2\lambda)\,\ln(1-2\lambda)\right]$$

$\mathrm{Sud}(N,m)$



Corresponding spectra



Coefficients in the Sudakov exponent in the large–β_0 limit

$$\mathrm{Sud}(N,m) = \exp\left\{-\sum_{n=1}^{\infty}\sum_{k=1}^{n+1} C_{n,k}\ln^k N \left(\frac{\alpha_s^{\overline{\mathrm{MS}}}(m^2)}{\pi}\right)^n\right\}$$

The part in $C_{n,k}$ that is proportional to $(\beta_0)^{n-1}$ is known to all orders:

$$k \longrightarrow$$

$$n \downarrow$$

-1.56	0.67	0	0	0	0	0
1.24	0.90	1.39	0	0	0	0
61.17	28.32	8.28	3.38	0	0	0
1096.06	515.20	166.25	34.89	9.04	0	0
20399.23	10078.43	3231.40	793.25	131.33	25.95	0
444615.21	221481.03	73268.94	17791.58	3514.66	482.12	78.49
11342675.74	5665794.49	1883129.50	468180.33	91361.30	15080.79	1768.50
334032127.30	166960507.50	55609620.17	13867704.58	2760946.21	449959.01	63745.75

- $C_{n,k}$ **increase** for lower powers of $\ln N$, building up $\sum_{k=1}^{n+1} C_{n,k}\ln^k N \sim \mathbf{n!} f_n(N)$
- Truncation at fixed logarithmic accuracy is **not** a good approximation.
- **Renormalon divergence** sets in already at low orders — requires a prescription!

Dressed Gluon Exponentiation — the jet function

Borel representation of the Sudakov exponent:

$$\ln J_N(Q;\mu_F) = \int_0^1 dx \frac{x^{N-1}-1}{1-x}\left[\int_{\mu_F^2}^{(1-x)Q^2}\frac{d\mu^2}{\mu^2}\mathcal{A}\Big(\alpha_s(\mu^2)\Big) + \mathcal{B}\Big(\alpha_s((1-x)Q^2)\Big)\right]$$

$$= -\frac{C_F}{\beta_0}\int_0^\infty \frac{du}{u}\left(\frac{\Lambda^2}{Q^2}\right)^u \times \left[B_{\mathcal{J}}(u)\Gamma(-u)\left(N^u-1\right) + \left(\frac{Q^2}{\mu_F^2}\right)^u B_{\mathcal{A}}(u)\ln N\right],$$

we defined $B_{\mathcal{J}}(u) \equiv B_{\mathcal{A}}(u) - uB_{\mathcal{B}}(u)$ and used the Borel representation of the anomalous dimensions,

$$\mathcal{A}\Big(\alpha_s(\mu^2)\Big) = \frac{C_F}{\beta_0}\int_0^\infty du \left(\frac{\Lambda^2}{\mu^2}\right)^u B_{\mathcal{A}}(u); \qquad \mathcal{B}\Big(\alpha_s(\mu^2)\Big) = \frac{C_F}{\beta_0}\int_0^\infty du \left(\frac{\Lambda^2}{\mu^2}\right)^u B_{\mathcal{B}}(u),$$

$$\int_0^1 dx x^{N-1}(1-x)^{-1-u} = \frac{\Gamma(-u)\Gamma(N)}{\Gamma(N-u)} \simeq \Gamma(-u)\,N^u \times (1+\mathcal{O}(1/N)).$$

In the large β_0 limit $B_{\mathcal{J}}(u) = \mathrm{e}^{\frac{5}{3}u}\frac{\sin\pi u}{\pi u} \times \frac{1}{2}\left(\frac{1}{1-u}+\frac{1}{1-u/2}\right) \times \Big(1+\mathcal{O}(u/\beta_0)\Big).$

Infrared sensitivity appears as renormalon ambiguity in the Sudakov exponent
⇒ parametrically–enhanced power corrections $\mathcal{O}(N\Lambda^2/Q^2)$ in the exponent

Dressed Gluon Exponentiation — the soft function

Borel representation of the soft Sudakov exponent:

$$\ln S_N(Q;\mu_F) = \int_0^1 dx \frac{x^{N-1}-1}{1-x}\left[\int_{(1-x)^2Q^2}^{\mu_F^2}\frac{d\mu^2}{\mu^2}\mathcal{A}\Big(\alpha_s(\mu^2)\Big) - \mathcal{D}\Big(\alpha_s((1-x)^2Q^2)\Big)\right]$$

$$= \frac{C_F}{\beta_0}\int_0^\infty \frac{du}{u} {\frac{\Lambda^2}{Q^2}}^{u}\left[B_{\mathcal{S}}(u)\Gamma(-2u)\left(N^{2u}-1\right) + {\frac{Q^2}{\mu_F^2}}^{u} B_{\mathcal{A}}(u)\ln N\right],$$

where we defined $B_{\mathcal{S}}(u) \equiv B_{\mathcal{A}}(u) - uB_{\mathcal{D}}(u)$.

What does one gain?

- Resummation of running–coupling effects beyond the available logarithmic accuracy
- Upon choosing a prescription (e.g. PV) for the Borel integral, *the divergent sum is defined*.
- Cancellation of certain renormalon ambiguities can then take place.
- Landau singularities are absent.
- The pattern of power corrections (observable dependent) can be studied: singularities in $\Gamma(-2u) \Longrightarrow$ power corrections $(N\Lambda/Q)^k$ in the exponent, except for $B_{\mathcal{S}}(u) = 0$.

However, QCD perturbation theory gives the power expansion: $B_{\mathcal{S}}(u) = 1 + s_1 u + \cdots$

For DGE one needs to know $B_{\mathcal{S}}(u)$ also away from the origin — involves assumptions!

Soft anomalous dimensions in the large–β_0 limit

$$B_{\mathcal{S}}(u) = \mathrm{e}^{\frac{5}{3}u}\frac{\sin \pi u}{\pi u}\, b_{\mathcal{S}}(u) \;\times\; \Big(1 + \,\mathcal{O}(u/\beta_0)\Big)$$

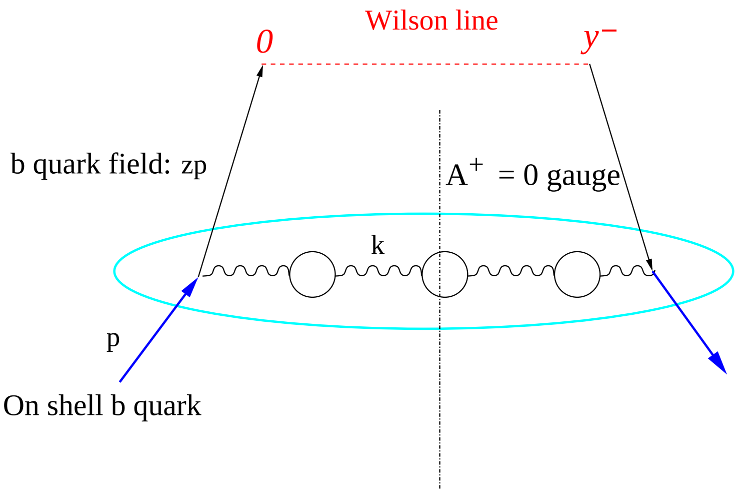
Observable	$b_{\mathcal{S}}(u)$	$B_{\mathcal{S}}(u) = 0$	power corrections
Drell-Yan (2)	$\frac{\Gamma^2(1-u)}{\Gamma(1-2u)}$	$u = \frac{1}{2}, \frac{3}{2}, \ldots$	$\left(\frac{\Lambda N}{Q}\right)^k$, k even
Heavy Jet Mass (1) / Thrust (2)	1		$\left(\frac{\Lambda N}{Q}\right)^k$, k integer
c parameter (2)	$\frac{\Gamma^2(1+u)}{\Gamma(1+2u)}$		$\left(\frac{\Lambda N}{Q}\right)^k$, k integer
Heavy Quark Fragmentation (1) Heavy Quark Distribution (1) $(Q^2 = m^2)$	$(1-u)\frac{\pi u}{\sin \pi u}$	$u = 1$	$\left(\frac{\Lambda N}{m}\right)^k$,$k \neq 2$

The quark distribution function

$$F_{\text{PT}}(N;\mu) \quad \underset{\text{large N}}{\longrightarrow} \quad \left\langle b(p_b) \left| \left[\bar{\Psi}(y)\gamma^+ \,\Phi_y(0,y)\,\Psi(0)\right]_\mu \right| b(p_b) \right\rangle \Big|_{ip_b^+ y^- \longrightarrow N} = H(m_b,\mu)\,\mathcal{S}\left(\frac{N\mu}{m_b}\right)$$

$$\mathcal{S}\left(\frac{N\mu}{m_b}\right) = \exp\left\{\frac{C_F}{\beta_0}\int_0^\infty \frac{du}{u}\left(\frac{\Lambda^2}{\mu^2}\right)^u \left[B_{\mathcal{S}}(u)\Gamma(-2u)\left(\left(\frac{N\mu}{m_b}\right)^{2u} - 1\right) + B_{\mathcal{A}}(u)\ln\left(\frac{N\mu}{m_b}\right)\right]\right\}$$

$$\text{with } B_{\mathcal{S}}(u) = e^{\frac{5}{3}u}\,(1-u) \times \left(1 + \mathcal{O}(u/\beta_0)\right)$$
$$= 1 + s_1 u + s_2 u^2/2! + \cdots$$



Renormalon in the exponent and their interpretation:

- Leading renormalon $u=\frac{1}{2}$, $\mathcal{O}(\Lambda N/m_b)$, is related to the *mass* of $\langle b(p_b)|$: $e^{-i\,\delta m\,y^-} = e^{-\delta m\,N/m_b}$
- Higher renormalons $u \geq \frac{3}{2}$, $(\Lambda N/m_b)^k$ with $k \geq 3$, correspond to the difference between the momentum distribution in the on-shell quark and the (unambiguous) distribution in the meson:

$$F(N;\mu) = \left\langle B(P_B) \left| \left[\bar{\Psi}(y)\gamma^+ \,\Phi_y(0,y)\,\Psi(0)\right]_\mu \right| B(P_B) \right\rangle \Big|_{iP_B^+ y^- \longrightarrow N} + \mathcal{O}(1/N)$$

Cancellation of the leading renormalon ambiguity

Owing to kinematic power corrections, the resummed E_γ spectrum is not influenced by the $u = \frac{1}{2}$ $\mathcal{O}(N\Lambda/m_b)$ ambiguity of the perturbative Sudakov exponent:

$$\frac{1}{\Gamma_{\text{tot}}}\frac{d\Gamma}{dE_\gamma} = \frac{2}{m_b}\int_{c-i\infty}^{c+i\infty}\frac{dN}{2\pi i}\left(\frac{2E_\gamma}{m_b}\right)^{-N} H(m_b)\,\underbrace{J(m_b^2/N;\mu)S_{\text{PT}}(m_b/N;\mu)}_{\text{Sud}(m_b,N)\;-\;\text{ambiguous}}$$

$$\simeq \frac{2}{M_B}\int_{c-i\infty}^{c+i\infty}\frac{dN}{2\pi i}\left(\frac{2E_\gamma}{M_B}\right)^{-N} H(m)J(m_b^2/N;\mu)\,\underbrace{S_{\text{PT}}(m_b/N;\mu)\,\text{e}^{-(N-1)\bar{\Lambda}/m_b}}_{u=\frac{1}{2}\text{ prescription independent}}$$

The cancellation is exact in all the moments, but it requires

- renormalon resummation in the Sudakov exponent
- renormalon resummation in $\bar{\Lambda} = M_B - m_b$ using the same prescription.

Sudakov resummation beyond logarithmic accuracy

$$\mathrm{Sud}(m,N)\big|_{\mathrm{PV}} = \exp\left\{\frac{C_F}{\beta_0}\mathrm{PV}\int_0^\infty du\, T(u)\left(\frac{\Lambda^2}{m^2}\right)^u \times\frac{1}{u}\left[B_{\mathcal{S}}(u)\Gamma(-2u)\left(N^{2u}-1\right) - B_{\mathcal{J}}(u)\Gamma(-u)\left(N^{u}-1\right)\right]\right\}.$$

What do we know about $B_{\mathcal{S}}(u)$?

- NNLO in the full theory: $B_{\mathcal{S}}(u) = 1 + s_1\frac{u}{1!} + s_2\frac{u^2}{2!} + \cdots$
- Renormalon cancellation in $\mathrm{Sud}(m,N)\,\mathrm{e}^{-(N-1)\bar{\Lambda}/M}$ implies: $B_{\mathcal{S}}(u=1/2)$ is equal in magnitude and opposite in sign to the residue of the $u=1/2$ renormalon in $m/m_{\overline{\mathrm{MS}}}$, which can be determined from the known NNLO expansion in $\overline{\mathrm{MS}}$ within a few percent.
- All orders in the large–β_0 limit: $B_{\mathcal{S}}(u) = \mathrm{e}^{\frac{5}{3}u}\left(1-u\right) + \mathcal{O}(1/\beta_0^2)$. The vanishing of $B_{\mathcal{S}}(u)$ at $u=1$ is assumed to hold in general.

$\bar{B} \longrightarrow X_s\gamma$ spectrum: from moment space to E_γ

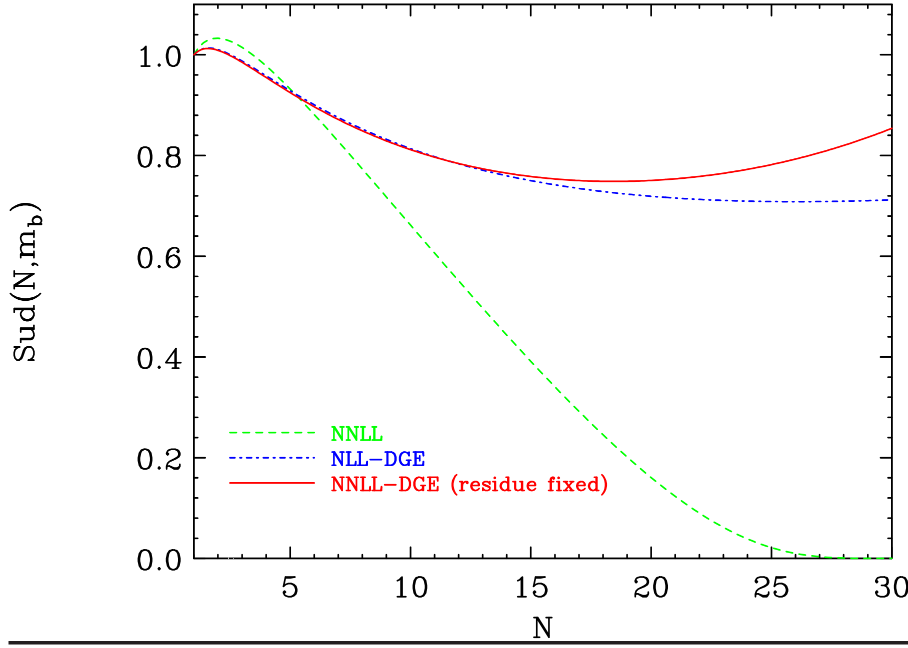
$$\mathrm{Sud}(m,N)|_{\mathsf{PV}} \;=\; \exp\left\{\frac{C_F}{\beta_0}\mathrm{PV}\;{}_{0}^{\infty}\, du\, T(u)\quad \frac{\Lambda^2}{m^2}^{\;u}\right.$$

$$\left.\times\frac{1}{u}\left[B_{\mathcal{S}}(u)\Gamma(-2u)\left(N^{2u}-1\right)-B_{\mathcal{J}}(u)\Gamma(-u)\left(N^{u}-1\right)\right]\right\}.$$

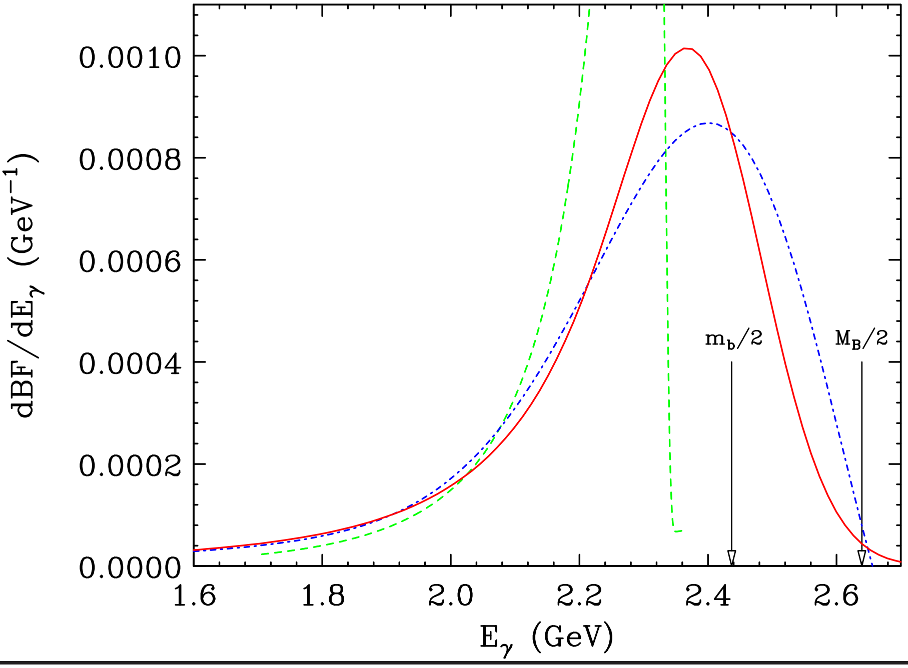
$$\frac{d\Gamma(E_\gamma)}{dE_\gamma} = \frac{m_{\mathsf{PV}}}{2}\;{}_{c-i\infty}^{c+i\infty}\frac{dN}{2\pi i}H(m)\;\mathrm{Sud}(m,N)|_{\mathsf{PV}}\left(\frac{2E_\gamma}{m_{\mathsf{PV}}}\right)^{-N}$$

Modified support properties:

$\mathrm{Sud}(N,m)|_{\mathsf{PV}}$ with various approx. for $B_{\mathcal{S}}(u)$



Corresponding spectra



$B_{\mathcal{S}}(u)$ away from the origin

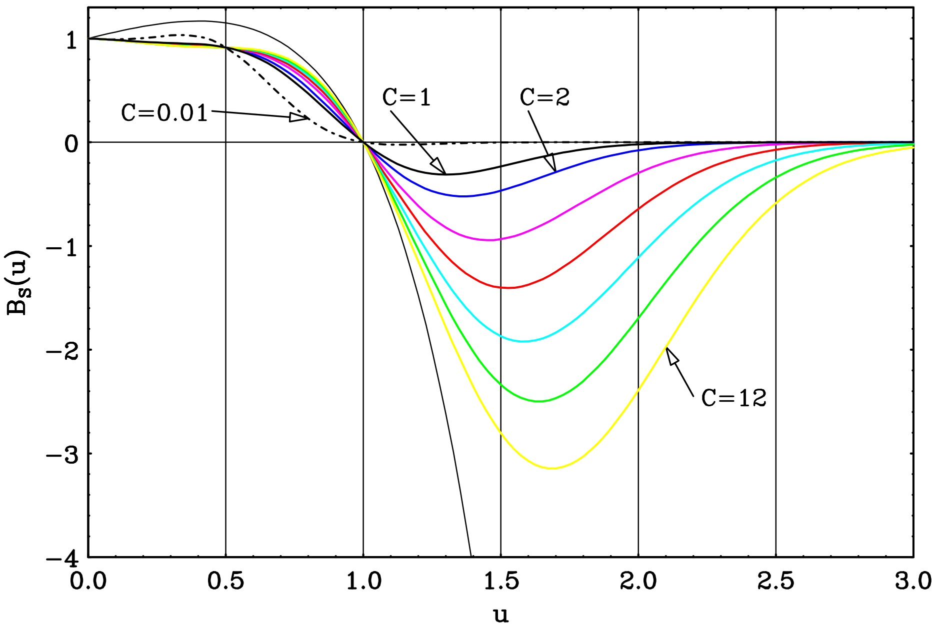
Ansatz for $B_{\mathcal{S}}(u)$ that is consistent with the known $\mathcal{O}(u^2)$ result in QCD (and the large–β_0 limit):

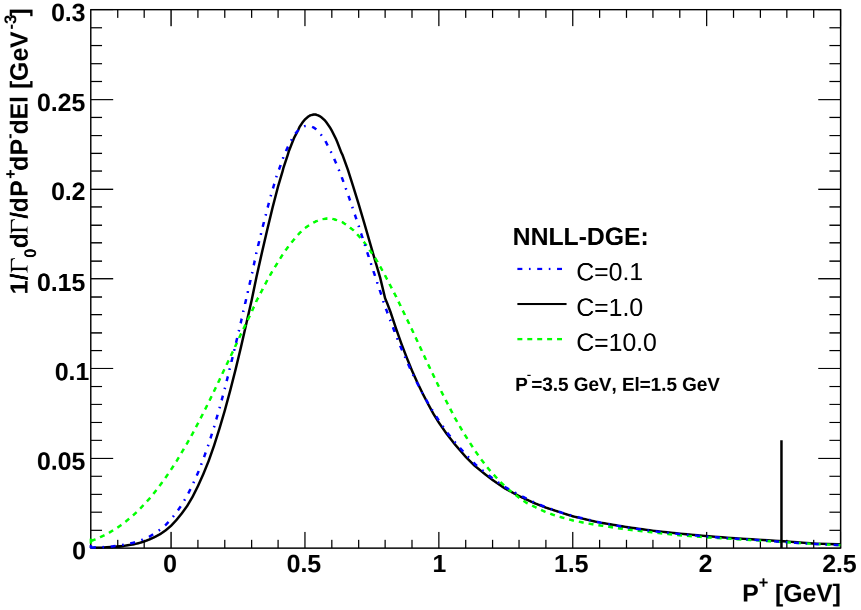
$$B_{\mathcal{S}}(u) = e^{\frac{5}{3}u}(1-u) \times \exp\left\{ c_2 u + \frac{1}{2}\left[c_3 - c_2^2 + \frac{C_A}{\beta_0}\left(\frac{5}{18}\pi^2 + \frac{7}{9} - \frac{9}{2}\zeta_3 \right) \right] u^2 \right\} \times W(u)$$

$$W(u) \equiv e^{t_1 u + \frac{1}{2} t_2 u^2} \left(1 - t_1 u + \frac{1}{2}(t_1^2 - t_2) u^2 \right) = 1 + \mathcal{O}(u^3).$$

Here $t_{1,2}$ are fixed requiring:

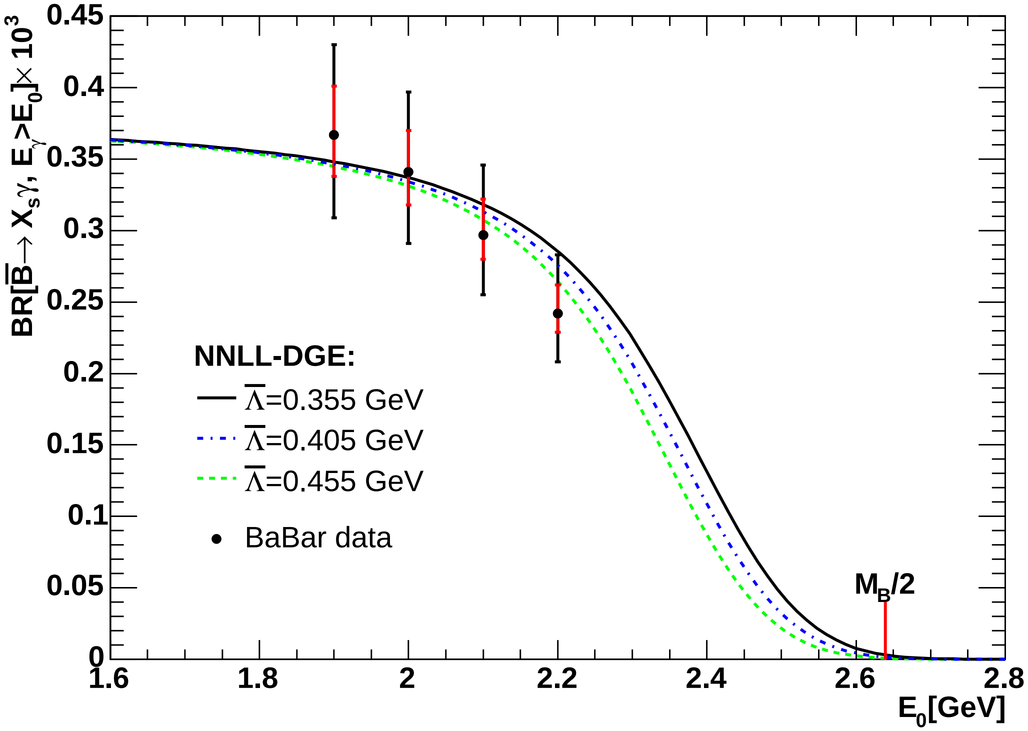
$$B_{\mathcal{S}}(u = 1/2) \ = \ 0.914 \ \pm 3\% \, (\text{computed}); \quad B_{\mathcal{S}}(u = 3/2) \ = \ -0.23366 \ \times \ C,$$





comparison to data: $\bar{B} \longrightarrow X_s\gamma$ branching fraction

- Theoretical uncertainty on the total BF $\sim 10\%$
- Experimental cuts on E_γ *do not* significantly increase the overall uncertainty.
- The measured BF is consistent with the Standard Model.

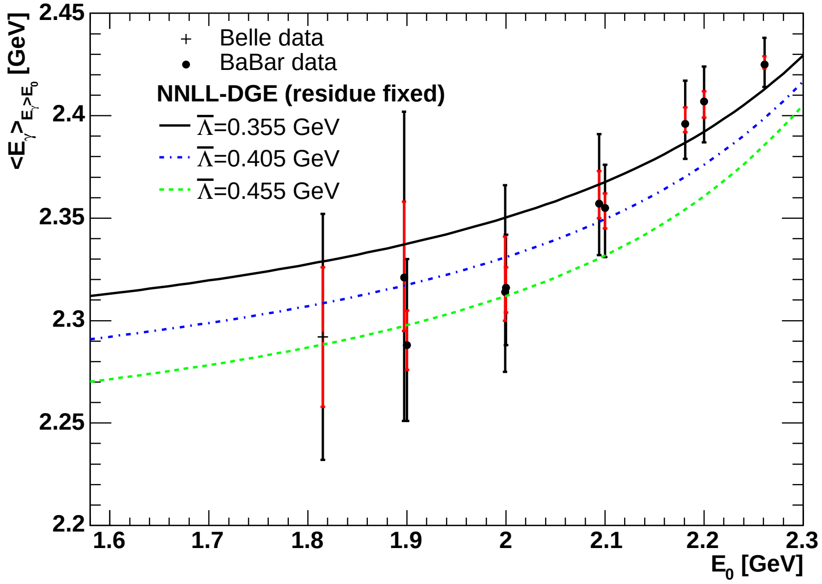


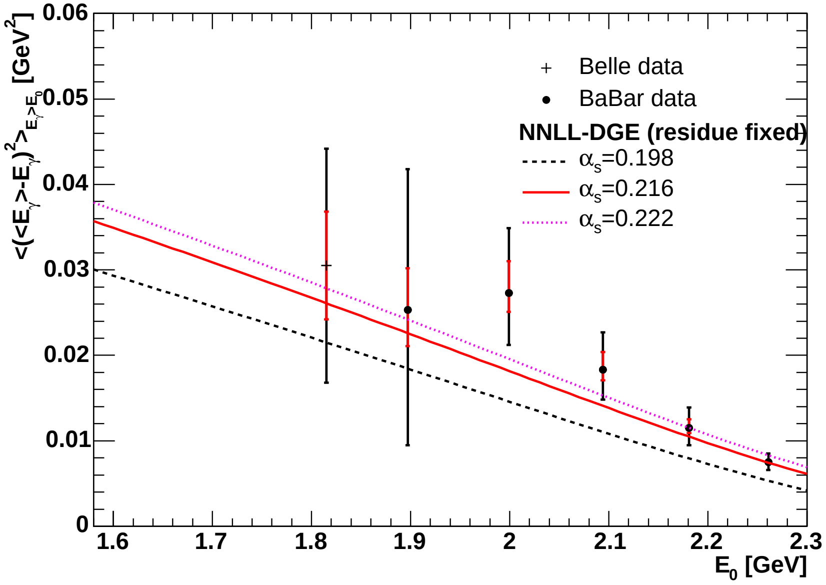
- Possible determination of m_b!

comparison to data: cut moments in $\bar{B} \longrightarrow X_s\gamma$

$$\Big\langle E_\gamma \Big\rangle_{E_\gamma > E_0} \equiv \frac{1}{\Gamma(E_\gamma > E_0)} \quad_{E_0} dE_\gamma \frac{d\Gamma(E_\gamma)}{dE_\gamma} E_\gamma$$

$$\left\langle \left(\Big\langle E_\gamma \Big\rangle_{E_\gamma > E_0} - E_\gamma \right)^n \right\rangle_{E_\gamma > E_0} \equiv \frac{1}{\Gamma(E_\gamma > E_0)} \quad_{E_0} dE_\gamma \frac{d\Gamma(E_\gamma)}{dE_\gamma} \left(\Big\langle E_\gamma \Big\rangle_{E_\gamma > E_0} - E_\gamma \right)^n .$$





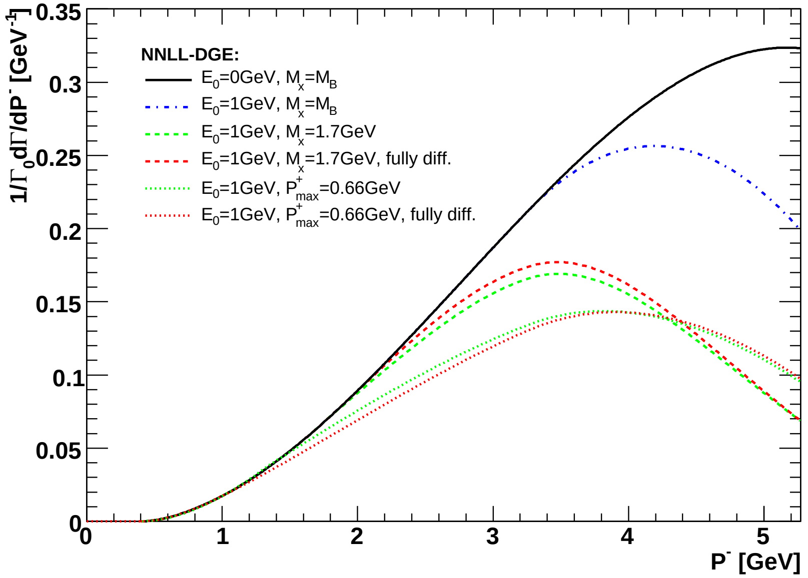
- The comparison suggests that power corrections are indeed small.
- In future: possible measurement of power corrections.

Integrated $\bar{B} \longrightarrow X_u l \bar{\nu}$ spectrum

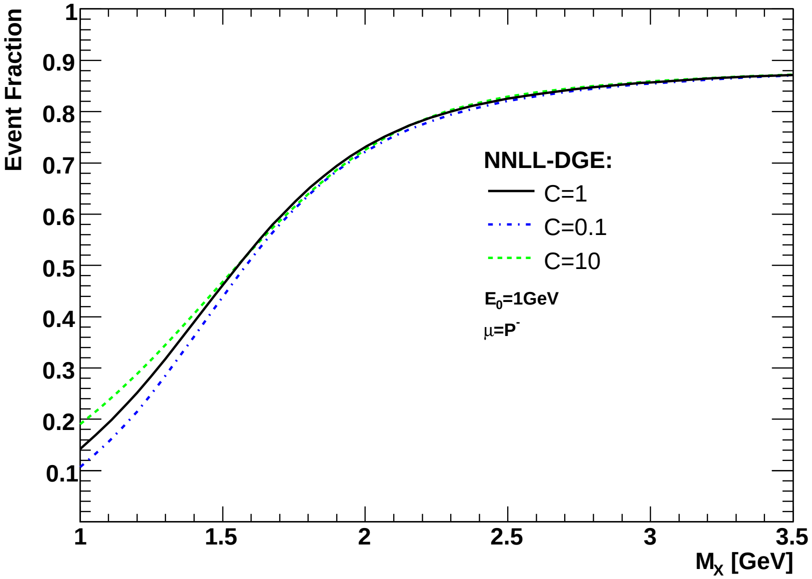
Integrating the spectrum with given experimental cuts:

- Hadronic Mass Cut: $P^+P^- < (1.7\,\mathrm{GeV})^2$, $E_l > 1\,\mathrm{GeV}$
- Small Lightcone Component Cut: $P^+ < 0.66\,\mathrm{GeV}$, $E_l > 1\,\mathrm{GeV}$

The effect of cuts on the P^- spectrum



Sensitivity of the Event Fraction to C



Extracting for $|V_{ub}|$ from Belle data

$$\Delta\mathcal{B}(\bar{B} \longrightarrow X_u l \bar{\nu} \text{ restricted phase space}) = \tau_B \Gamma_{\text{tot}} \left(\bar{B} \longrightarrow X_u l \bar{\nu}\right) \times R_{\text{cut}}.$$

From Belle data

$$\begin{aligned} \Delta\mathcal{B}(P^+P^- < (1.7\,\text{GeV})^2,\ E_l > 1\,\text{GeV}) &= 1.24 \cdot 10^{-3} \quad (\pm 13.4\%) \\ \Delta\mathcal{B}(P^+ < 0.66\,\text{GeV}\ ,\ E_l > 1\,\text{GeV}) &= 1.10 \cdot 10^{-3} \quad (\pm 17.2\%) \end{aligned}$$

and the computed event fraction

$$\begin{aligned} R_{\text{cut}}(P^+P^- < (1.7\,\text{GeV})^2\ ,\ E_l > 1\,\text{GeV}) &= 0.615 \qquad (\pm 9.6\%) \\ R_{\text{cut}}(P^+ < 0.66\,\text{GeV}\ ,\ E_l > 1\,\text{GeV}) &= 0.535 \qquad (\pm 15.2\%), \end{aligned}$$

we obtain

$$\begin{aligned} |V_{ub}| &= \left(4.35 \pm 0.28_{[\text{exp}]} \pm 0.14_{[\text{th}-\text{total}\,(m_b^{\overline{\text{MS}}})]} \pm 0.22_{[\text{th}-\text{cuts}]}\right) \cdot 10^{-3} \\ |V_{ub}| &= \left(4.39 \pm 0.36_{[\text{exp}]} \pm 0.14_{[\text{th}-\text{total}\,(m_b^{\overline{\text{MS}}})]} \pm 0.38_{[\text{th}-\text{cuts}]}\right) \cdot 10^{-3} \end{aligned}$$

Conclusions

- Resummed perturbation theory can be directly used as can approximation to inclusive B meson decay spectra, without a leading power non-perturbative function.
- The leading renormalon cancels out when using the **same prescription** for renormalons in both the Sudakov exponent and the pole mass that appears upon conversion to physical variables.
- Beyond the logarithmic accuracy at hand (NNLL), the Borel sum of the exponent is constrained by information on renormalon residues. For the quark distribution in an on-shell heavy quark $B_{\mathcal{S}}(u = 1/2)$ was **computed**(!) and $B_{\mathcal{S}}(u = 1)$ **vanishes**(?)
- Contrary to Sudakov resummation with fixed logarithmic accuracy,
  the DGE prediction is free of Landau singularities and stable.
- The DGE spectrum smoothly extends beyond the perturbative endpoint
  Its support is close to the physical one, provided that $B_{\mathcal{S}}(u)$ is not too large at intermediate u (e.g. around $u \sim 3/2$).
- Application to charmless semileptonic decay*:
  the event fraction associated with an invariant mass cut $P^+P^- < (1.7\,\mathrm{GeV})^2$
  is obtained with $\pm 10\%$ accuracy
  Consistent values for $|V_{ub}|$ are obtained from two different cuts.

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*The program can be found at: `www.hep.phy.cam.ac.uk/~andersen/BDK/B2U`